2008
ELECTROWEAK INTERACTIONS
AND UNIFIED THEORIES
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XLIIIrd Rencontres de Moriond
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Electroweak Interactions and Unified Theories

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The XLIIIrd Rencontres de Moriond

**ELECTROWEAK INTERACTIONS AND UNIFIED THEORIES**

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The XLIIIrd Rencontres de Moriond were held in La Thuile, Vallée d'Aoste, Italie.

The first meeting took place at Moriond in the French Alps in 1966. There, experimental as well as theoretical physicists not only shared their scientific preoccupations, but also the household chores. The participants in the first meeting were mainly French physicists interested in electromagnetic interactions. In subsequent years, a session on high energy strong interactions was added.

The main purpose of these meeting is to discuss recent developments in contemporary physics and also to promote effective collaboration between experimentalists and theorists in the field of elementary particle physics. By bringing together a relatively small number of participants, the meetings help to develop better human relations as well as a more thorough and detailed discussion of the contributions.

Our wish to develop and to experiment with new channels of communication and dialogue, which was the driving force behind the original Moriond meetings, led us to organize a parallel meeting of biologists on Cell Differentiation (1980) and to create the Moriond Astrophysics Meeting (1981). In the same spirit, we started a new series on Condensed Matter Physics in January 1994. Meetings between biologists, astrophysicists, condensed matter physicists and high-energy physicists are organized to study how the progress in one field can lead to new developments in the others. I trust that these conferences and lively discussions will lead to new analytical methods and new mathematical languages.

The XLIIIrd Rencontres de Moriond in 2008 comprised three physics sessions, one nanophysics session and one Astrophysics session:

* March 1 - 8  "Electroweak Interactions and Unified Theories"
* March 1 - 8  "Venus Express Science Workshop"
* March 8 - 15  "QCD and High Energy Hadronic Interactions"
* March 8 - 15  "Quantum electronic transport and Nanophysics"
* March 15- 22  "Cosmology"
I thank the organizers of the XLIIIrd Rencontres de Moriond:


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I am also grateful to Enrico Belli, Andrea Righetto, Marcello Torre, Gioacchino Romani and the Planibel Hotel staff who contributed through their hospitality and cooperation to the well-being of the participants, enabling them to work in a relaxed atmosphere.

These Rencontres were sponsored by the European Union "Marie Curie Conferences and Training Courses" Activity, the Centre National de la Recherche Scientifique, the "Centre National d’ Études Spatiales", the Centres de Compétence en Nanoscience, the European Space Agency, the Institut National de Physique Nucléaire et de Physique des Particules (IN2P3-CNRS), the Commissariat à l’Energie Atomique (DSM and IRFU), the Fonds de la Recherche Scientifique (FRS-FNRS), Belgian Science Policy and the National Science Foundation. I would like to express my thanks for their encouraging support.

It is my sincere hope that a fruitful exchange and an efficient collaboration between the physicists, the astrophysicists and the nanophysicists will arise from these Rencontres as from the previous ones.

J. Trân Thanh Vân
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Electroweak Interactions and Unified Theories

I - Standard Model, precision Electroweak data at colliders, Search for the Brout-Englert-Higgs particle and theoretical variations
Standard Model Higgs Searches at the Tevatron (Low Mass)

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We present recent results of the Standard Model (SM) Higgs boson searches from the Collider Detector at Fermilab (CDF) and DØ experiments using 1.0-2.1 fb$^{-1}$ data from proton-antiproton collisions with $\sqrt{s} = 1.96$ TeV at the Fermilab Tevatron. For Higgs mass ($m_H$) less than 135 GeV/c$^2$ (so-called low mass region), $W^\pm H \to t^\pm tH$, $ZH \to t\bar{t}H$, and $ZH \to \nu\bar{\nu}H$ processes are the most sensitive channels and these results are summarized in this report. In addition, a new CDF search based on $H \to \tau^+\tau^+ + 2\text{jets}$ mode is described. Both CDF and DØ have found no evidence for the Higgs boson, and therefore set upper limits on the Higgs production cross section. The latest Tevatron combined limits are also presented.

1 Introduction

The Higgs boson is an essential component in the Standard Model (SM), which provides masses to fundamental particles through electroweak symmetry breaking, but yet undiscovered as of now. Theory indicates that the Higgs should be a scalar and charge-neutral boson but no predictions on its mass. However, combining the results from the direct searches at LEP and the SM global fits based on electroweak data, the mass of the SM Higgs boson is bounded in the range of 114.4 to 144 GeV/c$^2$ at 95% confidence limit (CL)$^1$.

The CDF and DØ experiments at the Tevatron are currently the only places capable of searching for the Higgs boson. The dominant Higgs production mode at the Tevatron ($p\bar{p}$ collision with $\sqrt{s} = 1.96$ TeV) is via gluon fusion ($\sigma \approx 1$ pb) which is about five to even ten times larger than the associated production of $W^\pm H$ and $ZH$. The fourth highest production is given by so-called Vector Boson Fusion (VBF) process with slightly lower production rate than $ZH$. In the low mass region ($m_H < 135$ GeV/c$^2$), the Higgs immediately decays to $b\bar{b}$ at 70$\sim$80% or $\tau^+\tau^-$ at 7$\sim$8% respectively. Although $gg \to H \to b\bar{b}$ mode is expected to provide the largest signal yield, it is overwhelmed by QCD multijet background. Therefore the most
relevant channels at the Tevatron are $W^\pm H \to l^\pm \nu \bar{b}b$, $ZH \to l^+l^-\bar{b}b$, and $ZH \to \nu\bar{b}b\bar{b}$ as shown in figure 1.

Figure 1: Tree level Feynman diagrams for signal processes used for low mass Higgs searches. From left to right, $W^\pm H \to l^\pm \nu\bar{b}b$, $ZH \to l^+l^-\bar{b}b$, and $ZH \to \nu\bar{b}b\bar{b}$ final states.

In addition to these channels, CDF newly performed the SM Higgs boson search using $H \to \tau^+\tau^- + 2\text{jets}$ from $WH$, $ZH$, $VBF$, and gluon fusion processes. In this report, all results and combined limits are briefly summarized.

2 Recent RunII Results

Both CDF and DØ have established the search for the SM Higgs as one of their highest priorities and already presented elsewhere and published (e.g., in ²). This proceeding therefore mainly focuses on describing recent updates for analyzes with 1.0~2.1 fb⁻¹ of data, depending on channels.

For $H \to \bar{b}b$ associated with $W/Z$ boson, $b$ quark identification is one of the key elements in event selection. $b$ quarks produce jets in detection and can be identified by either secondary vertex algorithm (SecVtx btagging) or jet probability algorithm (JetProb btagging) based on the large $B$ hadron lifetime. The former algorithm looks for a displaced vertex from the primary vertex of an event and the latter requires a low probability that all tracks contained in a jet originated from the primary vertex, based on the track impact parameters. These btagging requirements are crucial to reduce large $W/Z + \text{jets}$ and QCD multijet backgrounds where non-$b$jet (gluon jet and $u-,d-,c-,s-$) jets contribution is dominated. However since the Higgs signal is still much smaller than the backgrounds even after btagging, all analyzes presented here employed multivariate techniques (Artificial Neural Net (ANN) or Boosted Decision Tree (BDT)) to further improve signal-to-background ratio.

Although significant improvements with enormous efforts have been achieved by both CDF and DØ, no single analysis by itself would reach the sensitivity needed for discovery or exclusion. Thus all channels studied must be combined.

2.1 Search for $W^\pm H \to l^\pm \nu\bar{b}b$

The event selection is imposed by having one high $p_T$ isolated lepton ($e$ or $\mu$), missing transverse energy ($E_T^\text{miss}$) from the neutrino (from $W$ decay), and two or more jets (exact two jets in CDF) expected from Higgs decay. CDF and DØ results presented here are based on 1.9 fb⁻¹ and 1.7 fb⁻¹ data respectively. From previous results, CDF updated the analysis to include forward electron with $\eta$ up to $\sim$2.0 and also added 1 $b$tag category with neural net flavor separation (ANN$_{btag}$) in addition to double btagged events³. Candidate events are therefore classified into three categories: 2 SecVtx, 1 SecVtx + 1 JetProb, and 1 ANN$_{btag}$. Those three are exclusive to each other and have different signal purity. Thus they are analyzed separately but combined to obtain the final sensitivity in the likelihood calculation. DØ uses events with not only exact
two jets but also three jets and performed sample splitting in a similar way of CDF (1 ANN tight bjet and 2 loose ANN bjets, depending on operation point with respect to fake rate and signal efficiency where they have been selected based on the optimal combined sensitivity to a $WH$ signal)\(^4\). Dominant backgrounds are $W + b\bar{b}$ and $t\bar{t}$ events. Both CDF and DØ use ANN technique for the final fit to distinguish signal from such backgrounds. Figure 2 shows ANN output distribution (CDF) and dijet mass distribution (DØ) for 2 btagged events. No significant excess is observed, so upper cross section limits at 95% CL are set. For $m_H = 115$ GeV/$c^2$, DØ sets a limit of 1.4 pb (1.2 pb expected) and CDF sets a limit of 1.1 pb (1.0 pb expected) while SM prediction on $\sigma(p\bar{p} \rightarrow WH) \times BR(H \rightarrow b\bar{b})$ is 0.13 pb.

![CDF Run II Preliminary (1.9 fb$^{-1}$)](image1)

![DØ Preliminary](image2)

Figure 2: Left: ANN distribution from CDF. Right: Dijet mass distribution from DØ. Both are for events with 2 btagged jets. Higgs signal (times 10) is also shown for comparison.

### 2.2 Search for $ZH \rightarrow t^+t^-b\bar{b}$

Since this analysis was done in Summer 2007 based on relatively small data (1.0-1.1 fb$^{-1}$) compared to other analyses summarized in this proceeding, only core parts of the analysis are briefly described here. More details can be found elsewhere\(^5\). The signature for this channel is two high $p_T$ isolated leptons and two high $E_T$ jets that could be identified as $b$ jets. The main backgrounds are $Z + jets$ with mis-btagging, $Z + bb$ and $t\bar{t}$. DØ uses one dimensional ANN with ten kinematic variables while CDF’s ANN analysis is two dimensions with eight variables to distinguish signal from $Z + jets$ and $t\bar{t}$ separately. CDF also employed ANN to improve the energy resolution of dijet invariant mass by correcting energies of two leading jets independently according to their projection onto the $E_T$ direction. This correction improves dijet mass resolution from 14% to 9% as shown in figure 3. No excess is found in signal region and so for $m_H = 115$ GeV/$c^2$, DC sets a limit of 1.4 pb (1.6 pb expected) and CDF sets a limit of 1.3 pb (1.3 pb expected) while SM prediction on $\sigma(p\bar{p} \rightarrow ZH) \times BR(H \rightarrow b\bar{b})$ is 0.8 pb.

### 2.3 Search for $ZH \rightarrow \nu\bar{\nu}b\bar{b}$

In this channel, there is no leptons in the final state and so the signature is two high $E_T$ jets which could be btagged and high $E_T$ due to the two neutrinos from $Z$ decay. The main backgrounds arise from $Z + jets$, $Z + bb$, $t\bar{t}$, and QCD. This analysis is challenging because except for $Z \rightarrow \nu\bar{\nu} + bb$, all backgrounds are basically fakes originated from either mis-btagged
jet or jet mis-reconstruction and/or jet energy mis-measurement that mimic bjet and large $E_T$, respectively. However there is a strong advantage that it gains some acceptance from $WH$ when the lepton is lost in detection. DØ uses 2.1 fb$^{-1}$ (the largest dataset in all analyzes presented in this report) and employs the Boosted Decision Tree (BDT) technique which is basically a machine learning technique to extend a simple cut-based analysis into a multivariate technique by continuing to analyze events that fail a particular criterion. The 95% CL limits are set with respect to the SM prediction as shown in figure 4. At 115 GeV/c$^2$, the limit is 7.5×SM (8.4×SM expected)$^b$. In CDF, two separate ANN have been developed for this analysis. The first focuses on discerning real $E_T$ from fake $E_T$ generated by mis-measurement by using track information. The second ANN is trained to optimize the separation of both $ZH$ and $WH$ events from QCD and $t\bar{t}$ backgrounds as shown in figure 4. Using 1.7 fb$^{-1}$ data, CDF found no excess. Therefore upper limits are set. The limit is 8.0 times over the SM prediction (8.3 expected) at $m_H = 115$ GeV/c$^2$.$^7$.

Figure 3: Left: Invariant mass distribution of two btagged jets (CDF). The red/yellow histograms is the invariant mass for the Higgs signal(times 50) after/before the ANN jet correction. Right: ANN output distribution (DØ).

Figure 4: Left: ANN output distribution for 2 btagged events (CDF). Right: Upper limits (relative to the SM prediction) set by DØ $ZH \rightarrow \nu\ell b\bar{b}$ analysis.
2.4 Search for $WH + ZH + VBF + ggH \rightarrow \tau^+\tau^- + 2jets$

All analyzes presented so far use $H \rightarrow b\bar{b}$ decay to take an advantage of the high branching ratio while it is smaller for $H \rightarrow \tau\tau$ (about 7-8%). CDF however performed the first dedicated search for the SM Higgs boson using $\tau$ decay. To recover disadvantage of $H$ decay rate, four signal processes are considered and simultaneously searched as shown in figure 5. The signature is one isolated lepton ($q_{lep}$) from leptonic $\tau$ decay, one hadronic $\tau$ ($q_{had}$) from another $\tau$, and two or more jets. No other cuts are imposed in the event selection but instead ANN is trained in order to incorporate possible kinematic variables and their correlations to maximize discrimination power against backgrounds. Main backgrounds are $Z \rightarrow \tau\tau + jets$ and jet faking $q_{had}$ (mainly from $W + jets$ and QCD). After all jet $\rightarrow q_{had}$ backgrounds are estimated and modeled by data-driven way (using the same sign ($Q(q_{lep}) \times Q(q_{had}) > 0$) data), three ANN are trained: 1) Mixed signal vs $Z \rightarrow \tau\tau$, 2) Mixed signal vs QCD, 3) Mixed signal vs $t\bar{t}$, where mixed signal is a mixture of four signal events. Then given a event, these three ANN scores are calculated and the minimum of the three is selected for making the final ANN distribution. With 2 fb$^{-1}$ CDF data, the simultaneous search was performed but there was no clear excess in signal region as shown in figure 5. Therefore 95% CL limits are set. The result at $m_H = 115$ GeV/$c^2$ is 30×SM (24 expected)$^9$.

![Diagram](image)

Figure 5: Left: Signal Feynman diagrams used in this analysis. Right: ANN distribution with 2.6 fb$^{-1}$ CDF data. The Higgs signal (times 3C) is shown in blue line.

3 Tevatron Combined Limits

Figure 6 shows the current Tevatron limits (at 95% CL) on the Higgs production cross section to the SM predictions. In this combination, all CDF and DØ results are included except for the new CDF's $H \rightarrow \tau^+\tau^- + 2jets$ result. Also for DØ ZH $\rightarrow \nu p\bar{p}$ analysis, only 0.9 fb$^{-1}$ data is currently combined. At $m_H = 115$ GeV/$c^2$, the observed limit is 6.2×SM (4.3 expected)$^{10}$.

4 Conclusions

We have reviewed new preliminary results for the SM Higgs boson searches performed by the CDF and DØ collaborations with 1.0-2.1 fb$^{-1}$ dataset accumulated from $p\bar{p}$ collisions at $\sqrt{s} = 1.96$ TeV. No signal is found in both experiments and the 95% CL limits are set for a wide range of Higgs mass. At $m_H = 115$ GeV/$c^2$, combined result now is only about 5 times far from the SM
prediction for the exclusion. Furthermore, there are some results not yet combined and also more improvements can be expected and already planned by adding new decay channels, improving b-tagging quality and using advanced technique and so on. On the other hand, CDF and DØ expect to collect at least 6 fb⁻¹ data by 2009. Therefore together with expected improvements, there might be a great chance to discover or exclude the SM Higgs even in this challenging low mass region by the end of Run II.

Acknowledgments

I would like to thank the members of the CDF and DØ collaborations for their work and effort in achieving the results presented in this report. Also I would like to thank the Moriond conference organizers for the invitation and great hospitality during the conference and the funding agencies for making this work possible.

References

Searches for non-SM Higgs Bosons at the Tevatron

Andrew Haas for the DZero and CDF collaborations
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Extensions of the Standard Model (SM) predict Higgs phenomenology which can be quite different from that expected within the SM. This contribution discusses the latest results from searches for Higgs bosons by the DZero and CDF experiments at the Tevatron in several non-SM scenarios: supersymmetry, left-right symmetric (Higgs-triplet), and fermiophobic.

1 Introduction

The Tevatron accelerator at Fermilab has performed very well in recent years, delivering over 3.5 fb$^{-1}$ of data. The DZero and CDF experiments are using this data to study many questions at the forefront of high-energy physics. A central goal is the understanding of electro-weak symmetry breaking, which gives the W and Z bosons mass while leaving the photon massless. In the SM, electro-weak symmetry breaking is delivered via the Higgs mechanism, which predicts a single neutral spin-0 boson - the Higgs boson. Extensions to the SM, such as supersymmetry, predict different productions and/or decays of the Higgs boson and often a richer spectrum of multiple Higgs bosons.

We will discuss first the most anticipated extension to the SM, supersymmetry. Both DZero and CDF have specifically tuned searches for the new neutral Higgs boson(s) of supersymmetry, in the $bh \rightarrow bbb$, $h \rightarrow \tau\tau$, and $bh \rightarrow b\tau\tau$ channels. Next we show the results for left-right symmetric or Higgs-triplet models, which predict a doubly-charged Higgs boson, $H^{++}$. Finally, DZero has considered the case where the Higgs prefers not to decay to fermions, the so-called fermiophobic scenario, leaving open the $h \rightarrow \gamma\gamma$ channel as the main decay mode for a low-mass Higgs boson.
2 Searches in the MSSM at high $\tan\beta$

In two-Higgs-doublet models of electro-weak symmetry breaking, such as the minimal supersymmetric extension of the standard model (MSSM)\textsuperscript{1}, there are five physical Higgs bosons: two neutral $CP$-even scalars, $h$ and $H$, with $H$ being the heavier state; a neutral $CP$-odd state, $A$; and two charged states, $H^\pm$. The ratio of the vacuum expectation values of the two Higgs fields is defined as $\tan\beta = v_u/v_d$, where $v_u$ and $v_d$ refer to the fields that couple to the up-type and down-type fermions, respectively. At tree level, the coupling of the $A$ boson to down-type quarks, such as the $b$-quark, is enhanced by a factor of $\tan\beta$ relative to the SM, and the production cross section is therefore enhanced by $\tan^2\beta^2$. At large $\tan\beta$, this is also true either for the $h$ or $H$ boson depending on their mass. At high $\tan\beta$, the $h/H$ and $A$ decay roughly 90% to $bb$ and 10% to $\tau\tau$.

The dominant decay to $bb$ is unfortunately drowned by QCD background, but due to the enhanced coupling to $b$-quarks the $h/H$ and $A$ are also produced in association with one or more $b$-quarks, opening the channel $bh \rightarrow bbb$. Both DZero and CDF have performed searches for an excess in the di-jet invariant mass spectrum of the $bbb$ final state. CDF’s latest results are shown in Figure 1, using $1.9$ fb\textsuperscript{-1} of data. The dijet mass spectrum of the heavy flavor multi-jet background is derived from double-tagged data in a manner that accounts for tagging biases and kinematic differences introduced by the addition of the third tag. No excess is observed for any di-jet invariant mass window, so limits are placed in the Higgs mass vs. $\tan\beta$ parameter space.

The subdominant decay of the $h/H$ or $A$ to $\tau\tau$ is much cleaner, so both DZero and CDF have searched for direct $h \rightarrow \tau\tau$ excesses, see Figures 2 and 3. The main background is from $Z \rightarrow \tau\tau$, which is essentially irreducible.

At DZero, a set of Neural Networks (NN) are trained to identify tau decays from jet backgrounds, for each of 3 tau types (charged pion-like, pion + EM shower-like, and 3-prong). One of the taus is required to decay to a muon, for triggering and to reduce QCD background. The QCD background is determined by comparing same-sign vs. opposite-sign candidates. Some loose selection cuts remove W backgrounds, such as requiring the visible W mass $< 20$ GeV. Finally, a set of NN’s (one for each tau type) is used to separate signal from backgrounds. Good agreement is seen between data and expected background at high NN output, so limits are placed on the signals’ cross-sections and interpreted in the Higgs mass vs. $\tan\beta$ parameter space. CDF has performed similar analyses, but using in addition the $e+\tau$ and $e+\mu$ decay channels, as well as more data. No NN separation is employed, however.

By adding the requirement that there be an associated $b$-quark in the production, followed by the clean di-tau decay, $bh \rightarrow b\tau\tau$ has the highest signal / background of any high $\tan\beta$ MSSM
Figure 2: (top) NN outputs for all tau types (top left) and each tau type (see text) individually in the DZero \( h \rightarrow \tau \tau \) analysis. (bottom) Visible mass distributions in the CDF \( h \rightarrow \tau \tau \) analysis for the \( e\mu \) and lepton + \( \tau \) channels.

Figure 3: D0 (left) and CDF (right) limits in the \( m_A - \tan \beta \) parameter space of the MSSM from the \( h \rightarrow \tau \tau \) analyses.
Figure 4: Limits in the $m_A - \tan\beta$ parameter space of the MSSM from the DZero $bh \rightarrow b\tau\tau$ analysis.

Figure 5: (left) Di-muon invariant mass spectrum in events with 3 muons. (right) Limits on the $H^{++}$ mass for left- or right-handed Higgs bosons from the latest DZero search.

Higgs search. It also suffers from a low cross-section times branching ratio however. DZero has performed a search using just 344 $pb^{-1}$ of data in this channel. Di-tau events are normalized to the $Z \rightarrow \tau\tau$ peak, and then an additional b-tagged jet is required with $p_T >15$ GeV. For a Higgs mass of 120 GeV and $\tan\beta=80$, there are 5.3 signal events expected, with just 6.3 expected background events. Only 3 events are observed in data, and limits are placed in the Higgs mass vs. $\tan\beta$ plane, as seen in Figure 4, competitive with the other channels.

3 Searches for $H^{++}H^{--}$

Many models of beyond-SM Higgs physics predict a doubly-charged Higgs boson, $H^{++}$, such as left-right symmetric models, Higgs-triplet models, and little-Higgs. DZero has recently updated its search for pair-produced $H^{++}H^{--}$, using 1.1 $fb^{-1}$ of data. 3 muons are required, with $p_T >15$ GeV. At least one pair must have an invariant mass $>30$ GeV and $\Delta\phi <2.5$ radians, to reduce backgrounds from QCD and Z decays. 3 events are observed in data, for an expectation of 3.1 events from backgrounds, and thus limits are set on the $H^{++}$ mass as shown in Figure 5. Left-handed doubly-charged Higgs bosons, $H_L^{++}$ have a larger production cross-section than right-handed, $H_R^{++}$, by about a factor of 2, due to their different coupling to the intermediate Z boson. Thus, the limit on the mass of $H_L^{++}$ ($>150$ GeV) are higher than on $H_R^{++}$ ($>127$ GeV).
4 Searches for $h \rightarrow \gamma \gamma$

Although the SM Higgs boson decays to the di-photon final state with a small branching ratio, this decay channel could be subject to a large enhancement in models where the Higgs does not couple strongly to fermions. In this case, fermion masses could arise from some other source. DZero has recently completed a search using $2.3 \, fb^{-1}$ of data for $h \rightarrow \gamma \gamma$. 2 isolated photons with $p_T > 25$ GeV are required. QCD jets faking photons are estimated by looking at the calorimeter shower-shape correlations between the two photons in each event, and samples are divided into QCD di-jet, photon+jet, and di-photon backgrounds. $Z \rightarrow ee$ also contributes near the Z mass. Data agree very well with estimated backgrounds, and no significant excess is seen at any di-photon invariant mass range. Limits are therefore set on the cross-section times branching ratio for $h \rightarrow \gamma \gamma$ and are compared to the expected cross-section times branching ratio in the SM, as shown in Figure 6. The current analysis can exclude down to about 50 times the SM event rate at 120 GeV and in fact also contributes non-negligibly to the SM Higgs search.

5 Conclusions

CDF and DZero have searched for Higgs bosons in several models beyond the SM: the MSSM at high $\tan\beta$ in b and $\tau$ channels, for doubly charged Higgs, and for fermiophobic Higgs. No significant excesses or deviations indicative of non-SM Higgs signatures have been seen so far. But the Tevatron experiments expect to more than double their data samples in the next two years, hone their search strategies, and enlarge the variety of models considered.

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References

W/Z Production: Asymmetry, Z(p_T), and W+charm at the Tevatron

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We report on W/Z production processes at the Tevatron, including the W charge asymmetry, the Z boson rapidity and p_T distributions, and the W+charm final state. The measurements test the Standard Model and yield constraints on parton distribution functions. The data samples used range from 0.3 to 2.1 fb^{-1} of pp collisions at \sqrt{s}=1.96 TeV.

1 W/Z production at the Tevatron

1.1 Electroweak Physics at Tevatron

Investigations of electroweak processes at the Tevatron have focussed on three different areas: (a) the W mass and width, (b) diboson production, and (c) details of W/Z production processes. Precise measurements of the W mass constrain the mass of the Higgs boson, and measurements of the W width can be compared with Standard Model (SM) predictions. Diboson production provides information on gauge boson self-interactions and is sensitive to physics beyond the SM. Understanding diboson production is also relevant to new particle searches because it is the background to Higgs and SUSY particle production. Measurement of details of W/Z production processes, such Z boson rapidity and p_T distributions and the W charge asymmetry provide information on the up and down parton distribution functions (PDFs) in the proton, while the W+charm cross section is sensitive to the PDF of strange quarks.

1.2 Detectors

We report on results from CDF and D0 Tevatron experiments at Fermilab. The CDF detector has a larger tracking volume which yields better momentum resolution for charged tracks than the D0 detector in the central region. The D0 detector has a larger muon acceptance and more coverage in the forward rapidity region.1,2
1.3 W Boson and Lepton Charge Asymmetry

$W^+(W^-)$ events are produced in the collisions of $u(\bar{u})$ and $\bar{d}(d)$ quarks. Since, on average, the $u(\bar{u})$ quark carries a higher momentum fraction than the $\bar{d}(d)$ quark, the $W^+(W^-)$ is boosted along the proton (anti-proton) direction. This results in a $W$ charge asymmetry. The $W$ charge asymmetry ($A(y_W)$) is defined as:

$$A(y_W) = \frac{d\sigma(W^+)/dy_W - d\sigma(W^-)/dy_W}{d\sigma(W^+)/dy_W + d\sigma(W^-)/dy_W}$$  \hspace{1cm} (1)$$

At LO, it can be described in terms of the momentum fractions of the $u$ and $d$ quarks:

$$A(y_W) \approx \frac{d(x_2)/u(x_2) - d(x_1)/u(x_1)}{d(x_2)/u(x_2) + d(x_1)/u(x_1)}$$  \hspace{1cm} (2)$$

where $x_{1,2} = \frac{M_W^2}{s} \epsilon^{\mp y_W}$. The $A(y_W)$ is sensitive to slope of $d(x)/u(x)$.

Experimentally, the lepton charge asymmetry is measured:

$$A(\eta) = \frac{d\sigma(\ell^+)/d\eta - d\sigma(\ell^-)/d\eta}{d\sigma(\ell^+)/d\eta + d\sigma(\ell^-)/d\eta}$$  \hspace{1cm} (3)$$

The measured lepton charge asymmetry is a convolution of the $W$ charge asymmetry, and the asymmetry from the V-A interaction. Figure 1 shows the $W$ boson and the lepton rapidity distributions and the resulting $W$ and lepton charge asymmetries.

At large lepton pseudorapidities, the V-A interaction distorts the boson production asymmetry. DØ and CDF collaboration use different approaches in the measurement of the $W$ asymmetry. The DØ collaboration uses the traditional method of measuring the charge asymmetry in the muon decay channel. The data sample consists of $0.3 \text{ fb}^{-1}$. Figure 2 shows the lepton charge asymmetry versus muon pseudorapidity. The central region measurements by both CDF and DØ ($|\eta_\mu| < 1.3$) have been used to constrain PDFs in previous fits. The high rapidity region ($|\eta_\mu| > 1.3$) probes the PDFs at higher $x$. The high rapidity region is still statistics limited.

A new measurement by the CDF collaboration extracts the $W$ charge asymmetry from the lepton charge asymmetry. The $W$ boson rapidity ($y_W$) can be determined from the 4-momentum of the final state neutrino. However, the momentum of the neutrino in z-direction ($P_\nu^z$) is not measured and only its transverse momentum is known (the missing transverse energy of the event). The $P_\nu^z$ is calculated as follows:

$$M_{W}^2 = (E_\ell + E_\nu)^2 - (P_\ell + P_\nu)^2$$  \hspace{1cm} (4)$$
Figure 2: The left plot shows the lepton charge asymmetry in muon pseudorapidity measured in DØ, and the right plot shows the W boson charge asymmetry in boson rapidity measured in CDF.

Figure 3: Z boson rapidity distributions. The left plot is the normalized of Z boson rapidity distribution measured at DØ, and the right plot is the Z boson rapidity measured at CDF. The theory prediction is normalized to the measured total $\sigma$.

where $M_W$ is constrained to the W boson mass. There are two solutions for $P_\lambda^W$. Each $P_\lambda^W$ solution is weighted by a factor which includes the $W^\pm$ production cross section ($d\sigma/dy_W$), the V-A angular distribution function in the center-of-mass frame ($1-\cos^2\theta^*$), and a $W^\pm$ transverse momentum ($P_T^W$) factor to account for higher order QCD corrections to the V-A $\cos\theta^*$ term. The dependence on input $d\sigma/dy_W$ is removed by iterating the measurement. Each iteration is used further to constrain $d\sigma/dy_W$. Here, the larger acceptance electron sample with 1 fb$^{-1}$ is used.

The W charge asymmetry for 1 fb$^{-1}$ data is shown in Figure 2. There is good agreement with the NNLO prediction using MRST2002 PDFs. The experimental error is smaller than the current uncertainty in the PDFs, and the new data can be used to further constrain PDFs in future fits.

1.4 The Z Boson Rapidity Distribution: $d\sigma/dy$

In the Drell-Yan process, the quark and anti-quark carry parton momentum fractions, $x_1$ and $x_2$. The difference in the parton momentum fractions $(x_{1,2})$ determine the rapidity of the final state Z boson $(y_Z)$. Since high $y_Z$ corresponds to high $x_1$ and low $x_2$, this region probes PDFs at high x. In addition, the $d\sigma/dy$ distribution tests QCD theory predictions at higher orders. In higher order (NLO or NNLO), gluons splitting in the initial state also contribute to the rapidity distribution. Both DØ and CDF measure $d\sigma/dy$ using the dielectron final state. In both measurements, $y_Z$ measurements extend up to 2.9. The DØ measurement, shown in Figure 3, with 0.4 fb$^{-1}$ shows good agreement with the NNLO MRST prediction. The CDF measurement, also shown in Figure 3, uses 2.1 fb$^{-1}$. Figure 4 shows the ratio of the CDF data to the NLO calculation with NLO CTEQ6.1M PDFs.
Figure 4: Ratio of data and theory prediction for the $Z$ boson rapidity distribution in CDF. Here, the NLO calculation with NLO CTEQ6.1M PDFs is used. The theory prediction is normalized to the measured $\sigma$.

Figure 5: The left plot shows the normalized differential cross section of $Z$ boson $p_T$ up to $p_T < 300\text{GeV}/c$. The right plot shows the fractional difference between data and the theory predictions. "KF" in the scale factor(K-factor) of NLO to NNLO theory calculation.

1.5 $Z(p_T)$ Distribution

The $Z$ boson $p_T$ which can be measured over a wide range of values provide a test of QCD. In the $p_T > 30\text{GeV}/c$ region, where the $p_T$ originates from the radiation of energetic gluons, the perturbative QCD calculation\cite{NNLO} gives reliable predictions. In the $p_T < 30\text{GeV}/c$ region, where multiple soft gluon emission is dominant, the soft gluon resummation technique is used (ResBos)\cite{ResBos1,ResBos2}. The $D_{\bar{\phi}}$ measurement in the $Z$ with 0.98 $fb^{-1}$ of data in the $ee$ channel is shown in Figure 5. The ResBos curve is the gluon resummation calculation including PHOTOS\cite{PHOTOS}, which accounts for radiated photons in the final state. The "Rescaled NNLO" is the NNLO calculation rescaled to match the data at $p_T = 30\text{GeV}/c$. The right hand side of Figure 5 shows the fractional differences between the data and theory on a linear scale. In the $p_T < 30\text{GeV}/c$ region, the ResBos calculation describes the data well. For $p_T > 30\text{GeV}/c$, the data is higher than all predictions. The NNLO theory prediction agrees with the data only in shape.

$D_{\phi}$ also investigated the “Small-x broadening effect”\cite{Small-x1,Small-x2} that predicts a wider $p_T$ distribution in the large rapidity region. This effect modifies the resummation form factor in the small-x parton region. At the Tevatron, $Z$ bosons with $2 < |y| < 3$ probe the Bjorken x region, $0.002 < |x| < 0.006$. Therefore, a measurement of the $Z$ boson $p_T$ in the high rapidity region is sensitive to the modified form factor. Figure 6 shows the $p_T < 30\text{GeV}/c$ region for all $y_Z$ and for $|y_Z| > 2$. The standard ResBos calculation agrees well with the data in all $y_Z$. The $|y_Z| > 2$ region data does not favor an additional small-x form factor. This is the first test of the “small-x broadening effect” using the high $y$ region at the Tevatron.
Figure 6: The plots show the normalized differential cross section of $Z$ boson $p_T$ measured at DØ. The left plot shows the result for all boson rapidity, and the right plot shows the result in the high rapidity region ($|y_Z| > 2$), which is sensitive to the “small-x broadening effect”.

Figure 7: The cross section ratio $W+c$-jet to $W+jets$ measured at DØ.

1.6 W+Charm Cross Section

The $W$+charm process, $g + s \to W^- + c$, is a background to top, SUSY particle, and SM Higgs production. Since CKM matrix element, $|V_{cd}|^2$, suppresses the $d$-quark- gluon fusion production, the $W$+charm production probes the $s$-quark PDF. Strange quarks contribute to the processes, $pp/\bar{p}p \to sg \to W^- + c$ and $pp/\bar{p}p \to \bar{s}c \to H^-$, at the Tevatron and LHC. At hadron colliders, the $s$-quark distribution is probed at large $Q^2 \approx M_W^2$.

The DØ and CDF collaborations tag the $c$-jet by looking for the muon from the decay of the charm particle in the $c$-jet. The muon from the decay of the charm particle and the lepton from the $W$ boson decay have the opposite electric charge. The opposite charge requirement is used in the event selection. All leptonic channels are allowed for the $W$ boson decay. $Z \to \mu\mu$ events are rejected by requiring $M_{\mu\mu} < 70\text{GeV}/c^2$ for the muon decay channel. DØ measures the cross section ratio of the $W$+c-jet to $W$+jets processes, $\frac{\sigma(pp \to W+c+jet)}{\sigma(pp \to W+jets)}$, using 1 $fb^{-1}$ data. The fraction of $\sigma(W+c-jet)$ compared with the theory versus jet $pt$ is shown in Figure 7. The fraction of $\sigma(W+c-jet)$ for $p_T > 20\text{GeV}/c$ is $0.071 \pm 0.017$. The measured fraction in the electron channel is $0.060 \pm 0.021(stat.) \pm 0.002(sys.)$ and the fraction in the muon channel is $0.093 \pm 0.029(stat.) \pm 0.005(sys.)$.

CDF measures the total cross section of $W+c$-jet with 1.8 $fb^{-1}$ data. Both electron and muon channels are used for the $W$ boson selection. The $c$-jets with $p_T(c) > 20\text{GeV}/c$ and $|\eta(c)| < 1.5$ are identified using the semi-muonic decay of the charm particle in the jet. The measured cross section for the leptonic channel, $\sigma_W \times BR(W \to l\nu)$, is $9.8 \pm 2.8(stat.) \pm 1.4(sys.) \pm 0.6(lum) pb$. 
which agrees with the NLO calculation, $11.0^{+1.4}_{-3.0}\text{pb}$.

1.7 Conclusion

New high statistics measurements of W and Z production processes at the Tevatron are used to provide new constraints on nucleon PDFs. The W charge asymmetry and Z boson $d\sigma/dy$ measurement have been extended up to $y \approx 2.9$. The Z boson $d\sigma/dy$ measurement is with 2.1 $fb^{-1}$ of data, and 8 $fb^{-1}$ is expected by end of 2009.

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References

ELECTROWEAK MEASUREMENTS FROM HERA

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New preliminary electroweak results from the HERA lepton-proton collider experiments H1 and ZEUS are presented. These include new high \( Q^2 \) neutral current cross section measurements, limits on a possible quark radius in the search for contact interactions as well as on quark-Z coupling parameters, extracted in combined electroweak and QCD fits. Furthermore, new charged current cross section measurements as a function of the lepton-beam polarisation are presented, as well as charged current measurement results, using the combined HERA I data from H1 and ZEUS. Finally, measurements of the single \( W \) boson production cross section and the \( W \) boson polarisation fractions are presented.

1 Introduction

The lepton-proton collider HERA\(^1\) has facilitated measurements of electroweak (EW) interactions between quarks and leptons in deep inelastic scattering (DIS) at a centre of mass energy up to 320 GeV and four momentum of the exchanged boson squared \( (Q^2) \) up to 40000 GeV\(^2\). The data taking took place in the years 1994-2000 (HERA I) and 2003-2007 (HERA II). Two collider-mode detectors, H1\(^2\) and ZEUS\(^3\), have each collected approximately 0.5 fb\(^{-1}\) of data, divided in electron-proton \((e^-p)\) and positron-proton \((e^+p)\) data. Previously obtained HERA results have led to a significantly improved understanding of the proton substructure\(^4,5\) allowing for measurements of important EW parameters\(^6,7,8,9\) as well as searches\(^10,11\) for contact interactions. In these proceedings, updates are presented to these analyses along with new results from H1 regarding the single production of \( W \) bosons at HERA.

2 Neutral Current Cross Section Measurement

New H1 Neutral Current (NC) single differential cross section measurements\(^12\) as a function of \( Q^2 \), are presented in Figure 1. The analysis includes previously published\(^7,13,14\) HERA I and preliminary\(^15\) HERA II data, corresponding to an integrated luminosity of 270 pb\(^{-1}\) of \( e^+p \) data and 165 pb\(^{-1}\) of \( e^-p \) data. The measurement precision is better than 10\% for \( Q^2 \) up to 20000 GeV\(^2\). The data agree well with SM QCD expectations, which are based on parton distribution functions obtained using high energy HERA I data.\(^4\)
3 Derivation of Limits on the Quark Radius in Contact Interactions

In the search for contact interactions, both H1 and ZEUS derive limits on a possible quark radius using high $Q^2$ NC events.\(^{12,16}\) A form factor $f_Q$, as a function of $Q^2$ and a hypothetical quark radius $R_q$, is defined as $f_Q(Q^2, R_q) = 1 - \frac{1}{6} \langle R_q^2 \rangle Q^2$. This leads to an altered single differential cross section $d\sigma/dQ^2 = f_Q(Q^2, R_q) d\sigma_{SM}/dQ^2$ for contact interactions, which can be fit to the data, as shown in Figure 2. The 95% Confidence Level (CL) limits on the quark radius are determined from the fit to be $0.74 \times 10^{-18}$ m and $0.62 \times 10^{-18}$ m, by H1 and ZEUS, respectively. These limits confine a hypothetical quark radius to be less than or equal to the resolution power of HERA.

Figure 2: NC single differential cross sections as a function of $Q^2$ normalised to the SM expectation $d\sigma_{SM}/dQ^2$. The lines represent corrections due to the hypothetical quark radius $R_q$ at its 95% CL limit for H1 $e^+p$ data (top left), H1 $e^-p$ data (bottom left), and for ZEUS $e^+p$ (right).
Figure 3: Reduced CC cross sections (points), using the combined HERA I data of H1 and ZEUS, in bins of $Q^2$ for $e^- p$ (left) and $e^+ p$ (right) data. The curves are NLO QCD fits, performed by H1 and ZEUS to their own data.

4 Combination of H1 and ZEUS Charged Current HERA I Data

The H1 and ZEUS collaborations are combining their data to improve the precision of DIS measurements. New preliminary results, using combined HERA I data of H1 and ZEUS, are shown in Figure 3, depicting CC reduced cross sections in bins of $Q^2$. A good agreement is observed between the combined data and the QCD fits of each experiment to their own data. The $e^+ p$ data correspond to a total integrated luminosity of approximately $200 \, \text{pb}^{-1}$ and have a typical precision of 8%. The statistical gain in precision is most significant in the statistically limited $e^- p$ data set, where the combined data ($30 \, \text{pb}^{-1}$) leads to an increase of the precision to about 20%. The precision is expected to increase further with the future inclusion of the HERA II data.

5 Charged Current Cross Section Measurement using Polarised Lepton-Beams

During the HERA II running period, longitudinally polarised lepton-beams were used. The polarisation is defined as $P = \frac{N_R - N_L}{N_R + N_L}$, with $N_R$ ($N_L$) the number of right (left) handed leptons in the beam. The SM predicts a linear scaling of the CC cross section with the beam polarisation, due to the absence of a right handed neutrino. New ZEUS results, using data from the years 2006-2007, are shown in Figure 4, together with previously obtained HERA measurements. A good agreement between the measurements in the different data sets and the SM expectations is observed.

6 Measurement of the quark-$Z$ coupling

The cross section of NC DIS events composes of photon ($\gamma$) and $Z$ exchange diagrams. Due to the heavy $Z$ boson propagator, the contribution from pure $Z$ exchange is suppressed. The contribution from $\gamma Z$ interference, however, is still sensitive to the vector and axial-vector couplings of the $Z$ boson to the quark. Limits at 68% CL on these couplings are extracted, using combined EW and QCD fits where the coupling parameters pertaining to both the up and down quark are left free in the fit. In particular, the limits concerning the couplings to the up quark ($v_u$ and $a_u$) profit from including the polarised HERA II data. The results are shown in
7 Measurement of Single $W$ Boson Production

Single $W$ boson production in the SM is a rare process at HERA with a cross section of order 1 pb. In the case of leptonic $W$ boson decay, for which the branching ratio is about 30%, the event gives rise to a characteristic $\ell + p_T$ detector signature, consisting of an energetic isolated electron or muon ($\ell$) and large missing transverse momentum ($p_T^{miss}$). The full HERA I+II high energy data sample, collected with the H1 detector in the years 1994-2007 and corresponding to an integrated luminosity of 478 pb$^{-1}$, is analysed and 59 $\ell + p_T^{miss}$ events are selected compared to a SM expectation of 58.9 ± 8.2. This yield is presented in Figure 6 (left) as a function of the transverse momentum of the hadronic system ($p_T^X$). Notwithstanding an excess of the data over the MC prediction in the small region of phase space where $p_T^X > 25$ GeV, a good over-all agreement with the SM is observed. The single $W$ boson production cross section is determined to be $\sigma_W = 1.2 \pm 0.3$ (stat) ± 0.2 (sys) pb, which is in good agreement with the (NLO) SM expectation of 1.3 ± 0.2 pb. The quoted errors on the measured cross section include theoretical and experimental uncertainties.

8 Measurement of the $W$ Boson Polarisation Fractions

The measurement of the $W$ boson polarisation fractions is based on the $\ell + p_T$ data sample discussed in Section 7 and makes use of the $\cos \theta^*$ distributions in the decay $W \to e/\mu + \nu$. $\theta^*$ is defined as the angle between the $W$ boson momentum in the lab frame and that of the charged decay lepton in the $W$ boson rest frame. For the left handed polarisation fraction $F_-$, the longitudinal fraction $F_0$ and the right handed fraction $F_+ \equiv 1 - F_- - F_0$, the $\cos \theta^*$
distributions for $W^+$ bosons are given\cite{20} by
\[ \frac{d\sigma_W}{d\cos\theta^*} \propto (1 - F_- - F_0) \cdot \frac{3}{8} (1 + \cos\theta^*)^2 + F_0 \cdot \frac{3}{4} (1 - \cos^2\theta^*) + F_- \cdot \frac{3}{8} (1 - \cos\theta^*)^2. \] (1)

For $W^-$ bosons, the $\cos\theta^*$ distributions have opposite values. To allow the combination of both channels, $\cos\theta^*$ is multiplied with the sign of the lepton charge $q_\ell = \pm 1$. Therefore, from the $\ell + P_T$ data sample, only events for which a reliable measurement of the charge of the isolated lepton exists are used. The reconstruction of the $W$ boson rest frame is performed and the $W$ boson differential cross section as a function of the decay angle $\theta^*$ is derived and fit to the model defined in Equation 1. In the fit, the optimal values for $F_-$ and $F_0$ are simultaneously extracted using a $\chi^2$ minimisation method. The result is shown in Figure 6 (right) and found to be in good agreement with the SM. $F_-$ and $F_0$ are also extracted in fits where one parameter is fixed to its SM value. No deviations from the SM are observed and the values are determined to be:

\[ F_- = 0.58 \pm 0.15 \text{ (stat)} \pm 0.12 \text{ (sys)} \quad \text{SM:} \quad 0.61 \pm 0.01 \text{ (stat)}, \]
\[ F_0 = 0.15 \pm 0.21 \text{ (stat)} \pm 0.09 \text{ (sys)} \quad \text{SM:} \quad 0.19 \pm 0.01 \text{ (stat)}. \]

9 Summary

Preliminary H1 results of NC cross section measurements at high $Q^2$ have been presented, using the complete HERA II high energy data set. A good agreement with the QCD SM expectations is observed. In the search for contact interactions, H1 and ZEUS derive strong upper limits on a possible quark radius of respectively $0.74 \cdot 10^{-18}$ m and $0.62 \cdot 10^{-18}$ m at 95% CL. In addition, two dimensional limits at 68% CL on the vector and axial vector couplings of the $Z$ boson to the up quark were shown, using combined EW and QCD fits. New measurements of the charged current cross section as a function of the lepton-beam polarisation, using data taken in the years 2006-2007, have been presented. A good agreement is observed with previous measurements and with the SM, which forbids right handed charged currents. Recently derived CC cross section measurements are presented, using the combined HERA I data of both experiments. The combination of the data has led to significant improvements in the statistical
precision. The cross section measurements using the combined data are in good agreement with the previously established QCD fits of H1 and ZEUS to their own data. A single W boson production cross section measurement is performed by H1 using the full HERA I+II data and found to be $\sigma_W = 1.2 \pm 0.3 \text{ (stat)} \pm 0.2 \text{ (sys) pb}$, which is in good agreement with the (NLO) SM expectation of $1.3 \pm 0.2 \text{ pb}$. Finally, the W boson polarisation fractions are measured and found to be in good agreement with the SM.

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ON THE POSSIBLE LINKS BETWEEN ELECTROWEAK SYMMETRY BREAKING AND DARK MATTER

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The mechanism behind electroweak symmetry breaking (EWSB) and the nature of dark matter (DM) are currently very important issues in particle physics. Usually, in most models, these two issues are not or poorly connected. However, since a natural dark matter candidate is a weakly interacting massive particle or WIMP, with mass around the electroweak scale, it is clearly of interest to investigate the possibility that DM and EWSB are closely related. In the context of a very simple extension of the Standard Model, the Inert Doublet Model, we show that dark matter could play a crucial role in the breaking of the electroweak symmetry. In this model, dark matter is the lightest component of an inert scalar doublet which can induce dynamically electroweak symmetry breaking at one loop level. Moreover, in a large fraction of the parameter space of this model, the mass of the dark matter particle is essentially determined by the electroweak scale, so that the fact that the WIMP DM mass is around the electroweak scale is not a coincidence.

1 Introduction

If one think about what kind of new physics the Large Hadron Collider could observe, beside elucidating the origin of electroweak symmetry breaking, which is the first goal of this accelerator, there are at least 2 issues which come directly to mind. The first one is the physics which would cure the hierarchy problem(s) related to the scalar sector of the theory. The second one is the particle at the origin of the dark matter in the universe. The reason why one might observe the DM particle at LHC is not as clear at all as for the physics at the origin of EWSB, but it is at least what we expect in the most straightforward explanation for the relic DM density of the universe, which is the WIMP mechanism. If the DM relic density of the universe is due to the

\*Talks given by M.T. at this Moriond Conference and by T.H. at the 4th Dark Side of the Universe Conference, June 2008, Cairo, Egypt, based on Ref. 1.
simple freeze out of the pair annihilation of a stable thermal particle, and if the annihilation cross section is driven by gauge couplings (or more generally couplings of order unity), the DM mass which is e.g. necessary to have the right relic density as observed in the universe \((\Omega_{DM} \simeq 0.22^{2,3})\) turns out to be around the electroweak scale. This leads to a coincidence problem since what sets the DM mass is the observed DM relic density which a priori has nothing to do with the electroweak scale. In most models it is a coincidence. In this talk we consider the following two questions curiously not often considered. First could it be not a coincidence due to some deep reason? Second, if the DM particle is around the electroweak scale, could it play a direct role in the dynamic of EWSB? In the following, focusing on these phenomenological issues of DM and EWSB, we consider an extremely simple model, the inert Higgs doublet model, which shows that DM could have indeed a crucial role in EWSB and that the WIMP scale coincidence above might not be accidental.

2 The inert Higgs doublet model

The model we consider is extremely simple.\(^4\)–\(^9\) It is based on only 2 assumptions. First it assumes the existence of a second Brout-Englert-Higgs (Higgs for short) doublet, \(H_2\). Second it assumes a discrete symmetry, the simplest one is a \(Z_2\) symmetry, such that all SM particles are even under it, except the second Higgs doublet. To assume such a discrete symmetry has several virtues. It automatically leads to no flavor changing neutral current problems which in more general 2 Higgs doublet model are generic. Moreover if the \(Z_2\) symmetry is not spontaneously broken, which is the case for large fractions of the scalar potential parameters, it leads to a stable DM candidate in the form of the lightest \(H_2\) component. The doublet \(H_2 \equiv (H^+ (H_0 + i A_0)/\sqrt{2})^T\), since it is complex, has four components, 2 charged, \(H^\pm\), one neutral scalar, \(H_0\), and one neutral pseudoscalar, \(A_0\).

The most general scalar potential one can write contains 2 mass and five quartic terms:

\[
V = \mu_1^2 |H_1|^2 + \mu_2^2 |H_2|^2 + \lambda_1 |H_1|^4 + \lambda_2 |H_2|^4 + \lambda_3 |H_1|^2 |H_2|^2 + \lambda_4 |H_1^\dagger H_2|^2 + \frac{\lambda_5}{2} (|H_1^\dagger H_2|^2 + h.c.)
\]

with real quartic couplings. After \(SU(2) \times U(1)\) symmetry breaking, from the vacuum expectation value of \(H_1\), \(\langle H_1 \rangle = v/\sqrt{2}\) with \(v = -\mu_1^2/\lambda_1 = 246\) GeV, we get the following mass spectrum

\[
\begin{align*}
  m_h^2 &= \mu_1^2 + 3\lambda_1 v^2 \equiv -2\mu_1^2 = 2\lambda_1 v^2 \\
  m_{H^+}^2 &= \mu_2^2 + \lambda_3 v^2/2 \\
  m_{H_0}^2 &= \mu_2^2 + (\lambda_3 + \lambda_4 + \lambda_5) v^2/2 \\
  m_{A_0}^2 &= \mu_2^2 + (\lambda_3 + \lambda_4 - \lambda_5) v^2/2.
\end{align*}
\]

with \(h\) the Higgs boson from \(H_1\). To have a dark stable particle, i.e. neutral, \(H_0\) or \(A_0\), we therefore need \(\lambda_2 \equiv \lambda_3 + \lambda_4 + \lambda_5 < \lambda_3\) and/or \(\lambda_5 \equiv \lambda_3 + \lambda_4 - \lambda_5 < \lambda_3\).

The DM properties of this model have been studied in a series of papers which show that this model is perfectly viable and moreover testable. There are 4 types of processes which drive the relic density\(^8\) annihilation to a pair of gauge bosons, to a pair of Higgs boson, to a pair of fermion via a Higgs boson, and coannihilation to a fermion pair of the DM particle with the other neutral \(H_2\) component via a \(Z\) boson or with \(H^\pm\) via a \(W^\pm\). The cross sections are exactly the same for \(H_0\) and \(A_0\) so that both DM candidates are equally good. Annihilations

---

\(^8\)For example in the MSSM neutralino scenario there is no direct link between these 2 scales, due to the \(\mu\) problem. There exists however models where such exist as in the NMSSM.
to a pair of gauge bosons tend to be too fast to give the right relic density, except in 2 mass regimes. A low mass regime where the DM mass is below the W and Z mass thresholds; it requires $40 \text{ GeV} < m_{DM} < 75 \text{ GeV}$ (in this case the relic density is determined by the 2 processes with light fermions in the final state). And a high mass regime (also possible because asymptotically, for large DM mass, the annihilation to a pair of gauge bosons drops as $1/m_{DM}^2$); it requires $600 \text{ GeV} \lesssim m_{DM} \lesssim 100 \text{ TeV}$.

For direct detection the main process is elastic scattering of DM with a nucleon via a Higgs boson. For the low mass regime most of the parameter space cannot be probed by present experiments but will be by the future ones, see $^6_{8,6}$. Similarly in this regime, and for usual Navarro-Frank-White DM galactic density profile, most of the parameter space will be covered by the GLAST satellite experiment, see $^8_{8,9}$. This model is therefore testable. The high mass regime, on the other hand, leads to more suppressed rates.

### 3 Electroweak Symmetry Breaking induced by Dark Matter

Although EWSB can be perfectly induced in the SM by the scalar potential of the Higgs boson without the need of any additional particle, the inert Higgs doublet model offers the possibility to have a dynamical origin for the EWSB. It provides an example of DM model where due to the fact that DM is around the electroweak scale, it can easily have an important role for EWSB, by driving it at one loop through the Coleman-Weinberg mechanism$^{10}_{10,11}$. Consider a regime where $\mu_1^2$ would be positive, vanishing or more generally much less negative than its ordinary value in the SM $-\lambda v^2$. In this case there is no or very little EWSB at tree level. There is not either EWSB at one loop in the SM. This is due to the well-known fact that the one loop effective potential is dominated in the SM by the top loops which have the wrong sign for EWSB. These loops can lead to a potential with an extremum in $v$ but only to a maximum, i.e. they destabilize the Higgs vacuum. However in the inert Higgs doublet model the situation is totally different. There are additional scalar loops involving $H_2$. Neglecting gauge bosons loops as well as fermion loops other than top ones, using the $\overline{MS}$ prescription, we get the following effective Higgs potential

$$V_{\text{eff}}(h) = \mu_1^2 h^2/2 + \lambda_1 h^4/4 + \frac{1}{64\pi^2} \sum_i n_i m_i^4 \left( m_i^2/m_{\lambda_i}^2 - 3/2 \right)$$

where $n_i = \{1, 1, 1, 1, 2, -2, -12\}$ is the number of degrees of freedom for each species $i = \{h, H_0, G_0, A_0, h^\pm, H^\pm, t\}$ which couples to the Higgs boson with tree level masses given in Eq. (2), $m_{G_0}^2 = m_{h^\pm}^2 = \mu_1^2 + \lambda_1 v^2$, $m_{H_0}^2 = g_t^2 v^2/2$ (with $G_0$, $G^\pm$, the 3 would-be Goldstone bosons in $H_1$ and $g_t$ the top Yukawa coupling). Since they are scalar loops, the $H_2$ loops have the right sign, i.e. they restabilize the potential and can lead to a minimum in $v$, so that EWSB is driven by the DM inert Higgs doublet. Imposing that the effective potential has an extremum in $v = 246$ GeV, the Higgs mass at one-loop is given by

$$M_h^2 = \frac{d^2 V_{\text{eff}}}{dh^2} = m_h^2 + \frac{1}{32\pi^2} \left[ 6 \lambda_1 f(m_h^2) + \lambda_L f(m_{H_0}^2) + 2 \lambda_1 f(m_{E_0}^2) + \lambda_S f(m_{A_0}^2) + 4 \lambda_1 f(m_{h^+}^2) + 2 \lambda_1 f(m_{H^+}^2) + 36 \lambda_1^2 h^2 \log \frac{m_h^2}{\mu^2} + \lambda_L^2 h^2 \log \frac{m_{H_0}^2}{\mu^2} + 4 \lambda_1^2 h^2 \log \frac{m_{E_0}^2}{\mu^2} + \lambda_L^2 h^2 \log \frac{m_{A_0}^2}{\mu^2} + 8 \lambda_1^2 h^2 \log \frac{m_{h^+}^2}{\mu^2} + 2 \lambda_L^2 h^2 \log \frac{m_{H^+}^2}{\mu^2} - 36 g_t^2 h^2 f(m_t^2) - 12 g_t^4 h^2 \right]_{(h)=v}$$

with $f(m^2) = m^2(\log(m^2/\mu^2) - 1)$. 

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Since $H_2$ has no vacuum expectation value, there is no mixing between the scalars and it is straightforward to compute the contribution of one-loop corrections to the mass of the other scalars from the second derivative of the effective potential around the Higgs vev. This still requires to keep track of the dependence of the propagators on $h$, $H_0$, $A_0$ and $H^{\pm}$ though. The fact that there is no mixing also means that the extremum is necessarily a minimum if all masses are positive. The result is

$$M_{H_0}^2 = \frac{\partial^2 V_{\text{eff}}}{\partial h^2} = m_{H_0}^2 + \frac{1}{32\pi^2} \left[ \lambda_L f(m_{H_0}^2) + 6\lambda_2 f(m_{H_0}^2) \right]$$

$$+ \left( \lambda_S f(m_{G_0}^2) + 2\lambda_2 f(m_{A_0}^2) + 2\lambda_3 f(m_{H_0}^2) + 4\lambda_2 f(m_{H_0}^2) \right)$$

$$- 2\lambda_2 v^2 g(m_{H_0}^2, m_{H_0}^2 - 2\lambda_2 v^2 g(m_{G_0}^2, m_{A_0}^2) - (\lambda_4 + \lambda_5) v^2 g(m_{H_0}^2, m_{H_0}^2) \right] \bigg|_{(h)=v}$$

$$M_{A_0}^2 = \frac{\partial^2 V_{\text{eff}}}{\partial A_0^2} = m_{A_0}^2 + \frac{1}{32\pi^2} \left[ \lambda_L f(m_{A_0}^2) + 2\lambda_2 f(m_{H_0}^2) \right]$$

$$+ \left( \lambda_S f(m_{G_0}^2) + 6\lambda_2 f(m_{A_0}^2) + 2\lambda_3 f(m_{H_0}^2) + 4\lambda_2 f(m_{H_0}^2) \right)$$

$$- 2\lambda_2 v^2 g(m_{A_0}^2, m_{A_0}^2) - 2\lambda_2 v^2 g(m_{G_0}^2, m_{A_0}^2) - (\lambda_4 - \lambda_5) v^2 g(m_{H_0}^2, m_{H_0}^2) \right] \bigg|_{(h)=v}$$

$$M_{H^{\pm}}^2 = \frac{\partial^2 V_{\text{eff}}}{\partial H^{\pm} \partial \bar{H}^{\pm}} = m_{H^{\pm}}^2 + \frac{1}{32\pi^2} \left[ \lambda_3 f(m_{H_0}^2) + 2\lambda_2 f(m_{A_0}^2) + \lambda_3 f(m_{G_0}^2) \right]$$

$$+ 2\lambda_2 f(m_{A_0}^2) + 2(\lambda_3 + \lambda_1) f(m_{H_0}^2) + 8\lambda_2 f(m_{H_0}^2) - \frac{1}{2} (\lambda_4 + \lambda_5) v^2 g(m_{H_0}^2, m_{H_0}^2)$$

$$- 2\lambda_3 v^2 g(m_{H_0}^2, m_{H_0}^2) - \frac{1}{2} (\lambda_4 - \lambda_5) v^2 g(m_{H_0}^2, m_{H_0}^2) \right] \bigg|_{(h)=v}.$$  

with $g(m_1^2, m_2^2) = [f(m_1^2) - f(m_2^2)]/(m_2^2 - m_1^2)$.

4 Constraints

In order that this dynamical mechanism of EWSB, driven by the DM doublet, does work, there are essentially 3 constraints:

1) EWSB. The general strategy is simple. The contribution of at least some of the loops with $H_2$ particles must be large enough to compensate the large, negative, contribution of the top quark. This requires that at least one of the $\lambda_{3,4,5}$ couplings must be large and positive. This will inevitably drive some of the scalar particle masses in the few hundred GeV range. Imagine that EWSB is driven by loop corrections of $H^{\pm}$ and $A_0$, with $\lambda_3 \simeq \lambda_S$. In this case the $\lambda_3,S$ contribution is relevant with respect to the top loop one provided $\lambda_{3,S} > 2g_t$. Asking that their contribution is large enough for the Higgs mass to be above $\sim 115$ GeV requires $\lambda_{3,S} > 5g_t^2$, approximately, i.e. fairly large but still perturbative quartic couplings. This gives $M_{H^{\pm},A_0} > 130$ GeV.

2) DM mass. Calculating the $H_0$ relic density using the one loop induced coupling $\lambda_{L}^{\text{eff}} = \frac{1}{v} \frac{\partial^2 V_{\text{eff}}}{\partial h \partial H_0} \equiv \frac{1}{v} \frac{\partial M_{H_0}^2}{\partial v}$, the low mass regime turns out to be still perfectly viable. Since at least one of the components of the inert Higgs doublet must be very heavy to break the electroweak symmetry while, in this case, the DM candidate must be lighter than $M_W$, this leads to large mass splittings between at least 2 of the inert Higgs components.

As for the large DM mass regime, it can be shown that it can work only for less phenomenologically interesting special cases. In the following we will consider only the low mass regime.

3) Electroweak precision measurements. The most important constraints on the model from electroweak precision measurements comes from the $\rho$ parameter or equivalently the Peskin-Takeuchi $T$ parameter. A doublet with large mass splitting gives a contribution

$$\Delta T = \frac{1}{32\pi^2 v^2} \left[ f(M_{H^{\pm}}, M_{H_0}) + f(M_{H^{\pm}}, M_{A_0}) - f(M_{A_0}, M_{H_0}) \right]$$  

(5)
Table 1: Instances of parameters with WMAP DM abundance. Also given are the relative contribution of Higgs mediated annihilation \((h_{BR})\) and gauge processes \((W_{BR})\).

<table>
<thead>
<tr>
<th>(\lambda_1)</th>
<th>(\lambda_2)</th>
<th>(\lambda_3)</th>
<th>(\lambda_4)</th>
<th>(\lambda_5)</th>
<th>(M_h)</th>
<th>(M_{H_0})</th>
<th>(M_{A_0})</th>
<th>(M_{H^\pm})</th>
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<th>(W_{BR})</th>
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</tr>
<tr>
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<td>120</td>
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</tr>
<tr>
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<tr>
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<td>64</td>
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<td>150</td>
<td>54</td>
<td>535</td>
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</table>

with \(f(m_1, m_2) = (m_1^2 + m_2^2)/2 - m_1^2 m_2^2/(m_1^2 - m_2^2) \ln(m_1^2/m_2^2)^6\). To give an idea of what is going on, the contribution from \(M_{H^\pm} \sim 450\) GeV and \(M_{DM} \sim 75\) GeV tree level masses gives \(\Delta T \sim 1\), while electroweak precision measurements impose \(|\Delta T| \lesssim 0.2\). There is however a nice and painless cure to this problem: as a quick inspection of Eq. (5) reveals, if either \(H_0\) or \(A_0\) is degenerate with \(H^\pm\), the contribution of the inert doublet to the \(\Delta T\) parameter vanishes identically. Physically, this is due to the existence of a custodial symmetry in the limit \(M_{H^\pm} = M_{A_0}\) or \(M_{H^\pm} = M_{H_0}\) (i.e. \(\lambda_5 = \pm \lambda_3\)). Technically, an exact or approximate custodial symmetry does not only avoid large corrections to the \(T\) parameter. It also implies that it is no fine tuning to take, for instance, the DM particle to be much lighter than the other components of the inert doublet (i.e. \(\lambda_L\) or \(\lambda_S\) much different from the other quartic couplings) as required by the EWSB and DM constraints.

From the three constraints above, we can now consider four cases, see the numerical examples of Table 1. Case I corresponds to a light \(H_0\) and to two heavy, nearly degenerate \(A_0\) and \(H^\pm\) (i.e. \(m_{H_0} << m_{A_0} \simeq m_{H^+}\) or \(\lambda_L << \lambda_S \simeq \lambda_3\)). Case II has a reversed hierarchy, i.e. \(m_{H_0} \lesssim m_{H^+} < m_{A_0} \) or \(\lambda_L \lesssim \lambda_3 < \lambda_S\)). The two last corresponds to \(A_0\) as the DM candidate, with \(m_{A_0} < m_{H_0} \simeq m_{H^+}\) (case III) and \(m_{A_0} \lesssim m_{H^+} < m_{H_0}\) (case IV). Cases III and IV can be obtained from cases I and II simply by switching \(H_0\) with \(A_0\). This leaves the relic density unchanged, so that Table 1 is relevant for these cases too.

All the examples of Table 1 have a DM abundance in agreement with WMAP data. As announced, we observe that some of the quartic couplings must be large. Also, in all the working cases the DM mass is below \(M_W\). In Case I (similarly case III), the DM abundance is determined by its annihilation through the Higgs particle only and thus depends on \(M_h\) and the effective trilinear \(hH_0H_0\) coupling, i.e. \(\lambda^{eff}_{L}\) above. For various, albeit large, couplings we found the correct abundance for DM masses in the range \(M_{H_0} \sim (10 - 72)\) GeV. Below this range, the Higgs mediated annihilation is too suppressed. For this calculation the one-loop contribution to \(\lambda^{eff}_{L}\) is important in some cases. In case II (resp. case IV) coannihilation through the \(W^+\) can play a role if the \(H^+ - H_0\) (resp. \(H^+ - A_0\)) splitting is not too large. Notice that the masses of \(H^\pm\) quoted in Table 1 are consistent with collider data because the \(H^+\) does not couple to fermions, is short lived and, if \(M_{H^\pm} > M_Z/2\), does not contribute to the width of the Z boson.

Imposing the perturbativity condition that the quartic couplings \(\lambda_{3,L,S}\) are smaller than \(e.g. 2\pi\) or \(4\pi\) gives \(M_h \lesssim 80\) GeV or \(M_h \lesssim 175\) GeV in Cases II and IV while for Cases I and III we have \(M_h \lesssim 150\) GeV or \(M_h \lesssim 350\) GeV. We have checked that these \(M_h\) bounds can be saturated, keeping \(\Omega_{DM} \sim 0.22\).

In the Table we considered only the case \(\mu_1 = \mu_2 = 0\) because it is a particularly clear and intriguing case. It shows an example of model where starting from no scale at all, through dimensional transmutation, one can generate all scales of the SM. This cannot be realized in the SM but can work adding to the SM the DM particle, which anyway has to be added to the SM, as in the inert Higgs doublet model. Moreover it shows clearly that it is possible to work in a regime where both DM mass scale and electroweak scale are directly related to a small.
unique scale. The later feature holds more generally as long as $\mu_{1,2}$ are small with respect to the electroweak scale. In view of the hierarchy problem it is however difficult to justify the strict $\mu_1 = \mu_2 = 0$ conformal case. Note also that this scenario of EWSB driven by DM can be realized over a large parameter range beyond the case $\mu_1, \mu_2 = 0$. In particular in the case $\mu_1^2 > 0$.

The existence of a second Higgs doublet has several consequences for colliders. The main ones is that if $m_{DM} < m_h/2$, the Higgs can decay invisibly to a pair of DM particles. This leads to a smaller branching ratio of $h \to bb$, and thus to a slightly lower bound on the Higgs mass from LEP data: $M_h > 105$ GeV instead of 114.4 GeV. Similarly the suppression of the visible branching ratios render more difficult but not impossible the search for the Higgs boson at LHC. Possibility of tests at LHC, by producing the inert Higgs doublet components, do exist.

5 Summary

We have shown that Dark Matter in the form of the lightest neutral component of a single inert scalar doublet could be responsible for EWSB. As a result of all constraints we get the bound on the mass of the Higgs $M_h \lesssim 350$ GeV while the mass of dark matter is in the range $M_{DM} \sim (10-72)$ GeV. Such a DM candidate is in a range of couplings that makes it accessible to both direct (ZEPLIN, Xenon,...) and indirect (GLAST) future searches (cf Figure 5 of §). Another interesting feature of our framework is that it provides a hint for why the DM mass would be around the electroweak scale, as required by the WIMP paradigm, i.e. $M_{DM} \propto v$ in our scenario in a large part of the parameter space.

Acknowledgments

We thank the FNRS-FRS for support.

References

We review the status of models of electroweak symmetry breaking in a slice of anti–de Sitter space. These models can be thought of as dual to strongly interacting theories of the electroweak scale. After an introduction to some generic issues in bulk theories in AdS$_5$, we concentrate on the model-building of the Higgs sector.

1 The Hierarchy Problem and Strong Dynamics

As the Large Hadron Collider (LHC) gets close to start taking data, we continue to ponder what new physics could appear at the TeV scale. In the standard model (SM) the electroweak symmetry is broken by a scalar doublet. This implies the existence of an elementary Higgs boson that must be relatively light ($< 1$ TeV) to unitarize electroweak scattering amplitudes, even lighter to satisfy electroweak precision constraints. However, its mass, and with it the electroweak scale, is unstable under radiative corrections. In order to keep it below the TeV scale and close to $v \simeq 246$ GeV, the bare mass parameter (presumably controlled by ultraviolet physics) must be finely adjusted to cancel against quadratically divergent loop corrections driven by the SM states. The need for this cancellation is highly unnatural and is called the hierarchy problem. In order for naturalness to be restored, new physics must cancel the quadratic divergences at a scale not far above the TeV scale.

The solution of the hierarchy problem is likely to shed light on the origin of electroweak symmetry breaking (EWSB). We can classify scenarios of new physics beyond the SM by how they solve the hierarchy problem. For instance, in supersymmetric theories at the weak scale$^1$, the quadratic divergences in the Higgs mass squared are canceled by the contributions of superpartners of the SM particles. Soft SUSY breaking allows for enough contributions to $m^2_h$ to trigger EWSB radiatively.
An alternative to solve the hierarchy problem is the possibility that some sort of new strong dynamics is present at or just above the TeV scale\(^2\). For instance, in Technicolor theories\(^3,4\) the new strong interaction becomes strong enough at the TeV scale to trigger the condensation of Techni-fermions. If some of these are \(SU(2)\)_L doublets, this triggers EWSB. There is no Higgs in this type of QCD-like scenario. The hierarchy problem is solved by dimensional transmutation, i.e. just as in the strong interactions the coupling becomes strong at low energies triggering EWSB naturally. The TeV scale is encoded in the running of this new strong coupling in the same way that the typical scale of hadronic physics (~1 GeV) is encoded in the running of \(\alpha_s\).

The trouble starts when one tries to construct the operators responsible for fermion masses. For this purpose one needs to extend the gauge group in Extended Technicolor (ETC) models. If at some energy scale \(\Lambda_{\text{ETC}}\) the ETC gauge bosons acquire masses, they generate four-fermion operators involving both fermions and Techni-fermions:

\[
\frac{g^2_{\text{ETC}}}{M^2_{\text{ETC}}} \bar{f}_L f_R \tilde{T}_R T_L \tag{1}
\]

At the lower TC scale \(\Lambda_{\text{TC}}\), the formation of the Techni-condensate \(\langle \tilde{T}_L T_R \rangle \sim \Lambda^3_{\text{TC}}\) results in a fermion mass of the order of

\[
m_f \simeq \frac{g^2_{\text{ETC}}}{M^2_{\text{ETC}}} \Lambda^3_{\text{TC}} \tag{2}
\]

Clearly, if one wants to explain the fermion mass hierarchy one needs either several very different ETC scales, or some non-standard type of dynamics, most probably both. Particularly troublesome for ETC models are heavier masses, say above the charm mass. To obtain the top mass the ETC mechanism fails given than the ETC scale should be right on top of the TC scale.

Several complications of the ETC/TC idea allow for the fermion mass hierarchy (tumbling\(^5\), walking TC\(^6\)) and even for the top quark mass (top-color assisted TC\(^7\)). In any case, it is clear that the dynamics associated with TC/ETC models must be quite different from that of a simple scaled up QCD-type theory. In addition to their problems with the generation of fermion masses, scaled up QCD TC models tend to predict a rather large \(S\) parameter

\[
S \simeq \frac{N}{6\pi} \tag{3}
\]

where \(N = N_{\text{tc}} N_{\text{D}}\) is the product of the number of Techni-colors and the number of weak doublets.

Despite all these problems, the idea that a new strong interaction at or just above the TeV scale is responsible for EWSB (and solves the hierarchy problem) remains attractive. Perhaps the new strong interaction is quite different from QCD and we do not have yet neither the theoretical tools nor the experimental guidance needed to build it.

2 Building Strongly Coupled Theories in AdS\(_5\)

Fortunately, there is a way to build a large array of strongly coupled theories of EWSB by using the AdS/CFT correspondence\(^8\). The conjecture relates a Type IIB string theory on AdS\(_5\) \(\times S^5\) with a four-dimensional \(N = 4\) SU(\(N\)) gauge theory, which is a conformal field theory. By extension, we can assume that at low energies we can describe the string theory by a higher dimensional field theory. As long as the AdS radius \(R_{\text{AdS}}\) is much larger than the string scale \(\ell_s\) the description in the higher-dimensional theory is weakly coupled. On the other hand, this leads to \(\tilde{g} N \gg 1\), where \(\tilde{g}\) is the gauge coupling of the Yang-Mills theory. Then, the description of a weakly coupled theory in AdS\(_5\) corresponds to a strongly coupled four-dimensional theory. We can think of the large \(N\) limit of the 4D gauge theory in terms of planar diagrams, which in turn are reminiscent of a loop expansion in the topology of the world-sheet in string theory.
In the original AdS/CFT correspondence, the boundary of AdS space is set at infinity, and is just a Minkowski 4D boundary. The corresponding 4D theory is exactly conformal, and therefore is not suitable for building models of EWSB. For our purposes, we want to consider an ultra-violet (UV) cutoff of the 4D theory, corresponding to a UV boundary at a finite coordinate in the extra dimension. Also, in order to obtain the description of an interesting strongly coupled theory, we require that the 4D gauge theory leads to non-trivial dynamics which triggers EWSB. This infra-red (IR) physics can be mimicked in the 5D theory by the appearance of an IR boundary. Thus, if we put the UV boundary at the origin of the extra dimension, \( y = 0 \), the statement of the AdS/CFT reads

\[
\int [D\phi_0] e^{i S_{UV}[\phi_0]} \int [D\phi_{CFT}] e^{i S_{CFT}[\phi_{CFT}]} + i \int d^4 x \phi_0 \mathcal{O} = \int [D\phi] e^{i S_{\text{bulk}}[\phi]} \tag{4}
\]

where \( \mathcal{O} \) is an operator in the strongly coupled 4D theory, and \( \phi_0 \equiv \phi(x, y = 0) \) is a UV boundary field which acts as a source for the 4D operator \( \mathcal{O} \). The correspondence in Eq. (4) states the the 4D action is equivalent to the effective action for the source field \( \phi_0 \) on the UV boundary, which is obtained by integrating over the bulk degrees of freedom of \( \phi(x, y) \):

\[
e^{i S_{\text{eff}}[\phi_0]} = \int_{\phi_0} [D\phi] e^{i S_{\text{bulk}}[\phi]} \tag{5}
\]

Then, n-point functions of the 4D theory can be obtained by using \( S_{\text{eff}}[\phi_0] \) as a generating functional. It is in this sense that the bulk degrees of freedom are determining the dynamics of the 4D theory.

We then see that for us to build models of strongly coupled theories, for instance to explain EWSB, we must specify a weakly coupled 5D theory in AdS\(_5\). The type of strongly coupled theories that we can build this way is not completely general. For instance, it requires that it be “large N”, which should in principle translate in the presence of narrow resonances. In what follows we consider the steps to build such theories.

2.1 Solving the Hierarchy Problem in a slice of AdS\(_5\)

The starting point to build theories of EWSB using holography is the Randall-Sundrum setup as a solution of the hierarchy problem\(^9\). We consider an extra dimension compactified on an orbifold, i.e. \( S_1/Z_2 \), with the metric

\[
ds^2 = e^{-2k y} g^{\mu \nu} dx_\mu dx_\nu - dy^2 ,
\]

where \( k \sim M_P \) is the AdS curvature. The orbifold compactification \( S_1/Z_2 \) results in a slice of AdS in the interval \([0, \pi R]\), with \( R \) the compactification radius. This metric is a solution of Einstein’s equations if we fine-tune the bulk cosmological constant to cancel the brane tensions. This choice of metric means that the graviton’s wave-function is exponentially suppressed away from the origin. In general, this metric exponentially suppresses all energy scales away from the origin. Then, if the Higgs field is localized a distance \( L = \pi R \) from the origin,

\[
S_H = \int d^4 x \int_0^{\pi R} dy \sqrt{g} \delta(y - \pi R) \left[ g^{\mu \nu} \partial_\mu H^\dagger \partial_\nu H - \lambda \left( |H|^2 - v_0^2 \right)^2 \right] \tag{7}
\]

where \( \lambda \) is the Higgs self-coupling and \( v_0 \) is its vacuum expectation value (VEV). The latter must satisfy \( v_0 \sim k \) for the theory to be technically natural. Taking into account the exponential
factors in $g^\mu\nu$ and $\sqrt{g}$, the renormalization of the Higgs field required to render its kinetic term canonical results in an effective four-dimensional VEV given by

$$v = e^{-k\pi R} v_0$$

Thus, in order for the weak scale $v \approx 246$ GeV to arise at the fixed point in $y = \pi R$, we need $kR \sim (10-12)$. Then, if the Higgs is for some reason localized at or nearly at this location, this setup constitutes a solution to the hierarchy problem.

In the original Randall-Sundrum (RS) proposal only gravity propagates in the extra dimension. However, this presents several problems, mostly associated with the fact that is not possible to sufficiently suppress higher-dimensional operators. For instance, operators mediating flavor violation might only be suppressed by the TeV scale, $k e^{-k\pi R}$. Grand Unified Theories (GUTs) might not be viable since we cannot effectively suppress proton decay. Many of these problems are solved when fermions and gauge bosons are allowed in the AdS$_5$ bulk. In fact, in order to solve the hierarchy problem, the only field that must remain localized near the $\pi R$ or TeV brane is the Higgs. In addition, allowing the standard model fields in the bulk opens up a large number of model building possibilities including viable GUTs$^{10}$ and the modeling of the origin of flavor$^{11,12}$ just to mention two very prominent cases. But most importantly, it gives us a tool to build models of EWSB that address its dynamical origin. In building bulk models of EWSB we must dynamically explain why is the Higgs localized near the TeV brane. Building such models is like constructing strongly coupled theories of EWSB but from a different perspective.

### 2.2 Bulk RS Theories and the Origin of Flavor

Writing a theory in the bulk require several ingredients. First, we must decide what the gauge symmetry should be. It turns out that the SM electroweak symmetry, $SU(2)_L \times U(1)_Y$, is not suitable for this since it results in a large value of the parameter $T$. The reason is that there is large explicit isospin violation in the bulk: in addition to the SM-sanctioned isospin violation proportional to $g'/g = \tan \theta_W$, the bulk adds the isospin violation of the Kaluza-Klein (KK) modes of the SM gauge fields. As a consequence, it is necessary to implement a gauge symmetry that would exhibit isospin symmetry in the bulk. A minimal extension of the SM with this feature is $SU(2)_L \times SU(2)_R \times U(1)_X$. This bulk symmetry may be broken by boundary conditions to the SM gauge group or directly to $U(1)_{\text{EM}}$. Additional discrete symmetries may be imposed to protect the $Z \to b\bar{b}$ coupling from large deviations.$^{15}$

One important consequence of writing a bulk theory, is that the KK states would start from masses of the order of $M_{KK} \sim O(1)$ TeV, independently of whether they correspond to fermions, gauge bosons or even scalars. This also means that the wave-function of the KK excitations in the extra dimension also peaks at $y = \pi R$, i.e. at the TeV brane. This is related to the fact that position in the bulk can be thought of as an energy scale, and is a generic feature independent of the model considered.

Another important issue is the localization of zero modes in the bulk, i.e. their effective 5D wave-function. The strength of couplings between zero-modes and KK modes and, since the Higgs is TeV-localized, of the zero-mode fermion Yukawas, are determined by this. Fermions propagating in the bulk can have a mass term (as long as we assume is an odd mass term). The natural order of magnitude of this fermion bulk mass is $k$, the only dimensionfull bulk parameter. We can then write the bulk fermion mass as

$$M_f = c_f k$$

where $c_f \approx O(1)$. Expanding the 5D fermion $\Psi(x,y)$ in KK modes and solving the equation of motion for the zero mode we see that its bulk wave-function, for instance for a left-handed zero
mode, behaves like \( F_{ZM}^L (y) \sim e^{(\frac{1}{2} - c_L) ky} \) \( (10) \)

Then, if \( c_L > 1/2 \) the zero-mode fermion is localized towards the Planck brane, whereas if \( c_L < 1/2 \) it would be localized near the TeV brane. For a right-handed zero mode, localization near the Planck bran occurs if \( c_R < -1/2 \), and near the TeV brane if \( c_R > -1/2 \). If a zero mode fermion has a large wave-function near the TeV brane it has also a large overlap with the Higgs, which has to be localized there. Thus, heavier fermions must have a wave-function towards the TeV brane, whereas light fermions must be essentially Plank-brane localized in order to explain their small Yukawa couplings. We see then that fermion localization in the AdS\(_5\) bulk provides a potential explanation for the hierarchy of fermion masses in the SM. This emerging picture of flavor requires that the top quark be highly localized towards the TeV brane. Furthermore, the left-handed third-generation doublet contains the left-handed b quark, which then also has to have a substantial TeV localization. The rest of the zero-mode fermions must be localized mostly near the Plank brane. Since the zero-mode gauge bosons are flat in the extra dimension, this does not affect the universality of the zero-mode gauge couplings. However, the KK modes of gauge bosons are TeV-brane localized, and then couple stronger to the TeV-brane localized fermions, i.e. the third generation. The couplings of light fermions to KK gauge bosons are nearly universal, making low energy flavor phenomenology viable, even in the presence of tree-level flavor violation. Also since light fermions are localized near the Plank brane, the scale suppressing higher dimensional operators responsible for proton decay is again \( M_P \).

Then any sign of the characteristic tree-level flavor violation would have to appear in the interactions of the third generation quarks with the KK gauge bosons. Although these would have potentially important effects in flavor physics\(^{16, 17}\), a direct observation of the KK gluon decay into a single top and a jet, coming most likely from \( G^{(1)} \rightarrow t\bar{b} \), would be an unambiguous signal of this theory of flavor \(^{18}\).

Finally, a word about electroweak precision constraints. Since the bulk gauge theory has an isospin symmetry built in, we need not worry about the theory generating a large \( T \) parameter. However, this class of models all share the same problem with the \( S \) parameter. They have an \( S \) parameter which is approximately\(^{13, 19, 20}\)

\[
S_{\text{tree}} \simeq 12\pi \frac{y^2}{M_{KK}^2},
\]

which results in a bound of about \( M_{KK} > 2.5 \) TeV. This is a feature we have to live with in most bulk RS constructions. We can interpret this in the dual 4D picture, as the fact that in the 4D theory we have a large \( N \), which enter in \( S_{\text{tree}} \sim \frac{N}{4} \), where \( N \) is typically the size of the 4D gauge group. This can in principle be made considerably smaller if the light fermions
are allowed out of the Plank brane region and have almost flat profiles in the extra dimensions. Essentially this decouples them from the KK modes and avoids $S_{\text{tree}}^{13}$. But this picture would lack a solution to the fermion mass hierarchy.

3 Electroweak Symmetry Breaking from AdS$_5$

We have arrived at a general picture of these kind of models in AdS$_5$:

- The Higgs is TeV-brane localized in order to solve the gauge hierarchy problem.
- Fermion localization explains the fermion mass hierarchy: light fermions are Plank-brane localized resulting in small Yukawa couplings. Heavier fermions are TeV-brane localized ($t_R$, $t_L$ and $b_L$).
- The bulk gauge symmetry must be enlarged to protect isospin, to be at least $SU(2)_L \times SU(2)_R \times U(1)_X$.

This is already quite a rich structure with a very interesting phenomenology and lots of model building possibilities beyond EWSB and fermion masses. However, if these models are dual to a strongly coupled theory in four dimensions, as the presence of the resonances (i.e. the KK modes) appears to suggest, then we must be a bit disappointed with the Higgs sector. What keeps the Higgs localized to the TeV brane? Do we have to have a Higgs at all? In what follows we briefly discuss three possibilities for the Higgs sector in these theories.

3.1 Higgsless EWSB in AdS$_5$

In the absence of a Higgs, the bulk gauge symmetry must be broken directly to $U(1)_{\text{EM}}$ by boundary conditions (BC) for the bulk gauge fields at the fixed point of the extra dimension. On the Plank brane the BC are such that the bulk gauge symmetry breaks as$^{14}$

$$SU(2)_R \times U(1)_X \longrightarrow U(1)_Y, \quad \text{at} \quad y = 0$$

On the other hand, on the TeV brane

$$SU(2)_L \times SU(2)_R \longrightarrow SU(2)_V, \quad \text{at} \quad y = \pi R$$

in such a way that it it preserves custodial symmetry. Then the gauge symmetry in the bulk has a gauged version of the SM custodial symmetry. This remains as a remnant global symmetry, resulting in the correct masses for the $W$ and the $Z$.

More problematic in these models is how to give fermions their masses. Zero-mode fermions can obtain isospin conserving masses through a TeV-brane localized mass term. Since $SU(2)_V$...
is not broken, isospin splitting is not generated by these terms. In order to achieve the freedom to have the correct mass spectrum one must introduce the following bulk spectrum (schematically and only for the third generation case):

\[
\begin{align*}
\Psi_L &= (t_L, b_L) \ (2, 1)_{1/6} \\
\Psi_R &= (t_R, b'_R) \ (1, 2)_{1/6} \\
\Psi'_R &= (t'_R, b_R) \ (1, 2)_{1/6}
\end{align*}
\]

Each of these bulk fermions contains both left and right handed components. We can choose the BCs so that the zero-modes of \(\Psi_L\) correspond to the third generation left-handed quark doublet, the zero mode of \(\Psi_R\) is the right-handed top \(t_R\), and the one corresponding to \(\Psi'_R\) is \(b_R\). Mass terms localized in the TeV brane gives rise to fermion masses. It is still true that the larger the zero-mode wave-function is at the TeV brane, the larger its mass would be. But the top mass cannot be so easily adjusted. The reason is that there is a tension between the TeV localization of the top, which if too extreme produces noticeable deviation in the \(Zb_Lb_L\) coupling, and the size of the isospin conserving TeV localized mass. The latter cannot be too large or it would induce a large mixing between \(b'_R\) and the b-quark’s zero mode through the mass term responsible for the top quark mass. This would again result in a deviation of \(Zb_Lb_L\). A way to circumvent this problem is to extend the custodial symmetry by a discrete symmetry, \(P_{LR}\) that relates the two \(SU(2)\)’s. In order to protect the \(b_L\) coupling, it must be included in a bi-doublet of \(SU(2)_L \times SU(2)_R\). Then the right-handed fermions could be in singlets or triplets of the \(SU(2)\)’s. For instance, if \(t_R \sim (1, 1)_{2/3}\), then it does not contain any fermion that could mix with the left-handed b. On the other hand, the field resulting in the right-handed zero mode would have to be a full \(O(4)\) triplet \(\Psi_R \sim (3, 1)_{2/3} \oplus (1, 3)_{2/3}\) resulting in a distinct spectrum of KK fermions.

But the most important signal of this scenario stems from the absence of the Higgs as a unitarizing field in \(VV\) scattering, with \(V = W^\pm, Z\). The unitarization of these amplitudes, unlike in other strongly coupled theories such as (4D) Technicolor, is the result of the presence of the narrow resonances that are the KK modes of the gauge bosons. The constraint of unitarization imposes sum rules on the couplings. For instance,

\[
g_{\gamma\gamma\gamma\gamma} = g_{\gamma\gamma\gamma Z}^2 + g_{\gamma\gamma\gamma\gamma}^2 + \sum_n (g_{\gamma\gamma\gamma V(n)}^2) \leq 4M_W^2
\]

The requirement of unitarity of gauge boson scattering amplitudes means that the KK modes of gauge bosons cannot be too heavy. For example, if one wants to preserve perturbative unitarity, they must be below the TeV scale. Higgsless models in AdS\(_5\) are then characterized by relatively low mass KK excitations.

### 3.2 Gauge-Higgs Unification

A remarkable mechanism to obtain a Higgs field naturally localized near the TeV, naturally light and suitable for EWSB is that in which the Higgs comes from a gauge field in 5D. In general, a 5D gauge field \(A_M(x, y), M = 0, 1, 2, 3, 5\), can be decomposed in a vector \(A_\mu(x, y)\) and a scalar component \(A_5(x, y)\). If we want to extract the Higgs \(SU(2)_L\) doublet from a gauge field in 5D, then the gauge symmetry in 5D has to be enlarged beyond the SM gauge symmetry. In order to illustrate how this works let us take a simple example, an \(SU(3)\) bulk gauge theory. We can use BCs to break this gauge symmetry as \(SU(3) \rightarrow SU(2)_L \times U(1)_Y\). By choosing the BCs...
appropriately, the gauge fields are
\[
 t^a A_\mu^a : \begin{pmatrix}
 (+, +) & (+, +) & (-, -) \\
 (+, +) & (+, +) & (-, -) \\
 (-, -) & (-, -) & (++, +)
\end{pmatrix}
\]
\[
 t^a A_5^a : \begin{pmatrix}
 (-, -) & (-, -) & (++, +) \\
 (-, -) & (-, -) & (++, +) \\
 (++, +) & (++, +) & (-, -)
\end{pmatrix}
\]
where \(a = 1, 2, 3\) is the \(SU(3)\) adjoint index and \(t^a\) are the \(SU(3)\) generators, and we use the fact that the BCs for the \(A_5^a(x, y)\) are always opposite from those for \(A_\mu^a(x, y)\). The signs correspond to the BCs in the Plank and TeV branes respectively. Thus, we see that the spectrum of zero-mode gauge bosons corresponds to \(SU(2) \times U(1)\) as the gauge symmetry. The symmetry has been reduced or broken by this choice of BCs. Furthermore, we see that the \(A_5^a(x, y)\)'s corresponding to the “broken” generators, i.e. the generators for which \(A_\mu^a(x, y)\) does not have a zero mode, have zero modes (i.e. \((+, +)\) BCs). In fact, these constitute four real degrees of freedom that can be seen to be a doublet of \(SU(2)\) and its adjoint. Then, we can identify this \(SU(2)\) doublet with the Higgs. In the case of an \(AdS_5\) metric, if we impose the unitary gauge, this results in
\[
 \partial_y (k e^{-ky} A_5(x, y)) = 0,
\]
which results in a scalar doublet with a profile localized towards the TeV brane. Thus, we achieved Higgs TeV-brane localization and extracted the Higgs from a gauge field in the bulk. These models are clearly related to Little Higgs theories, where the Higgs is a (pseudo-)Nambu-Goldstone boson (pNGB): the Higgs is associated to the broken generators of a symmetry, which from the 4D interpretation would be a global one. It cannot have a potential since shift symmetry (the remnant gauge symmetry for the \(A_5^a\) zero-modes) forbids it. Thus, these models should be dual to 4D theories of a composite Higgs, where the Higgs is a pNGB.

This simple \(SU(3)\) model of Gauge-Higgs unification has a lot of the features that we want. However, it does not have custodial symmetry in the bulk. The way to cure this is to simply enlarge the bulk gauge symmetry. The minimal realistic model of Gauge-Higgs unification in \(AdS_5\) requires that we start with \(SO(5) \times U(1)_X\) broken down to \(SO(4) \times U(1)_X\) on the TeV brane, whereas it is reduced to \(SU(2)_L \times U(1)_Y\) on the Plank brane. Just as in the previous case, in the unitary gauge only the \(A_5^a(x, y)\)'s associated with the broken generators, i.e. transforming in the coset space \(SO(5)/SO(4)\), have zero modes. These are arranged in a \(4\) of \(SO(4)\), or a bi-doublet of \(SU(2)_L \times SU(2)_R\).

Fermion masses are not very problematic and can be obtained by localization just as in the generic models with a TeV-brane localized Higgs. Regarding electroweak precision constraints, these models can evade them more efficiently. In particular, the \(S\) parameter can be made in agreement with experiment for gauge KK masses above 2 TeV. On the other hand, KK fermions could be quite a bit lighter, especially once the additional discrete custodial symmetry is introduced in order to keep the \(Z \to b_L b_L\) in check. For instance, possible embeddings are
\[
 5_{2/3} = (2, 2) \oplus (1, 1)
\]
or in a
\[
 10_{2/3} = (2, 2) \oplus (1, 3) \oplus (3, 1)
\]
With it, KK fermions can be as light as 500 GeV, and in some cases would have exotic charge assignments. The main reason for some KK fermions to be this light has to do with the need to get the top mass correctly. There are striking signals at the LHC, for instance from the pair production of charge \(5/3\) KK fermions, a distinct signature for the presence of the extended custodial symmetry.
3.3 Higgs from Fermion Condensation

Another alternative to dynamically generate the Higgs sector of the bulk Randall-Sundrum scenario is the condensation of zero-mode fermions. Since the localization of fermions near the TeV brane implies they must have strong couplings to the KK excitations of gauge bosons, it is possible that the induced four-fermion interaction is strong enough to result in fermion condensation, thus triggering EWSB. Among the SM fermions, the candidate would be the top quark: it has the strongest localization toward the TeV brane, therefore the largest coupling to KK gauge bosons, particularly the first excitation of the KK gluon. The four-fermion interaction would result in a top-condensation scenario. But we already know that this does not work if the scale of the underlying interaction is $O(1)$ TeV, as we expect the KK gluon mass to be. The problem is that the mass of the condensing fermion should be about 600 GeV in order for the condensation to result in the correct weak scale, $v \approx 246$ GeV. We can then consider the possibility of a fourth generation, one that is highly localized near the TeV brane, more than the top quark. This results in a effective four-fermion interaction -mostly mediated by KK gluons- that is attractive enough to trigger the condensation of at least one of the fourth-generation quarks. For instance, the up-type fourth-generation quark $U$ has a four-fermion interaction given by

$$\int dy \sqrt{g} \frac{g^L_{01} g^R_{01}}{M_{KK}^2} \left( \bar{U}_L \gamma_\mu t^A U_L \right) \left( \bar{U}_R \gamma^\nu t^A U_R \right),$$

where $g^L_{01}, g^R_{01}$ are the couplings of the left and right-handed $U$ quarks to the first KK mode of the gluon of mass $M_{KK}$, and $t^A$ are the QCD generators. After Fierz rearrangement, we can re-write this interaction as

$$\frac{g^2_U}{M_{KK}^2} \left\{ U^a_L U^b_R U^b_R U^a_R - \frac{1}{N_c} U^a_L U^b_R U^b_L U^a_R \right\},$$

where $a, b$ are $SU(3)_c$ indices, and we have defined

$$g^2_U \equiv g^L_{01} g^R_{01}. \quad (20)$$

The color singlet term in (19) is attractive, whereas the color octet is repulsive, as well as suppressed by $1/N_c$. There is a critical value of $g^2_U$ above which there forms a condensate $(\bar{U}_L U_R)$ leading to electroweak symmetry breaking and dynamical masses for the condensing fermions. This is

$$g^2_U > \frac{8\pi^2}{N_c}. \quad (21)$$

One can also write an effective theory in terms of a scalar doublet which becomes dynamical at low energies. So this theory gets a composite Higgs that is heavy and made of the already mostly-composite fourth-generation up quarks. The $U$ quark gets a large dynamical mass. All other zero-mode fermions, including the SM fermions and the other fourth-generation zero-modes, get masses through four-fermion interactions with the $U$ quark. These operators come from bulk higher dimensional operators such as

$$\int dy \sqrt{g} \frac{C^{ijkl}}{M_{KK}^3} \Psi^i_L(x,y) \Psi^j_R(x,y) \Psi^k_R(x,y) \Psi^l_L(x,y),$$

where $C^{ijkl}$ are generic coefficients, with $i, j, k, \ell$ standing for generation indices as well as other indices such as isospin, and the $\Psi(x,y)$’s can be bulk quarks or leptons. Upon condensation of the $U$ quarks these result in fermion masses that have the desired pattern as long as we choose the localization parameters appropriately.

These models have roughly the same $S$ parameter problem as the generic ones. But to the tree-level $S$, now we must add also the loop contributions coming from a heavy Higgs as well as from the fourth-generation. These are, however, not as large as $S_{\text{tree}}$. 

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The masses of the fourth-generation zero modes could be as small as 300 GeV due to mixing with KK modes, and as large as $\sim 600$ GeV. The Higgs is typically rather heavy, in the $(700 - 900)$ GeV range.

The phenomenology of these models is quite different from the three-generation RS models due to the fact that the fourth generation is the one that couples strongest to the KK gauge bosons, i.e. to the new physics in the $s$ channel. For instance the branching ratio of a KK gauge boson to $U^{(1)}\bar{U}^{(1)}$ is likely to be $5 - 10$ times larger than that for the $tt$ channel, which is the one always dominating in three-generation RS models. Also the KK gluon tends to be very broad and can only be seen as an excess in the production of the fourth-generation, which is dominated by QCD. In general, all KK gauge bosons will be considerably broader, although electroweak KK gauge bosons might have more manageable widths.

4 Conclusions

Model building in a slice of AdS$_5$ extends the Randall-Sundrum solution to the hierarchy problem, and opens up the possibility of addressing dynamically the origin of EWSB and fermion masses, among other things. Through the AdS/CFT correspondence we can see that these weakly coupled 5D theories are dual to strongly coupled 4D gauge theories. Thus, building bulk RS theories of EWSB corresponds, through holography, to certain strongly coupled electroweak sectors. Although the type of strongly coupled theory is not completely generic, this procedure gives us access to a large variety of strongly coupled theories of the electroweak sector.

In addition to solving the hierarchy problem, RS bulk models provide a framework to understand the hierarchy of fermion masses. This implies the presence of tree-level flavor violation with KK gauge bosons. Rigorous compatibility with low energy data may require some level of flavor symmetry in the bulk. However, this is not surprising since fermion localization already signaled some amount of flavor breaking in the UV. The central point is that the metric in AdS$_5$ provides the necessary large scale separation between the light and the heavy fermions.

Finally, there are several alternatives for the Higgs sector. Higgsless models are viable, although require some measure of fine-tuning to cancel contributions to $Z \rightarrow b_Lb_L$. Gauge-Higgs unification models are very promising and in best agreement with electroweak precision constraints. They provide a dynamical origin for the localization of the Higgs near the TeV brane, and in the Holographic dual correspond to a composite pNGB Higgs. Finally, another alternative to localize a composite Higgs near the TeV brane, is the condensation of fourth-generation quarks via the attractive four-fermion interaction mediated mostly by the KK gluon. This is a distinct possibility, a realization of the fourth-generation condensation proposed long ago by Bardeen, Hill and Lindner. These three Higgs sector possibilities have very different phenomenology and should be distinguishable at the LHC.

References

ELECTROWEAK CONSTRAINTS ON SEE-SA W MESSENGERS AND THEIR IMPLICATIONS FOR LHC

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We review the present electroweak precision data constraints on the mediators of the three types of see-saw mechanisms. Except in the see-saw mechanism of type I, with the heavy neutrino singlets being mainly produced through their mixing with the Standard Model leptons, LHC will be able to discover or put limits on new scalar (see-saw of type II) and lepton (see-saw of type III) triplets near the TeV. If discovered, it may be possible in the simplest models to measure the light neutrino mass and mixing properties that neutrino oscillation experiments are insensitive to.

1 Introduction

As it is well known, the original see-saw mechanism 1, nowadays called of type I, explains the smallness of the light neutrino masses $|m_\nu| \sim 1$ eV invoking a very heavy Majorana neutrino $M_N \sim 10^{14}$ GeV:

$$|m_\nu| \simeq \frac{v^2|\lambda|^2}{M_N} \simeq |V^*|^2 M_N,$$

where $|\lambda| \sim 1$ is the corresponding Yukawa coupling and $v \simeq 246$ GeV the electroweak vacuum expectation value. For reviews see 2,3. Alternatively, if the heavy scale is at the LHC reach $M_N \sim 1$ TeV, it requires a very small heavy–light mixing angle $|V| \sim 10^{-6}$. In its simplest form the model cannot be tested at large colliders, because the heavy neutrino $N$ is a Standard Model (SM) singlet and only couples to SM gauge bosons through its mixing $V$. Hence it is produced through the vertex $-g/\sqrt{2} \bar{\ell} \gamma^\mu V_{\ell N} P_L N W^-_{\mu}$, with $\ell$ a charged lepton, with a cross section proportional to $|V_{\ell N}|^2$, which is strongly suppressed. See Fig. 1-(I). There are two other types of see-saw mechanism giving tree level Majorana masses to the light neutrinos $\nu$, as shown
Figure 1: Examples of production diagrams for same-sign dilepton signals, $l^+l'^+X$, mediated by the three types of see-saw messengers.

Figure 2: See-saw mechanisms of type I, II and III. $\lambda_N$, $\lambda_\Delta$ and $\lambda_\Sigma$ are the Yukawa coupling matrices in the Lagrangian terms $-\bar{\nu}_L\phi^\dagger N_R$, $\bar{\nu}_L\lambda_\Delta(\bar{\nu}\cdot \Delta)l_L$, and $-\Sigma_R\lambda_\Sigma(\bar{\phi}^\dagger \frac{\nu}{2}\nu^j)l_L$, respectively, with $\bar{\nu}_L = -i\tau_2\hat{C}\nu_2$ and $\hat{C}$ the spinor charge conjugation matrix. Whereas $\mu_\Delta$ is the coefficient of the scalar potential term $\bar{\phi}^\dagger(\bar{\Delta}\cdot \Delta)\phi$.

In Fig. 2. In all cases the extra particles contribute at low energies to the dimension 5 lepton number (LN) violating operator

$$\mathcal{O}_5 = \mu_\Delta \bar{\nu}_L\phi^\dagger N_R, \bar{\nu}_L\lambda_\Delta(\bar{\nu}\cdot \Delta)l_L, \bar{\Sigma}_R\lambda_\Sigma(\bar{\phi}^\dagger \frac{\nu}{2}\nu^j)l_L$$

which gives Majorana masses to light neutrinos after spontaneous symmetry breaking. The see-saw of type II \(^5\) in Fig. 2 is mediated by an $SU(2)_L$ scalar triplet $\Delta$ of hypercharge $Y = 1$, implying three new complex scalars of charges $Q = T_3 + Y$: $\Delta^+, \Delta^0$. The see-saw of type III \(^6\) exchanges an $SU(2)_L$ fermion triplet $\Sigma$ of hypercharge $Y = 0$, assumed to be Majorana and containing charged leptons $\Sigma^\pm$ and a Majorana neutrino $\Sigma^0$. The main difference for LHC detection is that the see-saw messengers for these last two mechanisms can be produced by unsuppressed processes of electroweak size (Fig. 1). Their decay, even if suppressed by small couplings, can take place within the detector due to the large mass of the new particle. All three types of see-saw messengers produce LN conserving as well as LN violating signals, but the former have much larger backgrounds. On the other hand, same-sign dilepton signals, $l^\pm l'^\pm X$, do not have to be necessarily LN violating. Thus, in the example in Fig. 1–(II), the decay
coupling \( \lambda_\Delta \) needs not be very small because it is only one of the factors entering in the LN violating expression for \( \nu \) masses (see Table 1). In fact, this process is LN conserving as we can

<table>
<thead>
<tr>
<th>Coefficient</th>
<th>Type I</th>
<th>Type II</th>
<th>Type III</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \alpha_4 )</td>
<td>runs on</td>
<td>( \frac{2\mu_\Delta^2}{M_\Delta^2} )</td>
<td>runs on</td>
</tr>
<tr>
<td>( \frac{(\alpha_5)_{ij}}{\Lambda} )</td>
<td>( \frac{1}{2} \left( \frac{\lambda_\Delta^T \alpha_\Delta^T \alpha_\Delta^T \alpha_\Delta^T}{M_{\alpha \alpha}} \right) )</td>
<td>( -\frac{2\mu_\Delta (\lambda_\Delta)<em>{ij}}{M</em>\Delta^2} )</td>
<td>( \frac{1}{8} \left( \frac{\lambda_\Delta^T \alpha_\Delta^T \alpha_\Delta^T \alpha_\Delta^T}{M_{\alpha \alpha}} \right) )</td>
</tr>
<tr>
<td>( \frac{(\alpha_6)_{ij}}{\Lambda^2} )</td>
<td>( \frac{1}{4} \left( \frac{\lambda_\Delta^T \alpha_\Delta^T \alpha_\Delta^T \alpha_\Delta^T}{M_{\alpha \alpha}} \right) )</td>
<td>-</td>
<td>( \frac{3}{16} \left( \frac{\lambda_\Delta^T \alpha_\Delta^T \alpha_\Delta^T \alpha_\Delta^T}{M_{\alpha \alpha}} \right) )</td>
</tr>
<tr>
<td>( \frac{(\alpha_7)_{ij}}{\Lambda^2} )</td>
<td>( \frac{1}{2} \left( \frac{\lambda_\Delta^T \alpha_\Delta^T \alpha_\Delta^T \alpha_\Delta^T}{M_{\alpha \alpha}} \right) )</td>
<td>-</td>
<td>( \frac{3}{4} \left( \frac{\lambda_\Delta^T \alpha_\Delta^T \alpha_\Delta^T \alpha_\Delta^T}{M_{\alpha \alpha}} \right) )</td>
</tr>
<tr>
<td>( \frac{(\alpha_8)_{ij}}{\Lambda^2} )</td>
<td>( \frac{1}{4} \left( \frac{\lambda_\Delta^T \alpha_\Delta^T \alpha_\Delta^T \alpha_\Delta^T}{M_{\alpha \alpha}} \right) )</td>
<td>-</td>
<td>( \frac{4}{3} \left( \frac{\lambda_\Delta^T \alpha_\Delta^T \alpha_\Delta^T \alpha_\Delta^T}{M_{\alpha \alpha}} \right) )</td>
</tr>
<tr>
<td>( \frac{(\alpha_{\phi \phi})_{ij}}{\Lambda^2} )</td>
<td>( \frac{1}{2} \left( \frac{\lambda_\Delta^T \alpha_\Delta^T \alpha_\Delta^T \alpha_\Delta^T}{M_{\alpha \alpha}} \right) )</td>
<td>-</td>
<td>( \frac{4}{3} \left( \frac{\lambda_\Delta^T \alpha_\Delta^T \alpha_\Delta^T \alpha_\Delta^T}{M_{\alpha \alpha}} \right) )</td>
</tr>
</tbody>
</table>

conventionally assign LN equal to 2 to \( \Delta^{-} \). There are other processes that do violate LN, e.g. when one of the doubly-charged \( \Delta \) in Fig. 2–(II) decays into \( WW \). Then, what does violate LN is the corresponding \( \Delta WW \) vertex, which is proportional to the coupling of the only LN violating term in the fundamental Lagrangian \( \phi_l^*(\tilde{\sigma} \cdot \Delta)^\dagger \phi \), with total LN equal to 2. In the examples in Fig. 1–(I, III) LN is violated in the decay (mass) of the heavy neutral fermion.

In conclusion, all the three mechanisms produce same-sign dilepton signals, but only the last two are observable at LHC 7,8,9,10,11,12,13 in minimal models. Heavy neutrino singlets in particular non-minimal scenarios could also be observed, as described in Section 3.

In the following we first review the experimental constraints on the parameters entering the three see-saw mechanisms, and then the LHC reach for the corresponding see-saw messengers. Complementary reviews on this subject have been presented by other speakers at this Conference (see F. Bonnet, T. Hambye and J. Kersten in these Proceedings).

2 Electroweak precision data limits on see-saw messengers

The low energy effects of the see-saw messengers can be described by the effective Lagrangian

\[
\mathcal{L}_{\text{eff}} = \mathcal{L}_4 + \frac{1}{\Lambda} \mathcal{L}_5 + \frac{1}{\Lambda^2} \mathcal{L}_6 + \ldots,
\]

where \( \Lambda \) is the cut-off scale, in our case of the order of the see-saw messenger masses \( M \), and the different terms contain gauge-invariant operators of the corresponding dimension. The non-zero terms up to dimension 6 are 14,15

\[
\mathcal{L}_4 = \mathcal{L}_{\text{SM}} + \alpha_4 \left( \phi_l^* \phi \right)^2,
\]
\[ \mathcal{L}_5 = (\alpha_5)_{ij} (\bar{l}_L^i \gamma^\mu \phi^0 \gamma^\mu l_L^j) + \text{h.c.}, \]

\[ \mathcal{L}_6 = \left[ (\alpha^{(1)}_{\phi \ell})_{ij} \left( \phi^\dagger l_D^i \alpha \phi \left( \bar{l}_L^j \gamma^\mu \mu^\mu l_L^j \right) \right) + (\alpha^{(3)}_{\phi \ell})_{ij} \left( \phi^\dagger i \sigma \alpha \phi \left( \bar{l}_L^j \gamma^\mu \mu^\mu l_L^j \right) \right) \\
+ (\alpha_{\phi \ell})_{ij} \left( \phi^\dagger \phi \left( \bar{l}_L^j \phi e_R^j \right) + (\alpha^{(1)}_{\phi \ell})_{ijkl} \frac{1}{2} \left( \bar{l}_L^j \gamma^\mu \mu^\mu l_L^j \right) \right) + \text{h.c.} \right] \\
+ \alpha^{(1)}_{\phi} \left( \phi^\dagger \phi \left( \left(D_\mu \phi^\dagger \right) D^\mu \phi \right) + \alpha^{(3)}_{\phi} \left( \phi^\dagger D_\mu \phi \right) \left( \left(D^\mu \phi^\dagger \right) \phi \right) + \alpha_{\phi} \frac{1}{3} \left( \phi^\dagger \phi \right)^3, \]

where we choose the basis of Bliuchmüller and Wyler to express the result. \( l_L \) stands for any lepton doublet, \( e_R \) for any lepton singlet, and \( \phi \) is the SM Higgs doublet. In Table 1 we collect the explicit expressions of the coefficients in terms of the original parameters for each type of see-saw (see Fig. 2 and the table caption for definitions).

Only the dimension 6 operators can give deviations from the SM predictions for the electroweak precision data (EWPD). The operators of dimension 4 only redefine SM parameters. The one of dimension 5 gives tiny masses to the light neutrinos, and contributes to neutrinoless double \( \beta \) decay. An important difference is that the coefficient \( \alpha_5 \) involves LN-violating products of two \( \lambda \)'s or of \( \mu \) and \( \lambda \), while the other coefficients depend on \( \lambda^* \lambda \) or \( |\mu|^2 \). Therefore, it is possible to have large cancellations in \( \alpha_5 \) together with sizeable coefficients of dimension six. Type I and III fermions generate the operators \( O^{(1,3)}_{\phi \ell} \), which correct the gauge fermion couplings. Type II scalars, on the other hand, generate 4-lepton operators and the operator \( O^{(3)}_{\phi} \), which breaks custodial symmetry and modifies the SM relation between the gauge boson and EWPD are sensitive to all these effects and put limits on the see-saw parameters.

There are two classes of processes, depending on whether they involve neutral currents violating lepton flavour (LF) or not. The first class puts more stringent limits, but only on the combinations of coefficients entering off-diagonal elements. The second class is measured mainly at LEP and constrains the combinations in the diagonal entries. The LF violating limits are satisfied in types I and III if \( N \) and \( \Sigma \) mainly mix with only one charged lepton family. In Table 2 we collect the bounds from EWPD on the \( N \) and \( \Sigma \) mixings with the SM leptons \( V_{\nu L,\Sigma} \), and in Table 3 their product including the LF violating bounds. These

<table>
<thead>
<tr>
<th>Coupling</th>
<th>Only with ( \epsilon )</th>
<th>Only with ( \mu )</th>
<th>Only with ( \tau )</th>
<th>Universal</th>
</tr>
</thead>
<tbody>
<tr>
<td>( V_{\nu N} = \frac{v(\lambda_1)}{\sqrt{2M_N}} )</td>
<td>&lt; 0.055</td>
<td>0.057</td>
<td>0.079</td>
<td>0.038</td>
</tr>
<tr>
<td>( V_{\nu \Sigma} = -\frac{v(\lambda_2)}{\sqrt{2M_\Sigma}} )</td>
<td>&lt; 0.019</td>
<td>0.017</td>
<td>0.027</td>
<td>0.016</td>
</tr>
</tbody>
</table>

values update and extend previous bounds on diagonal entries for \( N \) (see also). Their dependence on the model parameters entering in the operator coefficients in Table 1 is explicit in the first column of Table 2. All low energy effects are proportional to this mixing, and the same holds for the gauge and Higgs couplings between the new and the SM leptons, responsible of the heavy lepton decay (and \( N \) production if there is no extra NP). An interesting by-product of a non-negligible mixing of the electron or muon with a heavy \( N \) is that the fit to EWPD prefers a Higgs mass \( M_H \) higher than in the SM, in better agreement with the present direct limit. This is so because their contributions to the most significative observables partially cancel, so that
both the mixing and $M_H$ can be relatively large without spoiling the agreement with EWPD. The new 90% CL on $M_H$ increases in this case up to \( \sim 260 \text{ GeV} \) (see also\textsuperscript{25,26}). In all other cases the limit stays at \( \sim 165 \text{ GeV} \).

In type II see-saw a crucial phenomenological issue is the relative size of \( \langle \lambda \Delta \rangle_{ij} \) and \( \mu_{\Delta} \) for \( M_{\Delta} \sim 1 \text{ TeV} \). The \( \nu \) masses are proportional to their product, \( (m_{\nu})_{ij} = 2v^2 \frac{\mu_{\Delta} \langle \lambda \Delta \rangle_{ij}}{M_{\Delta}} \), which gives the strength of the LN violation. If \( \mu_{\Delta} \) is small enough, \( \langle \lambda \Delta \rangle_{ij} \) can be relatively large and saturate present limits on LF violating processes, eventually showing at the next generation of experiments. If instead \( \langle \lambda \Delta \rangle_{ij} \) are very small, the flavour structure appears only in the \( \nu \) mass matrix. The present limits are reviewed in\textsuperscript{15}. Neglecting LF violating bounds (i.e., assuming that \( \langle \lambda \Delta \rangle_{ee} \) is small enough not to give a too large $\mu \rightarrow eee$ decay rate), \( \mu_{\Delta} \) and \( \lambda_{\Delta} \) are constrained by the $T$ oblique parameter and four-fermion processes, respectively. From a global fit to EWPD (see\textsuperscript{20} for details on the data set used) we obtain the following limits at 90% CL:

\[
\frac{\mu_{\Delta}}{M_{\Delta}^2} < 0.048 \text{ TeV}^{-1}, \quad \frac{|\lambda_{\Delta} \delta_{\mu e}|}{M_{\Delta}} < 0.100 \text{ TeV}^{-1}.
\]

### 3 Dilepton signals of see-saw messengers

The previous limits apply to any particle transforming as the corresponding see-saw messenger, independently of whether it contributes or not to light neutrino masses. As indicated above, in minimal models the tight restriction imposed by \( \nu \) masses (Eq. 1) gives much more stringent limits for the mixings of TeV-scale see-saw messengers. However, these limits can be avoided if additional particles give additional contributions to neutrino masses that cancel the previous ones, for instance if the fermionic messengers are quasi-Dirac, i.e. a nearly degenerate Majorana pair with appropriate couplings\textsuperscript{27}. The EWPD limits are in this case relevant for production and detection of type I messengers $N$, but the signals are different because they conserve LN to a very large extent\textsuperscript{14,28}. On the other hand, type II and III messengers with masses near the TeV can be produced and detected at LHC even in minimal models. Let us discuss the three types of see-saw mechanism in turn.

#### 3.1 Type I: Fermion singlets $N$

As already explained, a type I heavy neutrino $N$ with a mixing saturating the EWPD limit cannot be Majorana, unless extra fields with a very precise fine tuning keep the \( \nu \) masses small enough\textsuperscript{29}. Unnatural cancellations allowing for LN-violating signals are also possible in principle. In this case a fast simulation shows that LHC can discover a Majorana neutrino singlet with $M_N \simeq 150 \text{ GeV}$ for $|V_{\mu N}| \geq 0.054$ (near the EWPD limit)\textsuperscript{8}, assuming an integrated luminosity $L = 30 \text{ fb}^{-1}$.

Such a signal can be also observed for much smaller mixings and larger masses if there is some extra NP\textsuperscript{30}, especially if the extra particles can be copiously produced at LHC\textsuperscript{31}. This is
the case, for instance, if the gauge group is left-right symmetric and the new $W'_R$ has a few TeV mass. Then $pp \rightarrow W'_R \rightarrow t\bar{t} \rightarrow t\ell\ell W$ is observable, even with negligible mixing $V_{bN}$, for $M_N$ and $M_{W'}$ up to 2.3 TeV and 3.5 TeV, respectively,\(^{32}\) for an integrated luminosity $L = 30 \ fb^{-1}$. Similarly, if the SM is extended with a leptophobic $Z'$, the process $pp \rightarrow Z' \rightarrow NN \rightarrow t\ell\ell WW$ can probe $Z'$ masses\(^{33}\) up to 2.5 TeV, and $M_N$ up to 800 GeV.

3.2 Type II: Scalar triplets $\Delta$

$SU(2)_L$ scalar triplets can be produced through the exchange of electroweak gauge bosons with SM couplings, and then they may be observable for masses near the TeV scale (see for reviews\(^{3,31}\)). Although suppressed, their decays can occur within the detector for these large masses. In Fig. 1-(II) we display one of the possible processes. The search strategy and LHC potential depend on the dominant decay modes. These are proportional to the $\Delta$ vacuum expectation value $| < \Delta^0 > | \equiv v_\Delta$, as for example\(^9\) $\Delta^{\pm \pm} \rightarrow W^{\pm}W^{\pm}$, or to $( \lambda_\Delta )_{ij}$, as\(^{11}\) $\Delta^{\pm \pm} \rightarrow t\pm t^{(\gamma)} \pm$. $\Delta^{\pm \pm}$ can also decay into $\Delta^{\pm}W^{\pm}$ if kinematically allowed (see\(^{10}\)). All these different decay channels make the phenomenological analysis of single and pair $\Delta^{\pm \pm}$ production quite rich\(^{12}\). The EWPD limit in Eq. 7 translates into the bound $v_\Delta = \frac{v^2 |\mu_\Delta|}{\sqrt{M_\Delta^2}} < 2$ GeV. This is to be compared with $|m_{\nu_l}| = 2\sqrt{2}v_\Delta |\lambda_\Delta| \sim 10^{-9}$ GeV, which gives a much more stringent constraint for non-negligible $\lambda_\Delta$. Dilepton (diboson) decays are dominant for $v_\Delta < (<) v'_\Delta$ $\sim 10^{-4}$ GeV. If for instance $\lambda_\Delta$ is of the same size as the charged lepton Yukawa couplings $\sim 10^{-2} - 5 \times 10^{-6}$, $v_\Delta$ varies from $5 \times 10^{-8}$ to $10^{-4}$ GeV, below the critical value $v'_\Delta$, and $\Delta$ decays mainly into leptons. In this case the LHC reach for $M_{\Delta^{\pm \pm}}$ has been estimated, based on statistics, to be $\sim 1$ TeV for an integrated luminosity $L = 300 \ fb^{-1}$. In Fig. 3 we plot the invariant mass distribution $m_{\ell\ell}$ of same-sign dilepton pairs containing the lepton of largest transverse momentum for $M_\Delta = 600$ GeV. As this fast simulation analysis shows, the SM background is well separated from the signal, and the LHC discovery potential strongly depends on the light neutrino mass hierarchy. For the simulated sample we find 4 (44) signal events for the normal $\nu$ mass hierarchy NH (inverted IH), well separated from the main backgrounds: $t\bar{t}nj$ (1007 events), $Z\bar{b}b\bar{n}j$ (91 events), $tW$ (68 events), and $Z\bar{t}\bar{t}nj$ (51 events). We get rid of other possible backgrounds like $ZZnj$ requiring no opposite-sign dilepton pairs with an invariant mass in the range $M_Z \pm 5$ GeV. For larger $v_\Delta$ values, with dominant non-leptonic decays, the corresponding reach estimate based on statistics is $\sim 600$ GeV. Note that only in the leptonic case LHC is sensitive to the see-saw flavour structure. Near the critical value, one could in principle extract information on the structure and on the global scale of the see-saw.

Tevatron Collaborations have already established limits on the scalar triplet mass assuming that $\Delta^{\pm \pm} \rightarrow l^\pm l^\pm$ 100 % of the time: At the 95 % CL $M_{\Delta^{\pm \pm}} > 150$ GeV for $\Delta^{\pm \pm}$ only decaying to muons\(^{34}\), and an integrated luminosity $L = 1.1 \ fb^{-1}$.

3.3 Type III: Fermion triplets $\Sigma$

Not so much attention has been payed to the study of the LHC reach for $SU(2)_L$ fermion triplets $\Sigma$. Up to very recently a similar electroweak process, the production of a heavy vector-like lepton doublet\(^{35}\), had to be used to guess that LHC could be sensitive to $M_{\Sigma} \sim 500$ GeV. A dedicated study\(^{13}\) estimates that an integrated luminosity $L = 10 \ fb^{-1}$ should allow to observe LN violating signals (see Fig. 1-(III) for a relevant process) for $M_{\Sigma} < 800$ GeV. Vector-like fermion triplets couple to SM leptons proportionally to its mixing $V_{bN}$, which is $\leq 10^{-6}$ according to Eq. 1 if $\Sigma$ is at the LHC reach $\sim 1$ TeV. So, one can eventually improve the analysis using the displaced vertex signatures of their decays.
Figure 3: Same-sign dilepton invariant mass distributions for $M_\Delta = 600$ GeV and normal (NH) and inverted (IH) $\nu$ mass hierarchies, assuming an integrated luminosity $L = 300$ fb$^{-1}$.

4 Conclusions

Same-sign dilepton signals $l^\pm l^{(')}\pm X$ will allow to set significative limits on see-saw messengers at LHC, as illustrated in Table 4. The estimates for $M_\Delta$ and $M_\Sigma$ are mainly based on statistics, and a more detailed analysis is needed to confirm them.

Table 4: LHC discovery limit estimates for see-saw messengers, assuming an integrated luminosity $L = 30,300$ and $10$ fb$^{-1}$ for $N$, $\Delta$ and $\Sigma$, respectively. See Section 3 for a detailed explanation.

<table>
<thead>
<tr>
<th>LHC reach (in GeV)</th>
<th>$M_N$</th>
<th>$M_\Delta$</th>
<th>$M_\Sigma$</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>150</td>
<td>600 – 1000</td>
<td>800</td>
</tr>
</tbody>
</table>

Acknowledgments

We thank Ll. Ametller, S. Bar-Shalom, C. Biggio, T. Hambye, A. Soni and J. Wudka for discussions. F.A. thanks the organizers of the Rencontres de Moriond EW 2008 meeting for the excellent organization and the warm hospitality. This work has been supported by MEC project FPA2006-05294 and Junta de Andalucía projects FQM 101, FQM 437 and FQM03048. J.A.A.S. and J.B. also thank MEC for a Ramón y Cajal and an FPU grant, respectively.

References

The Standard model Higgs as the inflaton

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We describe how non-minimal coupling term between the Higgs boson and gravity can lead to the chaotic inflation in the Standard Model without introduction of any additional degrees of freedom. Produced cosmological perturbations are predicted to be in accordance with observations. The tensor modes of perturbations are practically vanishing in the model.

1 Introduction

This talk is based on the recent work, and closely follows it. Note, that the expression for the inflationary potential presented here differs from the one presented in the original work—both expressions coincide in the region relevant for inflation, while the expression given here has a wider range of validity (down to the Standard Model regime).

The fact that our universe is almost flat, homogeneous and isotropic is often considered as a strong indication that the Standard Model (SM) of elementary particles is not complete. Indeed, these puzzles, together with the problem of generation of (almost) scale invariant spectrum of perturbations, necessary for structure formation, are most elegantly solved by inflation. The majority of present models of inflation require an introduction of an additional scalar—the “inflaton”. Inflaton properties are constrained by the observations of fluctuations of the Cosmic Microwave Background (CMB) and the matter distribution in the universe. Though the mass and the interaction of the inflaton with matter fields are not fixed, the well known considerations prefer a heavy scalar field with a mass $\sim 10^{13}$ GeV and extremely small self-interacting quartic coupling constant $\lambda \sim 10^{-13}$ for realization of the chaotic inflationary scenario. This value of the mass is close to the GUT scale, which is often considered as an argument in favour of existence of new physics between the electroweak and Planck scales.

It was recently demonstrated in that the SM itself can give rise to inflation, provided non-minimal coupling of the Higgs field with gravity. The spectral index and the amplitude of tensor
perturbations can be predicted and be used to distinguish this possibility from other models for inflation; these parameters for the SM fall within the 1σ confidence contours of the WMAP-5 observations
.

To explain our main idea, let us consider the Lagrangian of the SM non-minimally coupled to gravity,

\[ L_{\text{tot}} = L_{\text{SM}} - \frac{M^2}{2} R - \xi H^\dagger H , \]

where \( L_{\text{SM}} \) is the SM part, \( M \) is some mass parameter, \( R \) is the scalar curvature, \( H \) is the Higgs field, and \( \xi \) is an unknown constant to be fixed later. The third term in (1) is in fact required by the renormalization properties of the scalar field in a curved space-time background, so, in principle, it should be added to the usual SM Lagrangian with some constant. Here, we will analyse the situation with large non-minimal coupling parameter \( \xi \gg 1 \), but still not too large for the non-minimal term to contribute significantly to the Planck mass in the SM regime (\( H \sim v \)), i.e. \( \sqrt{\xi} \ll 10^{17} \). Thus, we have \( M \approx M_P = (8\pi G_N)^{-1/2} = 2.4 \times 10^{18} \text{ GeV} \).

It is well known that inflation has interesting properties in models of this type, but most attempts were made to identify the inflaton field with the GUT Higgs field. In this case one naturally gets into the regime of induced gravity (where, unlike this paper, \( M = 0 \) and \( M_P \) is generated from the non-minimal coupling term by the Higgs vacuum expectation value). In this case the Higgs field decouples from the other fields of the model, which is generally undesirable. Here we demonstrate, that when the SM Higgs boson is coupled non-minimally to gravity, the scales for the electroweak physics and inflation are separate, the electroweak properties are unchanged, while for much larger field values the inflation is possible.

The paper is organised as follows. We start from discussion of inflation in the model, and use the slow-roll approximation to find the perturbation spectra parameters. Then we will argue in Section 3 that quantum corrections are unlikely to spoil the classical analysis we used in Section 2. We conclude in Section 4.

## 2 Inflation and CMB fluctuations

Let us consider the scalar sector of the Standard Model, coupled to gravity in a non-minimal way. We will use the unitary gauge \( H = h/\sqrt{2} \) and neglect all gauge interactions for the time being, they will be discussed later in Section 3. Then the Lagrangian has the form:

\[ S_J = \int d^4x\sqrt{-g}\left\{ - \frac{M^2 + \xi h^2}{2} R + \frac{\partial_{\mu} h \partial^{\mu} h}{2} - \frac{\lambda}{4} (h^2 - v^2)^2 \right\} . \]

This Lagrangian has been studied in detail in many papers on inflation, we will reproduce here the main results of. Compared to we present a better approximation for the inflationary potential here. To simplify the formulae, we will consider only \( \xi \) in the region \( 1 \ll \sqrt{\xi} \ll 10^{17} \), in which \( M \approx M_P \) with very good accuracy.

It is possible to get rid of the non-minimal coupling to gravity by making the conformal transformation from the Jordan frame to the Einstein frame

\[ \hat{g}_{\mu\nu} = \Omega^2 g_{\mu\nu} , \quad \Omega(h)^2 = 1 + \frac{\xi h^2}{M_P^2} . \]

This transformation leads to a non-minimal kinetic term for the Higgs field. So, it is convenient to make the change to the new scalar field \( \chi \) with

\[ \frac{d\chi}{dh} = \sqrt{\Omega^2 + \frac{3}{2} M_P^2 \left( \frac{d(\Omega^2)}{dh} \right)^2} = \sqrt{1 + (\xi + 6\xi^2) \frac{h^2}{M_P^2}} . \]
Finally, the action in the Einstein frame is

$$S_E = \int d^4 x \sqrt{-\hat{g}} \left\{ -\frac{\hat{M}_P^2}{2} \hat{R} + \frac{\partial_{\mu} \chi \partial^{\mu} \chi}{2} - U(\chi) \right\},$$  \hspace{1cm} (5)

where $\hat{R}$ is calculated using the metric $\hat{g}_{\mu\nu}$ and the potential is

$$U(\chi) = \frac{1}{\Omega(h(\chi))^4} \frac{\lambda}{4} (h(\chi)^2 - v^2)^2.$$  \hspace{1cm} (6)

For small field values $h, \chi < \frac{M_P}{\xi}$ the change of variables is trivial, $h \approx \xi$ and $\Omega^2 \approx 1$, so the potential for the field $\chi$ is the same as that for the initial Higgs field and we get into the SM regime. For $h, \chi \gg \frac{M_P}{\xi}$ the situation changes a lot. In this limit the variable change (4) is

$$\Omega(h)^2 \approx \exp \left( \frac{2 \chi}{\sqrt{6} M_P} \right).$$  \hspace{1cm} (7)

The potential for the Higgs field is exponentially flat for large $\xi$ and has the form

$$U(\chi) = \frac{\lambda M_P^4}{4 \xi^2} \left( 1 - \exp \left( -\frac{2 \chi}{\sqrt{6} M_P} \right) \right)^2.$$  \hspace{1cm} (8)

The full effective potential in the Einstein frame is presented in Fig. 1. It is the flatness of the potential at $\chi \gtrsim M_P$ which makes the successful (chaotic) inflation possible.

Basically, there are two distinct scales— for low field values $h, \chi \ll \frac{M_P}{\xi}$ we have the SM, for high field values $h \gg \frac{M_P}{\sqrt{\xi}} (\chi > M_P)$ we have inflation with exponentially flat potential (8) and the Higgs field is decoupled from all other SM fields (because $\Omega \propto h$, see Section 3). In the intermediate region $M_P/\xi \ll h \ll M_P/\sqrt{\xi}$ ($M_P/\xi \ll \chi < M_P$) the coupling with other particles is not suppressed ($\Omega \sim 1$), while the potential and change of variables are still given by (8) and (7).

Analysis of the inflation in the Einstein frame can be performed in the standard way using the slow-roll approximation. The slow roll parameters (in notations of $^{23}$) can be expressed

\begin{itemize}
\item The following two formulae have wider validity range than those in $^1$, which are valid only for $h \gg \frac{M_P}{\sqrt{\xi}}$.
\item The same results can be obtained in the Jordan frame $^{21,22}$.
\end{itemize}
analytically as functions of the field $h(\chi)$ using (4) and (6) (we give here the expressions for the case $\epsilon^2 \gg M_p^2/\xi \gg v^2, \xi \gg 1$, exact expressions can be found in\textsuperscript{16}),

$$
\epsilon = \frac{M_p^2}{2} \left( \frac{dU/d\chi}{U} \right)^2 \approx \frac{4M_p^4}{3\xi^2 h^4},
$$

$$
\eta = \frac{M_p^2}{2} \frac{d^2U/d\chi^2}{U} \approx \frac{4M_p^4}{3\xi^2 h^4} \left( 1 - \frac{\xi h^2}{M_p^2} \right),
$$

$$
\zeta^2 = \frac{M_p^2}{2} \frac{d^3U/d\chi^3}{U^2} \approx \frac{16M_p^4}{9\xi^3 h^6} \left( \frac{\xi h^2}{M_p^2} - 3 \right).
$$

Slow roll ends when $\epsilon \simeq 1$, so the field value at the end of inflation is $h_{\text{end}} \simeq (4/3)^{1/4} M_p/\sqrt{\xi} \simeq 1.07 M_p/\sqrt{\xi}$. The number of $e$-foldings for the change of the field $h$ from $h_0$ to $h_{\text{end}}$ is given by

$$
N = \int_{h_0}^{h_{\text{end}}} \frac{M_p^2}{U} \frac{dU/d\chi}{dh} \left( \frac{d\chi}{dh} \right)^2 dh \approx \frac{3}{4} \frac{h_0^2 - h_{\text{end}}^2}{M_p^2/\xi}.
$$

We see that for all values of $\sqrt{\xi} \ll 10^{17}$ the scale of the Standard Model $v$ does not enter in the formulae, so the inflationary physics is independent on it.

After end of the slow roll the $\chi$ field enters oscillatory stage with diminishing amplitude. After the oscillation amplitude falls below $M_p/\xi$, the situation returns to the SM one, so at this moment the reheating is imminent due to the SM interactions, which guarantees the minimum reheating temperature $T_{\text{reh}} \gtrsim (\frac{15\lambda}{8\pi^2 g^*})^{1/4} M_p/\xi \simeq 1.5 \times 10^{13}$ GeV, where $g^* = 106.75$ is the number of degrees of freedom of the SM. Careful analysis may give a larger temperature generated during the decay of the oscillating $\chi$ field, but definitely below the energy scale at the end of the inflation $T_{\text{reh}} < (\frac{2\lambda}{\pi^2 g^*})^{1/4} M_p/\sqrt{\xi} \simeq 2 \times 10^{15}$ GeV.

As far as the reheating mechanism and the universe evolution after the end of the inflation is fixed in the model, the number of $e$-foldings for the the COBE scale entering the horizon can be calculated (see\textsuperscript{23}). Here we estimate it as $N_{\text{COBE}} \simeq 62$ (exact value depends on the detailed analysis of reheating, which will be done elsewhere). The corresponding field value is $h_{\text{COBE}} \simeq 9.4 M_p/\sqrt{\xi}$. Inserting (12) into the COBE normalization $U/\epsilon = (0.027 M_p)^4$ we find the required value for $\xi$

$$
\xi \simeq \sqrt{\lambda \frac{N_{\text{COBE}}}{3 \times 0.027^2}} \simeq 49000\sqrt{\lambda} = 49000 \frac{m_H}{\sqrt{2} v}.
$$

Note, that if one could deduce $\xi$ from some fundamental theory this relation would provide a connection between the Higgs mass and the amplitude of primordial perturbations.

The spectral index $n_s = 1 - 6\epsilon + 2\eta$ calculated for $N = 60$ (corresponding to the scale $k = 0.002$ Mpc) is $n_s \simeq 1 - 8(4N + 9)/(4N + 3)^2 \simeq 0.97$. The tensor to scalar perturbation ratio\textsuperscript{9} is $r = 16\epsilon \simeq 192/(4N + 3)^2 \simeq 0.0033$. The predicted values are well within one sigma of the current WMAP measurements\textsuperscript{9}, see Fig. 2.

3 Radiative corrections

An essential point for inflation is the flatness of the scalar potential in the region of the field values $h \sim 10 M_p/\sqrt{\xi}$ ($\chi \sim 6 M_p$). It is important that radiative corrections do not spoil this property. Of course, any discussion of quantum corrections is flawed by the non-renormalizable character of gravity, so the arguments we present below are not rigorous.

\footnote{These formulas are valid up to the end of the slow roll regime $h_{\text{end}}$, while the formulas (10) and (11) in\textsuperscript{1} are applicable only for the earlier inflationary stages, $h^2 \gg M_p^2/\xi$, which is sufficient to calculate primordial spectrum parameters $n_s$ and $r$.}
There are two qualitatively different type of corrections one can think about. The first one is related to the quantum gravity contribution. It is conceivable to think\textsuperscript{24} that these terms are proportional to the energy density of the field $\chi$ rather than its value and are of the order of magnitude $U(\chi)/M_P^4 \sim \lambda/\xi^2$. They are small at large $\xi$ required by observations. Moreover, adding non-renormalizable operators $h^{4+2n}/M_P^{2n}$ to the Lagrangian (2) also does not change the flatness of the potential in the inflationary region.\textsuperscript{d}

Other type of corrections is induced by the fields of the Standard Model coupled to the Higgs field. In one loop approximation these contributions have the structure

$$\Delta U \sim \frac{m^4(\chi)}{64\pi^2} \log \frac{m^2(\chi)}{\mu^2},$$

where $m(\chi)$ is the mass of the particle (vector boson, fermion, or the Higgs field itself) in the background of field $\chi$, and $\mu$ is the normalization point. Note that the terms of the type $m^2(\chi)M_P^2$ (related to quadratic divergences) do not appear in scale-invariant subtraction schemes that are based, for example, on dimensional regularization (see a relevant discussion in\textsuperscript{25},\textsuperscript{26},\textsuperscript{27},\textsuperscript{28}). The masses of the SM fields can be readily computed\textsuperscript{13} and have the form

$$m_{\psi,A}(\chi) = \frac{m(v)}{v} \frac{h(\chi)}{\Omega(\chi)}, \quad m_{H}^2(\chi) = \frac{d^2U}{d\chi^2},$$

for fermions, vector bosons and the Higgs (inflaton) field. It is crucial that for large $\chi$ these masses approach different constants (i.e. the one-loop contribution is as flat as the tree potential) and that (14) is suppressed by the gauge or Yukawa couplings in comparison with the tree term. In other words, one-loop radiative corrections do not spoil the flatness of the potential as well. This argument is identical to the one given in\textsuperscript{13}.

### 4 Conclusions

Non-minimal coupling of the Higgs field to gravity leads to the possibility of chaotic inflation in SM. Specific predictions for the primordial perturbation spectrum are obtained. Specifically, very small amount of tensor perturbations is expected, which means that future CMB experiments measuring the B-mode of the CMB polarization (PLANCK) can distinguish between the described scenario from other models (based, e.g. on inflaton with quadratic potential).

At the same time, we expect that the Higgs potential does not enter into the string coupling regime, nor generates another vacuum up to the scale of at least $M_P/\xi \sim 10^{14}$ GeV, so we expect the Higgs mass to be in the window $130 \text{GeV} < M_H < 190 \text{GeV}$ (see, eg.\textsuperscript{29}), otherwise the inflation would be impossible.

The inflation mechanism we discussed has in fact a general character and can be used in many extensions of the SM. Thus, the $\nu$MSM of\textsuperscript{30,31,32,33,25,34,35,36,37,38,39,40} (SM plus three light fermionic singlets) can explain simultaneously neutrino masses, dark matter, baryon asymmetry of the universe and inflation without introducing any additional particles (the $\nu$MSM with the inflaton was considered in\textsuperscript{25}). This provides an extra argument in favour of absence of a new energy scale between the electroweak and Planck scales, advocated in\textsuperscript{27}.

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\textsuperscript{d}Actually, in the Jordan frame, we expect that higher-dimensional operators are suppressed by the effective Planck scale $M_P^2 + \xi h^2$. 
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UNPARTICLES,
A VIEW FROM THE HIGGS WINDOW

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In this Moriond talk I give an introduction to what unparticles are supposed to be and then show how coupling unparticles to the Higgs opens a window to the exploration of unparticle sectors. The second part is based on work done in collaboration with Antonio Delgado, Mariano Quirós and José Miguel No in 1,2.

1 Introduction to Unparticles

We start by assuming (with Georgi 3) that our theory contains a scale invariant sector. Such sector must look very different from the familiar Standard Model (SM). In the SM we have an explicit mass term in the Higgs potential, the origin of all mass scales. We cannot have such mass terms in the scale invariant sector so, one would naively think such sector is composed of massless fields. That is too naive as we will see later on.

In fact, not all SM masses come from the Higgs mass term: quantum running of couplings can generate mass scales as happens in QCD. So, another requirement of the scale invariant sector is that couplings there should not run at all, and this should happen in a non-trivial way, i.e. for non-zero coupling. In fact, there are examples in the literature of just that happening, like the Banks-Zaks (BZ) model4: a Yang-Mills theory with some particular choice of the number of colors and flavors. The gauge coupling reaches an IR fixed-point below some mass scale ΛU. This requires cancellations between different orders of perturbation theory and normally occurs at strong coupling. Below ΛU the theory becomes scale-invariant.

In order to explore the physics of such scale-invariant sector, Georgi considered the possibility of coupling it to the SM, e.g. assume there is a heavy sector, at the scale Mm, coupled both to the SM and the Banks-Zaks-like sector, see figure 1. The effective theory below this mass (Mm) will contain non-renormalizable interactions between SM operators and BZ operators:

\[ \frac{1}{M_m^8} O_{SM} O_{BZ} . \] (1)
In the BZ side interesting things happen when one reaches the scale-invariant fixed point below \( \Lambda_U \). Operators with some scaling dimension \( d_{BZ} \) get transformed into operators with dimension \( d_U \) that can be very different from an integer value.

\[
\mathcal{O}_{BZ} \rightarrow \Lambda_U^{d_{BZ}-d_U} \mathcal{O}_U .
\]  

(2)

For simplicity I only consider scalar operators, in which case unitarity demands \( d_U > 1 \) and I will take \( d_U < 2 \) because the most interesting effects take place in that range. The non-renormalizable term will now be a coupling between SM fields and unparticle operators:

\[
\frac{\Lambda_U^{d_{BZ}-d_U}}{M_m^{k}} \mathcal{O}_S \mathcal{O}_U .
\]  

(3)

One assumes that the parameter \( \epsilon \equiv \Lambda_U^{d_{BZ}-d_U}/M_m^{k} \) is small enough so that unparticle effects haven’t showed up so far (in fact, \( \epsilon \) is not dimensionless so the appropriate power of the low-energy scale relevant for the process in question should be included). The important point is that, \( d_U \) being a non-integer, \( \mathcal{O}_U \) cannot be interpreted in terms of particles in the usual way.

Once we have a coupling between SM fields and unparticles we can start discussing unparticle production, using this coupling as an insertion in some SM process. To first order in \( \epsilon \) one simply gets missing energy and momentum. The probability for such a process goes like \( \epsilon^2 \):

\[
\epsilon^2 |\langle SM_{out}|\mathcal{O}_S|SM_{in}\rangle|^2 |\langle U|\mathcal{O}_U|0\rangle|^2 ,
\]  

(4)

and the unparticle matrix element will be determined, by scale-invariance, to go like the appropriate power of \( x \):

\[
\langle 0|\mathcal{O}_U(x)\mathcal{O}_U^\dagger(0)|0\rangle = \int e^{-ipx} |\langle U,p|\mathcal{O}_U(0)|0\rangle|^2 p(p^2) \frac{d^4p}{(2\pi)^4} \sim x^{-2d_U} .
\]  

(5)

Going to momentum space we find an unparticle two-point correlator (or propagator) that goes like this unusual power of momentum: \( 3, 5 \)

\[
|\langle U,p|\mathcal{O}_U(0)|0\rangle|^2 = A_{d_U} \theta(p^0)\theta(p^2)(p^2)^{d_U-2} ,
\]  

(6)
(\(d_U \to 1\) would correspond to a normal particle propagator). Georgi compared this to the phase space factor for production of \(n\) massless particles (\(dLIPS_n = A_n s^{n-2}\)) to conclude that unparticles could be interpreted as a non-integer number of massless particles but I think this does not help intuition much.

I find more useful to obtain the spectral function representation of the unparticle propagator: it gives a scan of what is being propagated. For a normal particle the spectral function is just a Dirac delta at the particle’s mass, \(\rho(s) = \delta(s - m^2)\). For unparticles we get a continuous function, \(\rho(s) = s^{d_U-2}\), see figure 2. This leads to the following picture: the operator \(\mathcal{O}_U\) does not create particles out of the vacuum but rather a non-localized wave over the full range of \(p^2\).

Another useful way of looking at unparticles was proposed by Stephanov. He considered an infinite tower of scalars \(\varphi_n\), \((n = 1, \ldots, \infty)\), with masses-squared separated by a constant splitting \(\Delta^2, M_n^2 = \Delta^2 n\). In the continuum limit, \(\Delta^2 \to 0\), it is simple to show that one gets a scale-invariant spectrum. In fact, in the deconstructed action

\[
S = \int d^4x \sum_{n=1}^{\infty} \left[ \frac{1}{2} (\partial_{\mu} \varphi_n)^2 + \frac{1}{2} M_n^2 \varphi_n^2 \right],
\]

kinetic terms are scale invariant [under \(\varphi_n(x) \to \lambda \varphi_n(\lambda x)\)] while mass terms are not. However, when \(\Delta^2 \to 0\), replacing the discrete sum by an integral and taking \(\varphi_n/\Delta \to u(M^2)\), the continuum action including the mass term

\[
S = \int d^4x \int_0^{\infty} dM^2 \left[ \frac{1}{2} (\partial_{\mu} u)^2 + \frac{1}{2} M^2 u^2 \right],
\]

is indeed scale invariant [under \(u(M^2, x) \to u(M^2/\lambda^2, \lambda x)\)] thanks to the rescaling freedom in the integration variable. This is one example of scale invariant model not composed of massless fields. To make contact with unparticle operators one defines a linear combination \(O\) of the \(\varphi_n\) fields

\[
O = \sum_n F_n \varphi_n ,
\]

with the appropriate mass-dependent coefficients

\[
F_n^2 = \frac{A_{d_U}}{2\pi} \Delta^2 (M_n^2)^{d_U-2} ,
\]
such that the two-point correlator of the deconstructed operator \( O \) has the correct unparticle continuum limit. In this way Stephanov could rederive many of the phenomenological unparticle implications working in the more familiar deconstructed version and taking the continuum limit at the end.

2 Looking at Unparticles through the Higgs

In the second part of the talk I want to discuss the possibility of having a direct coupling between the unparticle operator and the Higgs:

\[
\kappa_U \mathcal{O}_U = \kappa_U |H|^2 \mathcal{O}_U .
\]

Notice that the Higgs is special in the sense that it offers the only possible renormalizable coupling with scalar unparticles\(^7\). One immediate difficulty one has to face is that, after electroweak symmetry breaking (EWSB), such coupling induces a tadpole for \( \mathcal{O}_U \) and therefore a VEV for it. It is a simple matter to obtain this VEV: simply get the VEVs \( v_n \) for the deconstructed scalars and integrate over \( M^2 \) to obtain:

\[
\langle \mathcal{O}_U \rangle = -\frac{\kappa_U v^2}{2} \int_0^\infty \frac{F^2(M^2)}{M^2} dM^2 \propto \int_0^\infty (M^2)^{d_U-3} dM^2 ,
\]

where \( F(M^2) \) is the continuum version of \( F_n \). For \( d_U < 2 \) one sees that this integral has an IR divergence. Its origin is clear: the tadpole goes to \( \infty \) when \( M_n \to 0 \) while the restoring term, which is \( M_n \) itself, goes to zero then.

However, this IR problem is easy to cure. For instance, if the Higgs also couples to quadratic terms for the \( \varphi_n \)’s

\[
\zeta |H|^2 \sum_n \varphi_n^2 ,
\]

this will give an extra mass squared \( \zeta v^2 \) to unparticles that will act as an IR cutoff in the previous integral:

\[
\langle \mathcal{O}_U \rangle = -\frac{\kappa_U v^2}{2} \int_0^\infty \frac{F^2(M^2)}{M^2 + \zeta v^2} dM^2 .
\]

Other more intriguing possibility uses quartic self-interactions of unparticles of the form

\[
\delta V = \xi \left( \sum_n \varphi_n^2 \right)^2 ,
\]

but I do not have time to discuss it here, see\(^2\). The structure of the unparticle continuum after the IR problem is solved is therefore a continuum above a mass gap \( m_g = \sqrt{\zeta v} \) (of EW size). Needless to say, the presence of such mass gap will affect dramatically the phenomenological (and constraints on) unparticles.

Let me focus here on some implications for Higgs physics. After EWSB the Higgs scalar will mix with the unparticle continuum. The mixed mass matrix in the deconstructed picture looks like this:

\[
\mathcal{M} = \begin{bmatrix}
  m_{h0}^2 & A_1 & \ldots & A_n & \ldots \\
  A_1 & M_1^2 + m_g^2 & \ldots & 0 & \ldots \\
  \vdots & \vdots & \ddots & \vdots & \ddots \\
  A_n & 0 & 0 & M_n^2 + m_g^2 & 0 \\
  \vdots & \vdots & \vdots & \vdots & \ddots 
\end{bmatrix}
\]

(\( \mathcal{M}^2 \))

Here \( m_{h0}^2 \) is the SM Higgs mass squared. Along the diagonal we have the unparticle tower with a mass gap and the first row and column are non-zero with \( A_n = v(\kappa_U F_n + 2\zeta v_n) \) and mix both
sectors. To analyze what happens in the presence of such mixing we derive the propagator for $h$ resumming $U$ insertions to obtain the following:

$$iP(p^2)^{-1} = p^2 - m_h^2 + v^2(\mu_U^2)^{2-d_U} \int_0^\infty \frac{(M^2)^{d_U-2}}{M^2 + m_g^2 - p^2} \left[ \frac{M^2}{M^2 + m_g^2} \right]^2 dM^2, \quad (17)$$

a propagator with a SM part and a more complicated term coming from unparticles. The same propagator can be obtained by diagonalizing the full propagator matrix. Here we have used

$$\langle \mu_U^2 \rangle^{2-d_U} \equiv \kappa_U^2 \frac{A_{d_U}}{2\pi}. \quad (18)$$

The first effect is a shift in the Higgs mass $m_h^2 \rightarrow m_h^2$, with $P(m_h^2)^{-1} = 0$. One can consider two qualitatively different cases depending on whether the shifted Higgs mass is below or above the mass gap. Let us start with a Higgs below the gap. Again it is quite useful to go to the spectral function representation of the modified propagator (17). It looks as shown in fig. 3: a Dirac delta at the pole Higgs mass and a continuum above $m_g^2$. In this plot we have chosen $m^2 = 0$, $\zeta = 1$, $d_U = 1.2$ and $\mu_U^2 = \mu_v^2 = v^2[A_{d_U}/(2\pi)]^{2-d_U}$. Explicitly the spectral function is

$$\rho(s) = \frac{1}{K^2(m_h^2)} \delta(s - m_h^2) + \theta(s - m_g^2) \frac{Q_H^2(s)}{[iP(s)^{-1}]^2 + \pi^2 Q_H^2(s)}, \quad (19)$$

where the functions $K^2(m_h^2)$ and $Q_H^2(s)$ can be found in $^1$. This spectral function has a very direct physical interpretation. For a given $s$ it gives the projection of the state with mass-squared $s$ onto the interaction eigenstate $h$:

$$\rho(s) \equiv \langle h|s|h \rangle = |\langle h|H \rangle|^2 \delta(s - m_h^2) + \theta(s - m_g^2)|\langle h|U, s \rangle|^2, \quad (20)$$

where $|h\rangle$, $|u, s\rangle$ are interaction eigenstates and $|H\rangle$, $|U, s\rangle$ are mass eigenstates. This is crucial for instance in studying the ZZ coupling as $h$ is the state that couples to ZZ. In this way we learn that the prefactor $R_h^2 = |\langle h|H \rangle|^2 = 1/K^2(m_h^2)$ of the delta function gives how much of the isolated pole is a pure Higgs. On the other hand, the spectral function above the mass gap,
Figure 4: Spectral function $\rho$ as a function of $s$ for a case with $m_h > m_g$. The parameters are $\mu^2_U = \mu^2_v$, $m^2 = 0$, $d_U = 1.2$ and $\zeta = 0.2$. All dimensions scaled by $m_g^2$.

$R_h^2(M^2) = |\langle h|U(s)\rangle|^2$, gives information about how the unparticle continuum couples to the $Z$ through Higgs mixing. Of course the following sum rule

$$R_h^2 + \int_0^\infty R_U^2(M^2)dM^2 = 1,$$

(21)

that implies that no “Higgsness” is lost in the mixing, is satisfied.

The case with a Higgs mass above the mass gap is also interesting. From the spectral function we learn that there is no delta function associated with the Higgs pole but rather a broad resonance due to the mixing of the Higgs with unparticles:

$$\rho(s) = \theta(s - m^2_h) \frac{Q^2_U(s)}{[tP(s) - 1]^2 + \pi^2 Q^2_U(s)}.$$  

(22)

In fact the Higgs is totally swallowed by this continuum, see fig. 4 (same parameters as in fig. 3 except for $\zeta = 0.2$). How big the width can be is shown by plot 5, where we show (as a function of $d_U$) the width, given by:

$$\Gamma_h = \frac{\pi Q^2_U(m^2_h)}{m_h K^2(m^2_h)},$$

(23)

and also the Higgs mass. We see that in some cases the width can be as large as the mass itself. Such effects will totally modify the expected Higgs collider phenomenology.

3 Conclusions

While almost everybody agrees that the idea of unparticles is extremely speculative it remains an intriguing theoretical possibility and we should keep it in mind for LHC. In this talk I focussed on the effects that one would expect from a direct coupling between the Higgs and an unparticle scalar operator. After dealing with an IR problem easy to solve I discussed how a mass gap is generated from EWSB. This gap would have important implications for unparticle phenomenology. I also showed how Higgs-unparticle mixing can greatly modify Higgs properties and how it can also give us a very interesting handle to explore the unparticle sector.
Figure 5: Width $\Gamma_h$ (from unparticle merging) and mass $m_h$ of the Higgs boson as a function of $d_U$ for $\mu_U^2 = \mu_v^2$, $m^2 = 0$ and $\zeta = 0.2$.

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References

Commissioning of CMS and early standard model measurements with jets, missing transverse energy and photons at the LHC

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We report on the status and history of the CMS commissioning, together with selected results from cosmic-muon data. The second part focuses on strategies for optimizing the reconstruction of jets, missing transverse energy and photons for early standard model measurements at ATLAS and CMS with the first collision data from the Large Hadron Collider at CERN.

1 CMS Commissioning

With the first collisions from the LHC expected soon, the CMS experiment at CERN has entered the final stage of commissioning. Nearly all of the detectors have been installed in the experiment and the last heavy structure of CMS was lowered into the experimental hall in January 2008 (see photograph in Fig. 1). The CMS collaboration has launched over the past years a series of combined data-taking exercises, so-called Global Runs, with increasing scope and complexity. More and more components have been integrated with the trigger and DAQ systems, and data from cosmic muons as well as high-rate random triggers have been used to prove readiness for LHC collisions.

One of the first highlights in the commissioning of the CMS detector was the "Magnet Test and Cosmic Challenge" (MTCC)\(^2\), which took place in 2006 when the CMS detector was operated in the assembly hall on the surface. The superconducting magnet of CMS required testing before lowering, providing a unique opportunity to operate all the subdetectors and sub-systems together and to take data with cosmic-ray muons. The participating systems included a 60° sector of the muon system, comprising gas detectors like the drift tubes (DTs), Cathode Strip Chambers (CSCs) and Resistive Plate Chambers (RPCs), both in the barrel part and endcaps of CMS. The tracking system comprised elements of the Silicon-Strip Tracker, and parts of the Electromagnetic (ECAL) and Hadron Calorimeter (HCAL) detected energy depositions of...
the traversing muons. In the second phase of the MTCC, the ECAL modules and the tracker elements were replaced by a specially designed mapping device to measure the three components of the magnetic field in the volume of the inner detectors with high precision, while HCAL and the muon systems continued to record cosmic-ray data.

Only six months after the conclusion of the MTCC, global data taking was resumed in the first of a series of Global Runs, this time in the underground experimental hall. With the end of 2007, CMS has recorded data from synchronized cosmic triggers from parts of all trigger detectors (CSCs, DTs, RPCs, ECAL and HCAL). The diagram on the right of Fig. 1 illustrates the progress made in the integration of subdetectors and subsystems in various Global Runs in 2007. One of the many results from these exercises is the confirmation of a single-hit resolution \( \delta x < 280 \mu m \) along \( \phi \) of the Muon Drift Tubes. One of the next goals before the LHC startup is the integration of the central tracker in the Global Runs.

2 Early Standard Model Measurements with Jets, Missing Transverse Energy and Photons

Already the first 10 to 100 pb\(^{-1}\) of recorded data at the LHC experiments will allow QCD measurements with minimum-bias events and will open the window to the Z- and W-boson and top-quark production. However, it is natural to expect that the first \( \sim 100 \text{pb}^{-1} \) of integrated luminosity will first be used to test and improve the understanding of the detector response and to calibrate and measure the performance of the physics objects used in the various searches and measurements to follow. Therefore, this article focuses on the strategies for the “commissioning” of the physics objects with the first data, using the standard model processes accessible with the recorded data set. Details about the longer term physics plans of ATLAS and CMS can be found in the references\(^1,3\).

Physicists from ATLAS and CMS are taking care that the experiments will make the best use of the early data to align and calibrate the detectors. Analyses are prepared to study and optimize the performance of the triggers and of physics objects. The strategies range from transverse-momentum balancing to exploiting well-known SM signatures, such as \textit{tag-and-probe}.

\( \phi \) is the azimuth angle of the CMS coordinate system with the \( z \) axis pointing in the direction of the beam and the origin in the center of the detector. \( \eta \) denotes the pseudorapidity \( \eta = -\ln[\tan(\theta/2)] \) where \( \theta \) is the polar angle of the CMS coordinate system.
with leptonically Z-boson decays. As the field is too large to cover in this article all the aspects of measurements with jets, electrons, photons and missing transverse energy, we report on a few selected recent studies from ATLAS and CMS with a focus on the reconstruction performance and strategies.

One of the first measurements to be performed at the LHC aims at the understanding of the spectrum of charged hadrons. These spectra have never been explored in hadron collisions at such high energies ($\sqrt{s} > 2$ TeV) and they are an important tool for the calibration and understanding of the detector response. Recent studies at CMS show that it is feasible to distinguish charged pions, kaons and protons with momenta up to 2 GeV/c and individually measure their spectra (see Fig. 2)\(^4\). Cross-sections and differential yields of charged particles (unidentified or identified pions, kaons and protons), produced in inelastic proton-proton collisions at $\sqrt{s} = 14$ TeV, can be measured with good precision with the CMS Pixel vertex detector and tracker system.

An important ingredient for the correct simulation of standard model processes at LHC is the understanding of the "underlying event", consisting of the beam-beam remnants (a soft component coming from the break-up of the two beam hadrons). Furthermore, it is sensitive to test multiple-parton interaction (MIP) tunes of QCD\(^5\). Technically, it is impossible to fully separate the underlying event from the hard scattering process, however, the observation of charged particles in the region "perpendicular" to the leading jet ($60^\circ < \Delta \phi_j < 120^\circ$ in the transverse projection) in QCD di-jet events can be used to distinguish between various MIP tunes that have been considered: DW\(^6\), DWT\(^7\), S0\(^8\) and Herwig\(^9\) (see Fig. 3)\(^10\).

Up to now, the analyses of the two experiments reconstruct jets mainly with variations of the iterative cone algorithm\(^11\) that forms jets from energy depositions in calorimeter towers in a cone of fixed size $\Delta R = \sqrt{\Delta \phi^2 + \Delta \eta^2}$. Calorimeter towers are the energy sums measured in the various depths of the calorimeter along $r$ at fixed $\eta$ and $\phi$. Even though iterative cone algorithms are well established and fast, which makes them particularly useful for triggering, new algorithms have been studied to improve infrared- and collinear-safety. Among those alternative algorithms are the Seedless Infrared-Safe Cone\(^12\) and Fast-$k_T$\(^13\) algorithms.

Rather than relying on the Monte Carlo simulation, the jet energy corrections can be extracted from transverse-energy asymmetry measured in di-jet events, a technique developed at the Tevatron\(^14\). In $2 \rightarrow 2$ events, transverse momenta of two jets are equal. This property can be used to scale the jet transverse momentum ($p_T$) at a given $\eta$ to a jet $p_T$ in a reference $\eta$.

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**Figure 2:** Left: Energy loss $dE/dx$ spectra for charged pions, kaons (K) and protons (p) in the Silicon Tracker (Pixels and Strips combined) of CMS. The scatter plot clearly shows the three separate bands for charged pions, kaons and protons (from left to right). Right: Charged-hadron spectra $dN/d\eta$ of CMS. The scatter plot clearly shows the three separate bands for charged pions, kaons and protons. Right: Separation power as a function of momentum for the CMS Preliminary simulation.
Figure 3: Densities $dN/dy/d\phi$ of charged particle multiplicity as a function of the azimuthal distance to the leading charged jet direction $\Delta \phi$ (left). The "transverse" region used for the measurement of the underlying event is defined as $60^\circ < |\Delta \phi| < 120^\circ$. The data points are shown for different trigger conditions (minimum-bias (MB) or single-jet calorimetric $p_T$ thresholds), sensitive to different ranges of jet momenta. Right: Densities $dN/dy/d\phi$ for tracks in the transverse region with $p_T > 0.9 \text{ GeV}/c$, as a function of the transverse momentum of the leading charged jet. The simulated "data" (open squares) correspond to an integrated luminosity of $100 \text{ pb}^{-1}$.

Figure 4: Jet energy resolution from a data-driven correction, based on a study of transverse energy asymmetry in di-jet events in CMS (left). ATLAS study of the noise contribution to $R = 0.7$ cone jets made from 3D clusters compared to standard jets based on calorimeter towers in QCD di-jet events (right).

region, in order to correct for the variation of the jet energy response as a function of $\eta$ based on data. However, as most events in the QCD di-jet samples have more than two reconstructed jets, which primarily originate from soft radiation and which degrade the momentum balance in the transverse plane between the two leading jets, the resolution obtained from the asymmetry is studied as a function of the maximum $p_T$ of the third jet. The final resolution is then extracted from the extrapolation of the $p_T$ of the third jet in the limit $p_T(3^{\text{rd}} \text{jet}) \rightarrow 0$. The left plot in Fig. 4 shows how well this data-driven method compares to Monte-Carlo truth as studied by CMS.

It has been shown for the ATLAS experiment, that the noise contribution in jets can be reduced when using topological cell clusters as inputs to the jet reconstruction rather than the total deposited energies in the calorimeter towers. The improvement from this attempt to reconstruct three-dimensional energy depositions is shown in Fig. 4 (right). As a result of the noise reduction shown in this figure, the angular and energy resolution and the jet efficiency is expected to improve for relatively low-momentum jets (below $p_T = 40 \text{ GeV}/c$).

It has also been shown that a clear signal of top-quark pair events, where one W boson from the top-quark decay produces a lepton (electron or muon) and a neutrino, while the other W boson decays into two quarks ($t\bar{t} \rightarrow Wb + W\bar{b} \rightarrow l\nu bq\bar{q}b$), can be extracted from the first $100 \text{ pb}^{-1}$ without the use of missing transverse energy or $b$-quark flavor tagging. With the help of the W-mass constraint, such a sample will be exploited to improve the jet-energy scale
The missing transverse energy ($\slashed{E}_T$) is one of the most complex physics objects at hadron collider experiments, because it combines information from all sub-detectors and requires homogeneous calorimeter coverage in the pseudorapidity region of $|\eta| < 5$. A good understanding of the measurement of $\slashed{E}_T$ is important, as it is present in signatures of physics beyond the standard model, but also required for the reconstruction of e.g. leptonically decaying W bosons. Various contributions of fake $\slashed{E}_T$ can degrade the performance: machine background from the accelerator and beam-gas interactions, noisy or dead calorimeter cells or regions, non-linearity in the hadronic response and finite energy resolution. Figure 5 illustrates an example on how to make analyses less sensitive to fake $\slashed{E}_T$: ATLAS has recently studied the reduction of the fake $\slashed{E}_T$ contribution from muons escaping the fiducial region of the detector by requiring angular separation of the direction of the missing transverse energy and high-momentum jets.\textsuperscript{15}

The first data recorded at the LHC experiments will be used to establish the calibration and uniformity of the response of the electromagnetic calorimeters. At first, the calibration can be obtained from the analysis of neutral pions. With increasing luminosity, events with Z bosons decaying into an electron and a positron will become more important, as they open the probed range to higher energies and provide the kinematical constraint from the well known mass of the Z boson. Figure 6 shows the expected improvements in the ECAL uniformity as a function of recorded data for ATLAS. The figure shows that for a relatively small data set of 100 pb$^{-1}$ integrated luminosity, the constant term in the energy resolution can be greatly improved and that the plateau can be reached for $> 300$ pb$^{-1}$.
3 Conclusions

The commissioning of the CMS detector is in its final phase. More and more subsystems are joining global data taking exercises that use muon signals from cosmic rays for coarse timing, calibration and alignment. Parallel to the commissioning of the detectors, the two multipurpose experiments ATLAS and CMS are preparing the commissioning of the basic physics objects, such as jets, photons and missing transverse energy, with first data. Various methods have been studied with more realistic simulations of the initial calibration and alignment and techniques are developed to establish the sound understanding of the detector response and calibration flows from the early data itself.

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SEARCHES FOR NEW PHYSICS IN PHOTON AND JET FINAL STATES

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Recent results from searches of physics beyond the standard model in pp collisions are reported, in particular, reactions involving high transverse momentum photons or jets in their final state. Data analyzed by the CDF and D0 experiments at the Run II of the Tevatron correspond to integrated luminosities between 1 and 2 fb$^{-1}$ depending on the analyses.

1 Introduction

At an energy frontier collider, the usual way to search for indices of physics beyond the standard model (SM) is to look for the collisions with the highest momentum-transfer particles. Typically, one chooses a particular model, and the event selection is optimized to enhance its contribution against the SM expectation. The absence of any deviation in data provides a limit on the production cross-section times the branching ratio for the channel under study, which is then translated into exclusion limits in the parameter space of this model. However, by nature, a new phenomena is unknown, and it exists a lot of models at disposal. This is the reason which motivates the "signature-based" search strategy which casts a wider look for deviations to the SM.

Both strategies will be reported here for final states with photons and jets.

2 Randall-Sundrum (RS) graviton

Many models with extra spatial dimensions have been proposed to solve the hierarchy problem. In the RS model$^1$, the SM brane and the Planck brane are separated by an extra dimension with a warped geometry. Only the graviton is allowed to propagate in this extra dimension. It appears as Kaluza-Klein (KK) towers in the SM brane. This model has only 2 parameters: $M_1$, the mass of the lowest KK excited mode, and $k/M_{Pl}$, a dimensionless coupling constant whose value should lie between 0.01 and 0.1.

KK towers couple to any boson or fermion pairs. CDF$^2$ looks separately at $\gamma\gamma$ and $ee$ final states, whereas D0$^3$ looks at both final states at once as they look similar in the electromagnetic calorimeter. Because of the spin 2 of the graviton, the ratio of the branching ratios to $\gamma\gamma$ and $ee$ final states is 2. Both experiments have analyzed about the same amount of data ($\sim 1 fb^{-1}$), and found no excess of events over the SM predictions (Drell-Yan and QCD where jets are misidentified as photons) excluding graviton masses below 900 GeV/c$^2$ for $k/M_{Pl} = 0.1$. Fig. 1 shows the excluded contour in the 2D parameter space as measured by D0.
3 Gauge mediated SUSY breaking (GMSB)

SUSY\(^4\) is a broken symmetry. Experimental signatures are determined through the manner and scale of the SUSY breaking. In the GMSB scenario, the lightest supersymmetric particle (LSP) is the gravitino, a very light and weakly interacting particle. The next to lightest supersymmetric particle (NLSP) is assumed in this analysis to be the neutralino which decays into the LSP and a photon. Assuming R-parity conservation\(^5\), SUSY particles are pair produced and the experimental signature will be 2 photons and missing energy from the 2 gravitinos. To get a quantitative result, the "Snowmass Slope SPS 8\(^{a}\) model\(^6\) is considered. All the GMSB parameters\(^a\) are fixed as a function of the effective energy scale \(\Lambda\) of SUSY breaking.

In this event topology, the SM background is the \(Z\gamma\gamma\) production where the \(Z\) boson decays into neutrinos. There is also important instrumental background from events with real \(E_T\) (W boson production) and fake \(E_T\) (QCD where jets are misidentified as photons). Fig. 2 (left) shows the \(E_T\) distribution. The observed distribution agrees well with the SM prediction; the entire spectrum is then used to set limits on the GMSB production cross section. Fig. 2 (right) shows the 95% C.L. cross section limit as a function of the effective scale \(\Lambda\) obtained by D0\(^7\). The observed limit on the signal cross section is below the prediction of the Snowmass Slope model for \(\Lambda < 61.5\) TeV, or for gaugino masses \(m_{\tilde{\chi}^0_i} < 125\) GeV/c\(^2\) and \(m_{\tilde{\chi}^\pm_i} < 229\) GeV/c\(^2\).

4 Large Extra Dimensions

The hierarchy problem can also be solved by postulating the existence of \(n\) new large extra dimensions as proposed first by Arkani-Hamed, Dimopoulos and Dvali\(^8\) (ADD); the extra volume serves to dilute gravity so that it appears weak in our 3D world as the graviton is the only particle allowed to propagate in the extra space. If the extra dimensions are compactified in a torus of radius \(R\), according to the Gauss law, one can relate the fundamental Planck mass scale \(M_D\), \(R\), the Planck mass and the number of extra dimensions by the relation \(M_{\text{Planck}}^2 = 8\pi M_D^{n+2} R^n\), allowing \(M_D\) to be compatible with the electroweak scale.

In this model, the graviton can be produced directly in the reaction \(q\bar{q} \rightarrow G\gamma\); G will remain undetected leaving a signature with a single photon and \(E_T\).

The only SM background is the \(Z\gamma\) production where the \(Z\) boson decays into a neutrino pair. In addition to the usual instrumental background coming from misidentification of electrons or jets into photons, the event topology is rather sensitive to a contribution from beam halos and

\(^a\)The messenger mass \(M_m = 2A\), the number of messengers \(N_\xi = 1\), \(\tan(\beta) = 15\), \(\mu > 0\).
cosmetics where muons produced photons by bremsstrahlung.

To fight the latter background, both experiments had to develop specific tools in addition to the usual ones based on the EM shower profile. Special hit finders in the tracker starting from the EM cluster increase the track veto efficiency. In addition, D0 uses a EM pointing tool thanks to its preshower detector, and CDF the timing system built within its EM calorimeter.

The results of CDF which has analysed about 2 fb⁻¹ of data, twice as much as D0⁹, are displayed in Fig. 3. The left plot shows a good agreement for the photon transverse between data and the sum of the various backgrounds. This allows to set limits on the fundamental scale \( M_L \) (right plot) as a function of the number of extra dimensions. For \( n > 4 \) the limits are comparable with the limits obtained in the monojet search, and better than the LEP combined result¹⁰.

![CDF Run II Preliminary](image)

**Figure 3:** The left figure shows the \( E_T \) distribution in the CDF monophoton search. The signal expected from the ADD model (\( n=4, m=0.8 \text{ TeV} \)) is added on top of the SM backgrounds. The right figure shows the exclusion limits for the ADD model obtained in this analysis, in comparison with the CDF jet and \( E_T \) result and LEP combined result.

5 Squarks and gluinos, stops

5.1 Squarks and gluinos

Squarks and gluinos can be copiously produced at the Tevatron if they are sufficiently light. The analysis is performed within the mSUGRA model¹¹. The final state is composed of jets with a large \( E_T \) due to the two escaping neutralinos, assumed to be the LSP. According to the relative mass of squarks and gluinos different event topologies are to be expected. If squarks are lighter than gluinos, a "dijet" topology is favored. On the contrary if squarks are heavier than gluinos, the final state contains at least 4 jets. Finally, the jet multiplicity is at least 3 if squarks and gluinos have similar masses. After a common event preselection, the three topologies have been studied and optimised separately. The left plot on Fig. 4 shows the D0 \( E_T \) distribution obtained in the "dijet" search, the right one is obtained by CDF in the "3-jet" search. D0 has analyzed 2.1 fb⁻¹ of data without finding any excess over the SM predictions. It allows to extend the exclusion domain in the squark gluino plane (Fig. 5). Using the most conservative hypothesis, D0¹²(CDF¹³) excludes a gluino lighter than 308(290) GeV/c².

5.2 Stop

Due to the large Yukawa coupling, there could be a large mixing in the 3rd generation of squarks. The lightest of the 2 stops could be the lightest squark and even the NLSP. Furthermore if its
mass is less than the sum of the masses of the b quark, the W boson, and the neutralino, the dominant decay mode is $t \rightarrow c\tilde{\chi}_1^0$, a flavor changing loop decay, assumed to be 100% in the analysis. The final state will then be 2 acoplanar charm jets and $E_T$. The analysis proceeds with 2 jets detected in the central part of the detector with a loose heavy quark tag for one of them. No excess of events has been observed\cite{114} in about 1 fb$^{-1}$ of data, which provides a lower limit for the stop mass at 149 GeV/c$^2$ for a neutralino mass of 63 GeV/c$^2$ (Fig. 6).

![Figure 4: Distributions of $E_T$ after applying all analysis criteria except the one on $E_T$ for the “dijet” (DC left) and “3-jets” (CDF right) squark-gluino analyses. Data (points with error bars) and the cumulative contributions from SM background, QCD background and signal MC are shown.](image)

![Figure 5: Excluded region in the squark and gluino mass plane; newly excluded domain by D0 is shown in dark shading. The region where no mSUGRA solution can be found is shown hatched.](image)

![Figure 6: Region in the $t$-$\tilde{\chi}_1^0$ mass plane excluded at the 55% C.L. by the DC search. The yellow band represents the theoretical uncertainty on the scalar top quark pair production cross section due to PDF and renormalization and factorization scale choice.](image)
6 Signature-based searches

6.1 Search for anomalous production of di photon events

In its quest for “signature-based” excess, CDF has searched for anomalous production of events in the $\gamma \gamma + E_T$ topology. In this analysis, use is made of the $E_T$ resolution model. The aim of this model is to discriminate events with large mismeasured $E_T$ from events with real $E_T$. It has been shown to provide a better background rejection power than a simple $E_T$ cut. This model is based on the assumption that individual particle’s energy resolution has Gaussian shape proportional to particle’s $\sqrt{E_T}$. Only two sources of fake $E_T$ are considered: soft unclustered energy (from underlying event and multiple interactions), and jets. The latter is responsible for most of the $E_T$, as it is collimated energy in contrast to the former which is spread out all over the calorimeter. According to this model, each event is given a $E_T$ significance value. Most of the QCD background is eliminated by requiring a significance above 5, leaving only the expected number of SM events with real $E_T$ (Fig. 7), and not much room for an extra signal.

Figure 7: Distribution of missing transverse energy significance for diphoton candidates.

Figure 8: The measured dijet mass spectrum and results of the fit to the parametrization form 1.

6.2 Search for dijet mass resonances

Many classes of models beyond the SM predict the existence of new massive particles decaying into 2 partons which would appear as resonances in the dijet mass spectrum. Such classes include excited quarks, techniparticles, new W’ or Z’ bosons, RS gravitons,... Jets are reconstructed by the cone-based midpoint jet algorithm with a cone radius of 0.7 and have central rapidity ($|y| < 1$). CDF has analysed about 1.1 $fb^{-1}$ of data and measured the dijet differential cross section (Fig. 8). The spectrum is fitted by the smooth parametrization:

$$\frac{d\sigma}{dm} = p_0 (1 - x)^{p_1} / x^{p_2} + p_3 \delta(x), \quad x = m/\sqrt{s}.$$  

This parametrization is found to fit well the dijet spectra from PYTHIA and HERWIG MC events as well as from NLO pQCD. As no evidence for existence of a new massive particle is observed, limits on new particle production cross sections can be derived as a function of the dijet mass. These limits are then translated into mass exclusion limits, see Table 1.

7 Conclusions

No hints of physics beyond the SM have been found so far. As the Tevatron is continuing to provide experiments with more data to analyze, the quest for indices will be pursued by CDF
Table 1: Mass exclusion ranges for several models.

<table>
<thead>
<tr>
<th>Model description</th>
<th>Observed mass exclusion range (GeV/c^2)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Excited quark (f = f' = f_0)</td>
<td>200-870</td>
</tr>
<tr>
<td>Color octet technicolor (top-color-assisted technicolor couplings)</td>
<td>263-1110</td>
</tr>
<tr>
<td>Axigluon and flavor universal technicolor (mixing of 2 SU(3)'s cot(\theta)=1)</td>
<td>263-1250</td>
</tr>
<tr>
<td>E6 diquark</td>
<td>200-630</td>
</tr>
<tr>
<td>W' (SM couplings)</td>
<td>280-840</td>
</tr>
<tr>
<td>Z' (SM couplings)</td>
<td>350-740</td>
</tr>
</tbody>
</table>

and D0. Some analyses presented in this talk have already been published, for the others, further details can be found at:

CDF  http://www-cdf.fnal.gov/physics/exotic/exotic.html
D0  http://www-d0.fnal.gov/Run2Physics/WWW/results/mp.htm

Acknowledgments

The author would like to thank the CDF and D0 working groups for providing the material for this talk, and the organizers of the Rencontres for a very enjoyable conference and the excellent organization.

References

1 Commissioning of ATLAS

At the time of this paper, everything except the endwall muon chambers has been installed. In this section we summarize the acquired understanding of the detector performance before start-up.

1.1 Detector status at start-up

The endwall muon chambers are the last components to be installed in the cavern (June 2008). With the full magnet tests performed soon afterwards in order to check the endcap toroids at their nominal field, the entirety of the ATLAS detector will have been assembled ready for the first collisions. Data taking has already been tested with cosmic rays throughout the two preceding years of commissioning as the different detectors have been installed. The complete acquisition and processing chain has been successfully exercised from the primary site at CERN and throughout worldwide distributed sites.

1.2 Inner detectors

The commissioning of the Inner Detector was initiated on the surface where cabling checks and debugging were performed. Once in the cavern, the noise levels were the same as those measured on the surface. The global alignment between the silicon strip (SCT) and transition radiation (TRT) trackers was measured with reconstructed cosmic ray tracks (as shown in Fig.1), where the results were consistent with precision survey measurements. Little time was available for cabling cross-checks of the innermost layer of pixels which was cabled just before ATLAS was closed. The number of non-working channels was measured to be small (0.3%).
1.3 Calorimeters

The calorimeters went down into the cavern beginning in 2006 and have been operating under stable conditions for a long period of time. For example, the liquid argon calorimeter was cold (88 K) with a measured temperature variation of 10 mK for over a year. The degree of understanding of the detector response is well illustrated by cosmic rays. Large signals as in Fig. 1 are frequently observed, interpreted as muons emitting a bremsstrahlung photon in the calorimeters. These events have allowed timing studies (channel intercalibration at the level of 2 ns) as well as improving the description of the pulse shape. In general, muons are minimum ionising particles and deposit energy according to a Landau distribution. By analyzing these distributions, the uniformity of response across EM readout channels could be cross-checked to the 2% level\(^1\). Taking advantage of the several months of commissioning running, many problems were identified and corrected, leaving only a small fraction of defective channels (~0.1% in the case of the LAr calorimeter).

![Figure 1:](image)

Figure 1: On the left is shown the trajectory of a cosmic muon reconstructed in the inner detectors. No zero suppression being performed, this display reveals that only a fraction of the TRT was recording data. The right plot shows a large amplitude signal (~13 GeV) in the LAr calorimeter which was also recorded during the commissioning. Compared to the predicted pulse, the residual in green is showing good agreement at the percent level.

1.4 Muon spectrometer

The commissioning of the muon spectrometer was aimed at testing its complex alignment system using optical sensors. The alignment is of crucial importance as a 1 TeV muon is bent only by 500 μm under the magnetic field produced by the toroids. In order to reach 10% accuracy on the momentum measurement, the chambers must be aligned with a precision of 50 μm. After a calibration of the system by moving and rotating modules, the reconstructed cosmic tracks have shown that the misalignments can be well monitored.

2 Early physics with leptons

With the first collisions, the main goal will be to calibrate the detectors \textit{in situ} using well-known physics samples. Table 1 summarizes the expected performances at start-up and those expected to be ultimately reached. The rediscovery of standard model at $\sqrt{s} = 10$ TeV will be followed by an effort to validate and tune the MC generators. Leptons will play a crucial
role in paving the way to new physics, taking half the trigger rate of the foreseen menu at $\mathcal{L} = 10^{31}\text{cm}^{-2}\text{s}^{-1}$. Assuming 3 months of operation in 2008 and an integrated luminosity around 100pb$^{-1}$, calibration procedures and possible early measurements are described in the following.

<table>
<thead>
<tr>
<th>Performance</th>
<th>Physics</th>
</tr>
</thead>
<tbody>
<tr>
<td>@ start-up</td>
<td>nominal</td>
</tr>
<tr>
<td>Physics tools</td>
<td></td>
</tr>
</tbody>
</table>

### Table 1: This table summarizes the detector performance concerning leptons which are expected at start-up. The values, mostly given for the ATLAS case, are compared to those ultimately reached using the physics tools mentioned in the last column.

<table>
<thead>
<tr>
<th>Performance</th>
<th>Physics</th>
</tr>
</thead>
<tbody>
<tr>
<td>EM energy uniformity</td>
<td>$\lesssim$1-2%</td>
</tr>
<tr>
<td></td>
<td>$\lesssim$2-4%</td>
</tr>
<tr>
<td>CMS</td>
<td>$H \rightarrow \gamma\gamma$</td>
</tr>
<tr>
<td>EM energy scale</td>
<td>$\sim$2%</td>
</tr>
<tr>
<td>Inner detector alignment</td>
<td>50-100$\mu$m</td>
</tr>
<tr>
<td>Muon system alignment</td>
<td>$&lt;200\mu$m</td>
</tr>
<tr>
<td>Muon momentum scale</td>
<td>$\sim$1%</td>
</tr>
</tbody>
</table>

2.1 First peaks: $(J/\psi, \Upsilon, Z) \rightarrow \mu\mu$

Assuming a 30% overall detector and machine efficiency at $10^{31}\text{cm}^{-2}\text{s}^{-1}$, around 16000 $J/\psi$ and 3000$\Upsilon$ in dimuons are accumulated per pb$^{-1}$, useful to perform checks, tracker alignment and momentum scale determination. After all cuts, $\sim$600 $Z \rightarrow \mu\mu$ events are also recorded. As shown in figure 2, the measured $Z$ boson mass is sensitive to misalignments of the tracker or the muon system as well as uncertainties on the magnetic field (distored B field): a misaligned tracker would affect noticeably the shape of the $Z$ peak.

![Figure 2: A peak of $\sim$6000 $Z$ obtained with 10pb$^{-1}$ is the starting point of many studies of the muon system as shown on the left in the case of CMS (the same effects exist in ATLAS). A clear $W$ peak is obtained as well (right). A cut at 50GeV on the transverse mass ($M_T$) gives a clean $W$ sample that can be used for $E_T$ studies.](image)

2.2 $Z \rightarrow ee$ calibration and energy scale

Simple analysis cuts are used to obtain the first $Z$ and $W$ samples with low background levels as in figure 2 (right). For example, they can be done without the tracker or restricted to the barrel region if required. The $Z$-mass constraint is a key tool for the commissioning of electron
reconstruction and will also be used to correct residual long-range non-uniformities. From the design of ATLAS LAr calorimeter, relatively large regions (0.2 × 0.4 in η × φ) are expected to be locally uniform and this was demonstrated from testbeam data\textsuperscript{2}. In order to intercalibrate them, \(~30000\) \(Z \rightarrow ee\) events will be sufficient to achieve the desired response uniformity of \(~0.7\%\) (see Fig.3). The situation in CMS \textsuperscript{3} is different due to the crystal-to-crystal response which varies by \(\geq 2\%\). Higher statistics (\(~10\,pb^{-1}\)) are required to perform a similar intercalibration.

![Figure 3: The ATLAS EM intercalibration as a function of the integrated luminosity (left) reaches 0.7% at about 160\,pb\(^{-1}\). In the right plot, the efficiency of selecting electrons in ATLAS is accurately evaluated using the tag&probe method described in the text (1% agreement w.r.t. MC generation).](image)

### 2.3 \(Z\) and \(W\) cross-sections

The expected precision for ATLAS and CMS in the measurements of the \(Z\) and \(W\) cross-sections (table 2) are limited by statistical uncertainties. Theoretical uncertainties remain at 2\%, mainly due to acceptance determination and PDFs. The luminosity uncertainty of 10\% will dominate until the dedicated detectors \textsuperscript{4,5} refine these measurements in 2009. Systematics are kept to the 1\% level with efficiency precisely determined using data-driven methods (tag&probe) as illustrated in Fig.3 in the case of electron identification. In simulated \(Z\) events, an electron candidate is tagged by fulfilling stringent isolation and reconstruction requirements. A simple object is then looked for on the opposite side, either a track or an EM cluster. If the invariant mass of the pair falls in the mass window of the \(Z\) (\(M_{ee}=M_\tau\pm 20\,\text{GeV}\)), the latter is known to be an electron and the efficiency can be computed in this way.

<table>
<thead>
<tr>
<th>process</th>
<th>(\Delta z_a)</th>
<th>ATLAS</th>
<th>CMS</th>
</tr>
</thead>
<tbody>
<tr>
<td>(Z/\gamma^*\rightarrow \mu\mu)</td>
<td>0.004(stat)+0.006(sys) \pm 0.02(syst) \pm 0.1(lumi)</td>
<td>0.007(stat)+0.009(sys) \pm 0.02(syst) \pm 0.01(lumi)</td>
<td></td>
</tr>
<tr>
<td>(W \rightarrow e\nu)</td>
<td>0.002(stat)+0.030(sys)</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

Table 2: The various contributions to the uncertainty of cross-section measurements are given for 50\,pb\(^{-1}\) of integrated luminosity. Higher systematic uncertainties are expected in the case of \(\sigma_W\), especially from the \(p_T\) measurement.
Figure 4: The electron rapidity distributions in $W$ decays can be used to constrain PDF distributions (left). Dielectron resonances at the TeV scale could be seen with a fraction of the statistics shown here for $1\ fb^{-1}$ (right).

### 2.4 Constrain PDFs with $W \rightarrow \ell\nu$

A region of the $x - Q^2$ plane becomes accessible at the LHC with a higher centre-of-mass energy ($x_{1,2} = \frac{M_{\ell\ell}}{\sqrt{s}}$). For instance, the $W$ production would involve $10^{-4} < x_{1,2} < 0.1$, a region dominated by sea-sea parton interactions. At these low $x$ values, the current PDF uncertainties remain large (4-8%\(^6\)). By adding LHC data into global fits, it would be possible to constrain further the PDFs. With a 5\% experimental precision, $e^{\pm}$ angular distributions could be used to discriminate between different PDF parametrizations. The principle is demonstrated by generating $10^6\ W \rightarrow \ell\nu$, equivalent to $150\ pb^{-1}$ of data. These simulated events are generated with the CTEQ6.1 PDF and full detector simulation. Statistical uncertainty being negligible, only a 4\% systematic error is introduced by hand, a level already attained with the $Z \rightarrow ee$ sample. These pseudo-data are included in the global ZEUS PDF fit in order to assess their impact. The uncertainty on the low-$x$ gluon shape parameter $\lambda \ [xg(x) \sim x^{-\lambda}]$ is reduced by 40\%.

### 2.5 Leptons for early top

The easiest channel to consider at the beginning is one lepton plus jets ($t\bar{t} \rightarrow b\ell b jj$) with $t\bar{t}$ combinatorics, $W$-jets and QCD as main backgrounds. The b-tagging will not be available at start-up and $p_T$ might be problematic as well. For these reasons, early top analyses are lepton-triggered. Nevertheless, a signal can be quickly seen with only $\sim 10\ pb^{-1}$, even with limited detector performance and a simple analysis. With $100\ pb^{-1}$, these early measurements of $\Delta r_{t\bar{t}}$ will reach $\sim 20\%$ and $\Delta m_{t\bar{t}}$ at $<10\ GeV$. It will be an excellent sample for light jet calibration and b-jet efficiency determination.

### 2.6 Early discoveries with leptons?

The discovery of a narrow resonance decaying to $e^+e^-$ (Fig.4) is accessible with as few as $70\ pb^{-1}$ if $M \sim 1\ TeV$. At these energies, ultimate calorimeter performance is not needed. Early searches in the case of dimuons decays would be less straightforward since the resolution worsens as $p_T$ increases. However, the generally lower instrumental background may make dimuons a discovery channel along with dielectrons.
3 Conclusions

ATLAS has been commissioning its detector with cosmic rays and is eagerly awaiting the first LHC collisions in 2008. The clean signatures involving leptons will have a major impact in the initial understanding of the detectors. Summarized in figure 5, the early physics program involves several activities with possible surprises along the way.

Figure 5: Summary of the physics program with the first year of LHC data taking.

References

Electroweak Interactions and Unified Theories

I I- Particle Searches, Supersymmetry, Theoretical issues
Collider searches for new physics and direct searches for dark matter are important topics at this conference. In this contribution, I discuss their potential synergy and complementarity by means of the minimal supersymmetric standard model with a neutralino dark matter candidate.

1 Introduction

Cosmological data ranging from the cosmic microwave background to rotation curves of spiral galaxies tell us that most of the mass in the Universe is provided by non-luminous, hence ‘dark’ matter\textsuperscript{1,2,3}. More precisely, the recent measurements from WMAP\textsuperscript{4} and SDSS\textsuperscript{5} imply a (dominantly cold) dark matter density of $\Omega h^2 \simeq 0.1$ to an accuracy of about 10%.\textsuperscript{a} The nature of this dark matter is one of the big open questions of present-day physics. Many lines of reasoning suggest, however, that it consists of a new weakly interacting massive particle, a so-called WIMP.

At the same time, we know that the Standard Model (SM) of particle physics, despite its tremendous success at energies up to $\sim 100$ GeV, is incomplete. In attempts to embed the SM in a more fundamental frame, theorists have come up with a wealth of Beyond the Standard Model (BSM) theories, which typically predict new particles and phenomena at the TeV energy scale. To probe this exciting new frontier is indeed the primary motivation to build the LHC!

It is even more exciting that the lightest of these new BSM particles is often stable by virtue of a new discrete symmetry (introduced to match electroweak precision measurements and/or the non-observation of proton decay) and hence provides a natural dark matter candidate.

The dark matter candidates such put forth by particle physics are quite numerous\textsuperscript{3} and contain, for example, the lightest supersymmetric particle in supersymmetry with R-parity con-

\textsuperscript{a}The exact mean value and error depend on the data combination and number of parameters fitted, see Ref.\textsuperscript{6}.
Table 1: Standard Model particles and their superpartners in the MSSM. 

supersymmetry; the lightest Kaluza–Klein (KK) excitation in models with extra dimensions and KK-parity; the lightest T-odd state in little Higgs models with T-parity; etc. Note that all these possibilities are generally testable in collider experiments. This creates a strong interplay between particle physics, astrophysics and cosmology, at both theoretical and experimental levels.

2 Supersymmetry

Of the existing BSM theories, supersymmetry (SUSY) is arguably the best motivated one. SUSY is a symmetry between fermions and bosons. A SUSY generator \( Q \) changes a fermion into a boson and vice versa:

\[
Q |\text{fermion} \rangle = |\text{boson} \rangle, \quad Q |\text{boson} \rangle = |\text{fermion} \rangle.
\] (1)

This is an extension of space-time to include anti-commuting coordinates \( x^\mu \rightarrow (x^\mu, \theta^\alpha) \) with \( \{ \theta^\alpha, \theta^\beta \} = \varepsilon^{\alpha\beta} \), combining the relativistic ‘external’ symmetries (such as Lorentz invariance) with the ‘internal’ symmetries of a field, such as weak isospin. It is in fact the unique(!) extension of the Poincaré algebra (the algebra of space-time translations, rotations and boosts).

From the phenomenological point of view, SUSY predicts a partner particle, a so-called ‘superpartner’ or ‘sparticle’, for every SM state. The particle content of the Minimal Supersymmetric Standard Model (MSSM), is given in Table 2. In its local gauge theory version, SUSY also includes spin-2 and spin-3/2 states, the graviton and its superpartner the gravitino and is hence potentially capable of connecting gravity with the other interactions (so-called supergravity or short SUGRA). A few more things are important to observe:

i) SUSY must be a broken symmetry, else SM particles and their superpartners would have equal mass. In order to still solve the hierarchy problem of the SM (i.e. to stabilize the electroweak scale against quadratically divergent radiative corrections) and to achieve gauge-coupling unification, one expects the superpartners to have masses of \( m \leq O(1) \text{ TeV} \).

ii) After electroweak symmetry breaking, we are left with three neutral Higgs bosons: two scalars \( h, H \) and one pseudoscalar \( A \). Moreover, particles with the same \( SU(3) \times U(1) \) quantum numbers mix, c.f. Table 2. In particular, the bino, wino and neutral higgsinos mix to mass eigenstates called neutralinos \( \tilde{\chi}_1, \ldots, 4 \) (with \( \tilde{\chi}_1 \) the lightest one by definition). 

\begin{table}
\begin{tabular}{|c|c|c|c|c|c|c|c|c|}
\hline
\textbf{Standard Model particles and fields} & \textbf{Supersymmetric partners} & \textbf{Interaction eigenstates} & \textbf{Mass eigenstates} \\
\hline
\textbf{Symbol} & \textbf{Name} & \textbf{Symbol} & \textbf{Name} & \textbf{Symbol} & \textbf{Name} \\
\hline
\( q = d, c, b, u, s, t \) & quark & \( \tilde{q}_L, \tilde{q}_R \) & squark & \( \tilde{q}_1, \tilde{q}_2 \) & squark \\
\( l = e, \mu, \tau \) & lepton & \( \tilde{l}_L, \tilde{l}_R \) & slepton & \( \tilde{l}_1, \tilde{l}_2 \) & slepton \\
\( \nu = \nu_e, \nu_\mu, \nu_\tau \) & neutrino & \( \tilde{\nu} \) & sneutrino & \( \tilde{\nu} \) & sneutrino \\
\( g \) & gluon & \( \tilde{g} \) & gluino & \( \tilde{g} \) & gluino \\
\( W^\pm \) & W-boson & \( W^\pm \) & wino & \( \tilde{\chi}_1, \tilde{\chi}_2 \) & chargino \\
\( H^- \) & Higgs boson & \( H^- \) & higgsino & \( \tilde{\chi}_1, \tilde{\chi}_2 \) & neutralino \\
\( H^+ \) & Higgs boson & \( H^+ \) & higgsino \\
\( B \) & B-field & \( B \) & bino \\
\( W^3 \) & W³-field & \( W^3 \) & wino \\
\( H^0_1 \) & Higgs boson & \( H^0_1 \) & higgsino \\
\( H^0_2 \) & Higgs boson & \( H^0_2 \) & higgsino \\
\( H^0_3 \) & Higgs boson & \( H^0_3 \) & higgsino \\
\hline
\end{tabular}
\end{table}
iii) If SUSY comes with a new conserved parity, so-called R-parity, under which SM particles are even and SUSY particles are odd, the lightest supersymmetric particle (LSP) is stable. In this case it has to be electrically and colour neutral and constitutes a natural dark matter candidate.

In the following I concentrate on the MSSM with a neutralino LSP. For gravitino dark matter, which has a quite different phenomenology, I refer to the contribution by F. Steffen in these proceedings. Sneutrinos in extensions of the MSSM are discussed by C. Arina in the YSF.

3 Collider searches

If low-scale supersymmetry is realized in Nature, experiments at the LHC have excellent prospects to discover it\(^\text{10,11,12}\). In particular, squarks and gluinos should be copiously produced at the LHC through the QCD interaction, with cross sections of \(\mathcal{O}(1)\) pb for masses around 1 TeV.

This is followed by (multi-step) decays into lighter sparticles. Squarks decay into gluinos plus jets, \(\tilde{q} \rightarrow q \tilde{g}\), if kinematically allowed, or into charginos/neutralinos plus jets, \(\tilde{q}_L \rightarrow q \tilde{\chi}_i^\pm\), \(q \tilde{\chi}_j^0\) and \(\tilde{q}_R \rightarrow q \tilde{\chi}_j^0\) (\(i = 1, 2; j = 1, ..., 4\)). Gluinos always decay into squarks, either in the two-body mode \(\tilde{g} \rightarrow q \bar{q}\) if kinematically open, or else \(\tilde{g} \rightarrow q \tilde{q} \tilde{\chi}_i^\pm\), \(q \tilde{\chi}_j^0\) via an off-shell squark. The charginos \(\tilde{\chi}_1^\pm\) and neutralinos \(\tilde{\chi}_2^0, 3^0, 4^0\) decay further, e.g., \(\tilde{\chi}_i^\pm \rightarrow W^\pm \tilde{\chi}_j^0\) or \(\tilde{\chi}_3^0 \rightarrow Z \tilde{\chi}_j^0\), until the LSP \(\tilde{\chi}_1^0\) is reached. The LSP, being stable and neutral, escapes undetected.

SUSY events are hence characterized by multiple hard jets, maybe accompanied by leptons, plus large missing transverse energy \(E_T^{\text{miss}}\). The significance of such a signal over the SM background is illustrated in the left plot in Fig. 1, which shows the number of events as a function of the ‘effective mass’ computed from the missing energy and the momenta of the hardest jets, \(M_{\text{eff}} = E_T^{\text{miss}} + \sum p_T^{\text{jets}}\). Note that the y-axis is log-scale! The \(M_{\text{eff}}\) distribution also provides a first estimate of the gluino/squark mass scale.

In certain scenarios and/or with high enough statistics, electroweak production of charginos and neutralinos, e.g., \(pp \rightarrow \tilde{\chi}_1^\pm \tilde{\chi}_1^\mp\), \(\tilde{\chi}_1^\pm \tilde{\chi}_2^0\), can also be important. The latter process can lead to the goldplated tri-lepton signal, which is also searched for at the Tevatron\(^\text{13}\). Moreover, for slepton masses up to 300 GeV, slepton-pair production can lead to detectable di-lepton signals.

The discovery of SUSY particles will be followed by detailed measurements of their masses and decay properties. Since the LSP escapes as missing energy, no mass peaks can be reconstructed. Instead, mass measurements exploit kinematic distributions in cascade decays\(^\text{14,15}\). For

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Figure 1: Left: \(M_{\text{eff}}\) distribution for a SUGRA point with gluino and squark masses of about 700 GeV (histogram) and Standard Model background (shaded) after cuts. Right: Di-lepton invariant-mass distribution from \(\tilde{\chi}_2^0 \rightarrow l^\pm l^\mp \rightarrow l^+ l^- \tilde{\chi}_1^0\) decays. From\(^\text{19}\).
The cosmological evolution of a thermal relic’s comoving number density, from $\frac{1}{2}$. The full line is the equilibrium abundance; the dashed lines are the actual abundance after freeze-out. As the annihilation cross section $\langle \sigma_A v \rangle$ is increased, the WIMP stays in equilibrium longer, leading to a smaller relic density.

instance, the invariant-mass distribution of the leptons stemming from the chain $\tilde{\chi}^0_2 \rightarrow l^+ l^- \tilde{\chi}^0_1$ has a triangular shape with a sharp endpoint at $M_{\tilde{l}\tilde{l}}^{\text{max}} = \left[ (m_{\tilde{\chi}}^2 - m_{\tilde{l}}^2)(m_{\tilde{l}}^2 - m_{\tilde{\chi}}^2) / m_{\tilde{l}}^2 \right]^{1/2}$, which can be measured very precisely, see the right plot in Fig. 1. If the leptons come from the three-body decay $\tilde{\chi}^0_2 \rightarrow l^+ l^- \tilde{\chi}^0_1$, $M_{\tilde{l}\tilde{l}}$ has a different shape and an endpoint at $m_{\chi} - m_{\chi}$.

Additional distributions can be constructed involving jets stemming from gluino and squark decays. This way the masses of the sparticles appearing in the decay chains can be reconstructed.

Let us finally come back to the dark matter question. The alert reader will have noticed that because of R-parity sparticles are produced in even numbers, and every sparticle decay terminates in the LSP. As a consequence, each SUSY event contains two LSPs. Moreover, if squarks and gluinos weigh about 1 TeV, we expect of the order of 100 events/day —at low luminosity. The LHC may hence well turn out as a dark matter factory, where the nature and properties of dark matter candidates may be studied in a controlled environment.

Typical precisions at the LHC are $O(10\%)$. Much higher precisions at the percent to permil level might be achieved at an International $e^+e^-$ Linear Collider (ILC). The determination of neutralino dark matter properties at LHC and ILC has been analyzed, e.g., in $^{17}$.

4 Relic density

The standard cosmological scenario assumes that the dark matter particle, let us call it $\chi$, is a thermal relic of the Big Bang as illustrated in Fig. 2: When the early Universe was dense and hot, $T \gg m_\chi$, $\chi$ was in thermal equilibrium; annihilation of $\chi$ and $\tilde{\chi}$ into lighter particles, $\chi\tilde{\chi} \rightarrow l\bar{l}$, and the inverse process $l\bar{l} \rightarrow \chi\tilde{\chi}$ proceeded with equal rates. As the Universe expanded and cooled to a temperature $T < m_\chi$, the number density of $\chi$ dropped exponentially, $n_\chi \sim e^{-m_\chi/T}$.

Eventually the temperature became too low for the annihilation to keep up with the expansion rate and $\chi$ ‘froze out’ with the cosmological abundance observed today.

The time evolution of the number density $n_\chi(t)$ is described by the Boltzmann equation,

$$dn_\chi/dt + 3Hn_\chi = -\langle \sigma_A v \rangle \left[ (n_\chi)^2 - (n_\chi^{\text{eq}})^2 \right], \quad (2)$$

where $H$ is the Hubble expansion rate, $n_\chi^{\text{eq}}$ is the equilibrium number density, and $\langle \sigma_A v \rangle$ is the thermally averaged cross section times the relative velocity of the annihilating particles. The relic density today turns out to be inversely proportional to the annihilation cross section,
\[ \Omega \chi h^2 \propto 1/(\sigma_A v) \]. Note that \( \langle \sigma_A v \rangle \) includes a sum over all possible annihilation channels for the LSP. These are annihilation into gauge boson pairs through \( t \)-channel chargino and neutralino exchange, and annihilation into fermion pairs through \( t \)-channel sfermion exchange and \( s \)-channel \( Z/\text{Higgs} \) exchange. Moreover, co-annihilation channels involving sparticles that are close in mass to the LSP have to be taken into account. For details of the calculation, see \textsuperscript{2,18}.

The relic density of the LSP hence depends on all the MSSM masses and couplings that enter the different annihilation/co-annihilation channels. On the one hand, this is often used to severely constrain SUSY models by demanding that the relic density of the LSP falls within the WMAP–SDSS range (see the YSF contribution by S. Sekmen for an example). On the other hand, if the masses and couplings of SUSY particles are measured precisely enough, \( \Omega \chi h^2 \) can be computed and compared to the cosmologically observed value.

Here notice that the standard picture heavily relies on two assumptions: i) that the initial temperature after inflation has been high enough to fully thermalize the LSP and ii) that the entropy per comoving volume has been constant below the freeze-out temperature. In non-standard scenarios with low reheat temperature and/or late entropy production, the relic density can be quite different from the value in the standard scenario. A precise determination of the LSP annihilation cross section from collider experiments, together with a confirmation that the LSP is indeed the cold dark matter through direct detection (see next section), will hence allow to probe these assumptions \textsuperscript{19}, i.e. probe the evolution of the early universe up to the freeze-out temperature \( T_f \sim m_\chi/20 \).

5 Direct detection

Experiments such as CDMS\textsuperscript{20}, XENON\textsuperscript{21}, ZEPLIN\textsuperscript{22}, EDELWEISS\textsuperscript{23}, CRESST\textsuperscript{24}, KIMS\textsuperscript{25} and COUPP\textsuperscript{26} aim at detecting WIMPs through their elastic scattering with nuclei. The current experimental limits and projected sensitivities are shown in Fig. 3, together with predictions from various MSSM scenarios. Principally one distinguishes two classes, spin-dependent and spin-independent interactions. On the partonic level, WIMP interactions with quarks and gluons in the nucleons contribute.

In the case of neutralino dark matter, the scattering off quarks can occur through \( t \)-channel exchange of \( Z \) or CP-even Higgs bosons, or \( s \)-channel exchange of squarks:

\begin{align*}
\chi^0 & \quad \chi^0 \\
q & \quad q \\
Z, H, h & \quad q \quad q \\
\chi^0 & \quad \tilde{q} \\
\chi^0 & \quad \chi^0
\end{align*}

The diagrams with \( Z \) and squark exchange contribute to the axial-vector (spin) interaction, \( \mathcal{L} \sim \bar{\chi} \gamma_5 \gamma^\mu \chi q \bar{q} \gamma^\mu q \). The Higgs and squark exchange diagrams contribute to the scalar (spin-independent) interaction, \( \mathcal{L} \sim \bar{\chi} \chi \bar{q} q \). The neutralino interaction with gluons proceeds through quark and squark loops and contributes to the spin-independent cross section. See \textsuperscript{2,27} for details. Note that since the neutralino is a Majorana particle, there is no vector interaction of the form \( \mathcal{L} \sim \bar{\chi} \gamma_\mu \chi q \bar{q} \gamma^\mu q \).

The effective neutralino–nucleon coupling hence depends on the neutralino mass and decomposition (i.e. the bino/wino/higgsino content) as well as on the Higgs and squark masses and couplings. Again, if the supersymmetric spectrum is known from collider experiments, the scattering cross section can be predicted.\textsuperscript{c} A word of caution is, however, in order here because the strange content of the nucleon is not known well; this induces a considerable uncertainty\textsuperscript{29} in the neutralino–nucleon cross section, in particular if Higgs exchange dominates. Finally, the

\textsuperscript{c}This cross section also determines the rate at which neutralinos would accrete in the Earth and Sun.
neutralino–nucleon cross sections $\sigma_{\chi n}$ and $\sigma_{\chi p}$ have to be translated to the neutralino–nucleus scattering cross section $\sigma_{\chi N}$ applying nuclear form factors.

The actual direct-detection rate depends, moreover, on the local dark matter density $\rho_\chi$ and the velocity distribution $f(v)$. Roughly speaking, the rate of events per day and per kg of detector material is $R \sim \rho_\chi \sigma_{\chi N} \langle v \rangle / (m_\chi m_N)$, with $m_N$ the target nucleus mass and $\langle v \rangle$ the average velocity of $\chi$ relative to the target. Typical estimates are $\rho_\chi = 0.22 - 0.73$ GeV/cm$^3$ and $\langle v \rangle = 230 \pm 20$ km/s. If $m_\chi$ and $\sigma_{\chi N}$ are known with good precision, the local density and velocity distribution can be tested.

At this point, let me stress the importance of direct detection for another reason: Although collider experiments may identify a dark matter candidate and precisely measure its properties, they will not be able to distinguish a cosmologically stable from a very long-lived but unstable particle. Therefore validation of the collider signal through direct detection is essential. The key to this is the WIMP mass, which may be determined in direct-detection experiments through the distribution of the recoil energy, $E_R \propto 2v^2 m_N/(1 + m_N/m_\chi)^2$. Note that $E_R$ is sensitive to small $m_\chi$ but becomes almost constant for $m_\chi \gg m_N$. Note also the velocity dependence, which is source of considerable uncertainty in a single experiment. This may be evaded by using multiple targets. Precisions are, however, still poor for $m_\chi \gg m_N$.

6 Conclusions

For conclusions, let me cite G. F. Giudice in “Theories for the Fermi scale”:

_It is impossible to overestimate the importance of discovering dark matter at the LHC. Such a discovery will imply a revision of the SM, it will strengthen the connection between particle physics, cosmology and astrophysics, and it will enormously enlarge our understanding of the present and past universe._

So be prepared for exciting times at future Moriond meetings.

Acknowledgments

I wish to thank the organisers for creating a remarkably pleasant and inspiring atmosphere. Fruitful discussions with other participants are also gratefully acknowledged.
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MODEL INDEPENDENT SEARCHES IN EP COLLISIONS

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The high energy program of the HERA collider ended in March 2007. In total the H1 and ZEUS experiments collected an integrated luminosity of about 1 fb$^{-1}$. Recent results of model independent searches for new physics from both experiments are presented. Specifically, studies of the events with an isolated lepton and missing transverse momentum and multi-lepton topologies, where H1 and ZEUS data are combined, and a general signature based search are discussed.

1 Introduction

At HERA electrons (or positrons) collide with protons at a centre-of-mass energy of $\sqrt{s} \simeq 320$ GeV. During the two running periods of HERA from 1994 to 2000 and from 2003 to 2007, respectively, the H1 and ZEUS experiments have each recorded $\sim 0.5$ fb$^{-1}$ of data in total, shared between $e^+p$ and $e^-p$ collision modes. These high energy electron-proton interactions provide a testing ground for the Standard Model (SM) complementary to $e^+e^-$ and pp scattering studied at other colliders, giving access to rare processes with cross sections below 1 pb. They are therefore used to pursue a rich variety of searches for new phenomena. Among them, signature based searches look for differences in precise comparisons between data and SM expectations in different event topologies. As an advantage, such model independent analyses do not depend on any a priori definition of expected signatures for exotic phenomena. Following this approach, final states corresponding to rare SM processes such as real $W$ boson or lepton pair production are investigated. A general scan at high transverse momenta ($P_T$) of all possible final states is also performed by H1.

2 Events with high $P_T$ isolated leptons

The production of $W$ bosons in $ep$ collisions at HERA has a cross-section of about 1 pb. The leptonic decay of the $W$ leads to events with an isolated high transverse momentum lepton (electron, muon or tau) and missing total transverse momentum. Of particular interest are events with a hadronic system of large transverse momentum ($P_T^X$). An abnormally large rate of high $P_T^X$ events is observed by the H1 experiment$^{1,2}$ in the electron and muon channels. In the analysis of all HERA I and HERA II data sets, which amounts to a total luminosity of 478 pb$^{-1}$, 24 events are observed at $P_T^X > 25$ GeV for a SM expectation of $15.8 \pm 2.5$. Amongst them only 3 events are observed in $e^-p$ collisions, in agreement with the SM expectation of $6.9 \pm 1.0$, while 21 events are observed in the $e^+p$ data for an expectation of $8.9 \pm 1.5$ (see table 1).
The ZEUS experiment has carried out a similar analysis using 492 pb$^{-1}$ of 1996–2007 data$^3$. The results are also shown in table 1. At $P_T^X > 25$ GeV the number of data events observed by ZEUS is in agreement with the SM expectation in both $e^+p$ and $e^-p$ collisions. A detailed comparison between efficiencies of the H1 and ZEUS detectors for the $W$ signal was performed. Both efficiencies are comparable in the central region. While H1 detection region extends to lower polar angle than ZEUS, most of the high $P_T^X$ events observed by H1 are within the range of the ZEUS acceptance.

The data samples of the H1 and ZEUS experiments have been used for a combined analysis performed in a common phase space$^4$. The combined data set corresponds to a total integrated luminosity of 0.97 fb$^{-1}$. A total of 87 events containing an isolated electron or muon and missing transverse momentum are observed in the data, compared to a SM expectation of $92.7 \pm 11.2$. At $P_T^X > 25$ GeV, a total of 29 events are observed compared to a SM prediction of $25.3 \pm 3.2$. In this kinematic region, 23 events are observed in the $e^+p$ data compared to a SM prediction of $14.6 \pm 1.9$. The observations in the $e^+p$ and $e^-p$ data sets are exemplified in figure 1 where the $P_T^X$ distributions of both data sets are displayed.

![Figure 1: Hadronic transverse momentum distribution of isolated lepton events observed by H1 and ZEUS in $(a)$ $e^+p$ and $(b)$ $e^-p$ data samples. The total SM expectation is represented by the open histograms and the contribution from $W$ production by the hatched histogram.](image-url)

The analysis of the tau decay channel is also performed by H1$^5$ on all HERA data with a total luminosity of 471 pb$^{-1}$. In this channel, the separation of the $W$ signal from other SM processes is more difficult and the purity and efficiency are lower than for the $e$ and $\mu$ channels. In total, 20 data events are observed compared to a SM expectation of $19.5 \pm 3.2$. One of the data events has $P_T^X$ above 25 GeV, compared to a SM expectation of $0.99 \pm 0.13$. An older
analysis of the tau channel performed by the ZEUS Collaboration\textsuperscript{6} on HERA I data reported an observation of two data events with $P_T^X > 25$ GeV, compared to a SM expectation of $0.2 \pm 0.05$.

3 Multi-lepton events

The main production mechanism for multi-lepton events is photon-photon collisions. All event topologies with high transverse momentum electrons and muons have been investigated by the H1 experiment\textsuperscript{7} using a total luminosity of 459 pb\(^{-1}\). The measured yields of di-lepton and tri-lepton events are in good agreement with the SM prediction, except in the tail of the distribution of the scalar sum of transverse momenta of the leptons ($\sum P_T$). In $e^+p$ collisions, 4 data events with at least two high $P_T$ leptons are observed with $\sum P_T > 100$ GeV compared to a SM prediction of $1.2 \pm 0.2$. No such events are observed in $e^-p$ collisions for a similar SM expectation of $0.8 \pm 0.2$.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure2.png}
\caption{Distribution of the scalar sum of the transverse momenta of leptons compared to expectations separately for events recorded by H1 in $e^+p$ (a) and $e^-p$ (b) collisions and for all H1 data (c).}
\end{figure}

The analysis of di-electron ($2e$) and tri-electron ($3e$) topologies is also carried out by ZEUS using 478 pb\(^{-1}\) of data\textsuperscript{8}. Two data events with an electron pair invariant mass above 100 GeV are observed in each $2e$ and $3e$ channel. These observations are in good agreement with the corresponding SM expectations of $1.9 \pm 0.2$ and $1.0 \pm 0.1$ in the $2e$ and $3e$ channels, respectively.

Analyses of the $2e$ and $3e$ topologies from the H1 and ZEUS experiments have been combined in a common phase space\textsuperscript{9}. The total integrated luminosity amounts to 0.94 fb\(^{-1}\). The measured event yields of di-electron and tri-electron events are in good agreement with the SM predictions. The distribution of the invariant mass $M_{12}$ of the two highest $P_T$ electrons in $2e$ and $3e$ channels is presented in figure 3. In the $2e$ ($3e$) channel, 5 (4) events with an invariant mass $M_{12} > 100$ GeV are observed compared to a SM expectation of $3.4 \pm 0.4$ ($1.8 \pm 0.2$). Combining the two channels, six events are observed with $\sum P_T > 100$ GeV, compared to a SM expectation of $3.0 \pm 0.3$. Five of those events are observed in $e^+p$ collisions where the SM expectation is of $1.8 \pm 0.2$, whereas one event is observed in $e^-p$ data for a SM prediction of $1.2 \pm 0.1$.  


Figure 3: Distribution of the invariant mass $M_{12}$ of the two highest $P_T$ electrons of di-electron (a) and tri-electron (b) events observed by H1 and ZEUS in $e^+p$ data. The points correspond to the observed data events and the open histogram to the SM expectation. The total error on the SM expectation is given by the shaded band. The component of the SM expectation arising from lepton pair production is given by the hatched histogram.

4 A general search for new phenomena

A broad range signature based search has been developed by the H1 Collaboration on HERA I data \(^{10}\). All final states containing at least two objects ($e$, $\mu$, $j$, $\gamma$, $\nu$) with $P_T > 20$ GeV in the polar angle range $10^\circ < \theta < 140^\circ$ are now also investigated in all HERA II data \(^{11}\). The observed and predicted event yields in each channel are presented in figure 4(a) and (b) for $e^+p$ and $e^-p$ collisions, respectively. The good agreement observed between data and SM prediction demonstrates the good understanding of the detector and of the contributions of the SM backgrounds.

A systematic scan of the distributions of the scalar sum of transverse momenta $\sum P_T$ and of the invariant mass $M_{all}$ of all objects is performed in each channel to look for regions of largest deviations from the SM. In order to quantify the level of agreement between the data and the SM expectation and to identify regions of possible deviations, the search algorithm developed in reference \(^{10}\) is used. All possible regions in the histograms of $\sum P_T$ and $M_{all}$ distributions are considered. A statistical estimator $p$ is defined to judge which region is of most interest. This estimator is derived from the convolution of the Poisson probability density function (pdf) to account for statistical errors with a Gaussian pdf to include the effect of non negligible systematic uncertainties \(^{10}\). The value of $p$ gives an estimate of the probability of a fluctuation of the SM expectation upwards (downwards) to at least (at most) the observed number of data events in the region considered. The region of greatest deviation is the region having the smallest $p$-value, $p_{\text{min}}$. The fact that the deviation could have occurred at any point in the distribution is taken into account by calculating the probability $\hat{P}$ to observe a deviation with a $p$-value $p_{\text{min}}$ at any position in the distribution. This $\hat{P}$ is a measure of the statistical significance of the deviation observed in the data. The event class of most interest for a search is the one with the smallest $\hat{P}$ value.

The overall degree of agreement with the SM can further be quantified by taking into account the large number of event classes studied in this analysis. Among all studied classes there is some chance that small $\hat{P}$ values occur. This probability can be calculated with MC experiments. A MC experiment is defined as a set of hypothetical data histograms following the SM expectation with an integrated luminosity equal to the amount of data recorded. The complete search algorithm and statistical analysis are applied to the MC experiments analogously to the data. The expectation for the $\hat{P}$ values observed in the data is then given by the distribution of $\hat{P}^{SM}$ values obtained from all MC experiments.
The $\hat{P}$ values observed in the real data in all event classes are compared in figure 5 to the distribution of $\hat{P}^{SM}$ expected from MC experiments. The comparison is presented for the scans of the $\sum P_T$ distributions. Due to the uncertainties of the SM prediction in the $j$-$j$-$j$-$j$ and $j$-$j$-$j$-$j$-$\nu$ event classes at highest $M_{sh}$ and $\sum P_T$ (see reference\textsuperscript{10}), where data events are observed, no reliable $\hat{P}$ values can be calculated for these classes. These event classes are not considered to search for deviations from the SM in this extreme kinematic domain. All $\hat{P}$ values range from 0.01 to 0.99, corresponding to event classes where no significant discrepancy between data and the SM expectation is observed. These results are in agreement with the expectation from MC experiments. The most significant deviation from SM predictions is observed in the $\mu$-$j$-$\nu$ event class in $e^+p$ collisions with a value of $-\log_{10}\hat{P}$ equal to 1.7. In the previous H1 analysis\textsuperscript{10} based on HERA I data, which is dominated by $e^+p$ collisions, the largest deviation was also found in this event class, with $-\log_{10}\hat{P} = 3$.

Figure 4: The data and the SM expectation in event classes investigated by the H1 general search. Only channels with observed data events or a SM expectation greater than one event are displayed. The results are presented separately for $e^+p$ (a) and $e^-p$ (b) collision modes.

Figure 5: The $-\log_{10}\hat{P}$ values for the data event classes and the expected distribution from MC experiments as derived by investigating the $\sum P_T$ distributions in $e^+p$ (a) and $e^-p$ (b) data.
5 Conclusions

The recent results of model independent searches for new physics performed at the HERA ep collider have been presented. All analyses fully exploit the complete high energy data sample, which amounts to $\sim 0.5 \text{ fb}^{-1}$ per experiment. No convincing evidence for the existence of new phenomena beyond the Standard Model has been observed. Among all event topologies investigated, the largest deviation to the SM expectation is observed by the H1 experiment for isolated lepton events in $e^+p$ collisions. After having analysed all data recorded by H1, this deviation corresponds to a $3 \sigma$ excess of atypical W-like events. This deviation is not confirmed by the ZEUS experiment.

References

THE FIRST A FEW fb$^{-1}$ : POTENTIAL FOR OBSERVING SUSY AND HIGGS

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We outline the expected sensitivity of the discovery for SUSY and Higgs on the first few fb$^{-1}$ operation of LHC. We also briefly introduce the global strategy for finding the SUSY particles, and proposed background estimation methods by data-driven analysis. This report contains both results from ATLAS and CMS experiments. With a few fb$^{-1}$ data, both experiments will discover SUSY particles over wide range of the SUSY parameters, and possibly discover the SM Higgs boson in the first few year operation of LHC.

1 Introduction

The LHC is a proton-proton collider machine with a center-of-mass energy of 14 TeV at CERN in Switzerland. The CMS and ATLAS detectors are located on their collision points in the LHC ring, aiming to cover a wide range of physics programs in the hadron collider. In this letter, we will focus on the potential for the early discovery of SUSY particles and Higgs bosons. Needless to say, understanding of the detector is fundamental for the early operation of LHC running. We however introduce the baseline analysis strategy proposed in Physics Readiness Reports$^{1,2}$ rather than having a discussion about the performance of the detectors.

The key issue of the physics analysis at LHC is how we handle a significant amount of the QCD backgrounds over a signal statistics for the individual search channel. One may require so high rejection to the lepton or photon identifications against QCD jets. One may also require to be very large missing transverse energy and high-energetic multi-jets in the final state to avoid those events. As an example, the typical SUSY events ($gar{g}$) or SM Higgs bosons ($m_H=150$GeV) would be possibly produced with the cross sections $\sigma \sim 10$ pb, which corresponds to the production rate of one event every two minutes in collision at the luminosity $10^{33}$cm$^2$s$^{-1}$. On the other hand, QCD di-jets events with $E_T \geq 100$ GeV is roughly produced with order of $\sigma \sim 1$ mb, which corresponds to 1 kHz at $10^{33}$cm$^2$s$^{-1}$. In $W$ or $Z$ boson production, the production rate
with $100 \sim 200$ Hz is expected. If the discovery is performed by simple counting approaches, the order of $10^{3-5}$ background rejection has to be achieved. This means it requires extremely high rejection performance of the backgrounds while keeping large acceptance of the signal events. In addition, the data-driven analysis methods not to rely on the Monte Carlo prediction by using the real data are desired to estimate the background contribution. Using unique event topology, one may make use of the kinematic variables. An example is the anti-tag method or masking method to evaluate the background shapes by using the similar kinematic configuration in the well-known data control sample. They are described later section.

In this paper, we start addressing the inclusive SUSY searches at first, then describe model specific SUSY searches and then followed by Higgs searches. The target luminosity for the discovery in this context is around $1 \sim 2$ fb$^{-1}$, so that a model independent and cut-based analysis is rather preferred with respect to the majority of importance for the understanding of the detectors.

2 SUSY searches

There are several reasons to motivate a new physics beyond the Standard Model: the electroweak symmetry breaking, hierarchy problems, gravitational force, and so forth. The Supersymmetry (SUSY) is one of the most attractive theory to resolve these issues. And this theory also inspires a good candidate of the cold dark matter. To satisfy the above condition, the supersymmetric particles tend to be degenerated at Tev-scale region. Suppose the $R$-parity is conserved, the model predicts new signatures that can be seen at the experiments as the unique event topology. The SUSY particles sequentially decay into another SUSY particles with lower mass eigenstates. Then, they leave at least two high mass undetected neutralinos afterwards. The typical SUSY events hence would have large missing transverse energy, multi-jets and leptons in the final state. Based upon the allowed parameter space, some benchmark points have been used for the analysis.

The event topology from various SUSY cascade chains is so complicated that the “inclusive” search, based on the signature-oriented analysis, is the first approach for the discovery. Since the inclusive search does not result in such a clear resonance peak, it is often hard to distinguish new phenomena from the SM background processes. The multi-jets searches are sometimes overwhelmed by the SM processes, so that the lepton(s) may be also required in the final state even through the rates are often suppressed. Both ATLAS and CMS analyze the inclusive SUSY events categorized by the number of leptons in the final state. In addition, as the baseline selection for the inclusive searches, a very large missing transverse energy ($\geq 100$ GeV) and high energy multi-jet ($\geq 100$ GeV, at least 3-jets) events are required. Note that one remarkable difference between ATLAS and CMS SUSY inclusive searches is that ATLAS categorizes the events by the number of leptons exclusively, while CMS categorizes inclusively. Readers may refer\textsuperscript{1,2} for details.

Figure 1 and 2 show the missing transverse energy and effective mass distributions for the typical SUSY events at 1 fb$^{-1}$. One can find a huge event excess ($\sim 5000$ events) of the SUSY events. Against of the rapid decrease of the SM backgrounds, the SUSY events tend to result in the large contribution in the tail region over a wide range of parameter space. Thus, understanding of the tail structure of those distributions for the SM backgrounds is the most important task. Since the large missing $p_T$ distribution caused by the SM processes most likely comes from the mis-measurement of the missing $p_T$ such that a jet goes in the crack region of detector and they are weakness part of the theoretical prediction, the data-driven background estimation methods are strongly desired.

The proposed background estimation methods by the data in ATLAS and CMS are as follows. The dominant background is $Z \rightarrow \nu \nu + \text{jets}$, $W^+ + \text{jets}$ and $t\bar{t} + \text{jets}$. The $Z \rightarrow \nu \nu + \text{jets}$
events are estimated by using $Z \rightarrow \mu \mu + \text{jets}$ events. The identified di-muon events are replaced to two neutrinos, then the missing transverse energy is re-calculated. Since the production mechanics are same in both processes, after the correction of the muon reconstruction efficiency and scaling to the neutrino branching ratio, the missing transverse energy is correctly estimated by the $Z \rightarrow \mu \mu$ events in data. Figure 3 justifies to reproduce the original missing $p_T$ distribution of $Z \rightarrow \nu \nu$ process by the $Z \rightarrow \mu \mu$ events after masking their muons. The normalization is simply given by the relative-ratio of the branching ratio of $Z \rightarrow \nu \nu$ to $Z \rightarrow \mu \mu$. For the $t\bar{t}$ and $W+\text{jets}$ backgrounds, the transverse mass distribution is used in the lepton+MET+jets events because the transverse mass has no correlation to the final effective mass distribution. In low transverse mass region, as shown in Figure 4, the $t\bar{t}$ and $W+\text{jets}$ events are dominated. The fraction of $W+\text{jets}$ events to $t\bar{t}$ events does not depend on the cut of the transverse mass distribution, so that this fraction is valid in the signal region. Thus, we can use low transverse mass region as the control sample. The shape is drawn by the control sample, then overall normalization is taken into account the low missing $p_T$ region in which the SUSY signal does not contribute. Finally, the QCD multi-jet events is estimated by using the various kinematical configurations or the anti-tag methods in the object identification. As mentioned, in most of case, the large missing $p_T$ by the QCD multi-jet events is from the mis-measurement of the energy in calorimeters. In such a case, the direction of the missing $p_T$ has strong correlation to the observed jets, where the direction of the JES missing $p_T$ is most likely back-to-back with respect to the leading jet. Thus, one of major estimation method is to use the $\phi$-correlation between jets and missing $p_T$. Using the phase space where QCD jet events are enhanced as control sample, the QCD shape and normalization are driven by the data control sample. One should also note that the semi-leptonically decayed $b$-jets are also major source of the large missing $p_T$ background. The situation is even more complicated, where the dedicated studies are necessary.

In the di-lepton mode, the end-point of the di-lepton invariant mass distribution is used to determine the mass difference of the SUSY particles. For instance, if a squark $\tilde{q}$ has the decay chain of $\tilde{\ell}_L \rightarrow \chi_0^0 \tilde{q} \rightarrow \chi_1^0 t \bar{t}$ and $\chi_0^0$ is LSP, the di-lepton invariant mass edge contains an information of the mass difference of $\chi_0^0$ to $\chi_1^0$ by simple kinematical requirement. Then, the end-point is measured by the flavor subtraction method. In the SM processes, the leptons from weak boson decay have the same branching ratio not depend on the lepton flavor. Thus, since most of leptons come from the weak boson decay when high $p_T$ leptons are required, the SM processes are canceled out by using the relation of $N_{ee} + \beta^2 N_{\mu\mu} - \beta N_{e\mu}$, where $\beta$ is the efficiency correction of muon to electron. We demonstrate this feature in Figure 5. The $t\bar{t}$ background are
Figure 3: Missing transverse energy for the two processes rescaled $Z \rightarrow \mu\mu$ and true $Z \rightarrow \nu\bar{\nu}$ (CMS). Different histogram (black-solid-marker) expresses before rescaling of the $Z \rightarrow \mu\mu$ events.

Figure 4: Transverse mass distribution of signal and background events (ATLAS). Low transverse mass region ($M_T \leq 100\text{GeV}$) is used as the control sample.

The only shown in the plot, but the SM contribution fairly cancels and then the large event excess by SUSY events is left. The signal event statistics is enough to determine the end-point of the mass distribution. This technique is very powerful not only for the background suppression like a requirement of the same signed di-lepton events but also for the direct measurement of the SUSY particles.

Figure 5: Di-lepton invariant mass distribution after flavor subtraction (CMS). The $t\bar{t}$ background are only shown in the plot.

Figure 6: Invariant mass distribution by $H \rightarrow b\bar{b}$ in the SUSY events (ATLAS).

In the particular parameter space, the SUSY particles also allow to decay into Higgs or $Z$ bosons. In this case, the search is to look for the mass resonance in addition to the large missing $p_T$ requirement. Requiring the di-lepton pair for the resonance peak, high mass resolution is obtained. As the other example, Figure 6 shows the mass distribution by $H \rightarrow b\bar{b}$ in the SUSY events. The large missing $p_T$ cut is already imposed, so that the SM background events do not contribute. This mode has extremely high potential to discover the SUSY events in this particular parameter space.

We present the expected discovery potential from ATLAS and CMS in Figure 7, where note that CMS shows only muon channel. With 1 fb$^{-1}$ data, both experiments covers a wide range of SUSY parameter space for the discovery.
3 Higgs searches

Despite the remarkable success of the Standard Model in high energy physics during the recent decades, nothing is known about the source of its fundamental theoretical basis, the Higgs mechanism. The search for the Higgs bosons has been considered to be the most important subject in LHC. In 14-TeV proton-proton collision at LHC, the SM Higgs boson is predominantly produced via gluon-fusion production process. The search channels thus depend on the Higgs mass and its decay mode. Since the recent top-quark and $W$-boson mass measurements at Tevatron Run II suggest to have a low-mass Higgs boson, the low-mass SM Higgs search is the primary topics for early discovery. The promising decay mode are $H \rightarrow \gamma \gamma$, $H \rightarrow ZZ \rightarrow 4l$ and $H \rightarrow WW \rightarrow l\nu l\nu$ modes for the early discovery. The mass resolutions are about 1~$\text{GeV}$ for the $H \rightarrow \gamma \gamma$ and $4l$ channels while the $l\nu l\nu$ channel has the largest production cross section although the invariant mass is not reconstructed. The achievement of a good performance of the lepton and photon identifications and energy resolution are essential for those analyses. We show an example of the reconstructed invariant mass for $4l$ channel in Figure 8. The number of events corresponds to 10 $fb^{-1}$ luminosity. We can see the large resonance peak in $4l$ channel. In $l\nu l\nu$ channel, the transverse mass is reconstructed and predicts the large excess of the signal events.

In addition to the gluon-fusion process, the second largest production mechanisms is the Vector-Boson fusion process (VBF). In the VBF process, the incoming quarks are scattered off the heavy vector bosons to form the jets in the forward and backward region. This provides additional means to suppress the backgrounds. The correlation between two forward jets is used to separate the signal events from the background events. The separation in the pseudorapidity $\Delta \eta$ for two jets strongly suppresses the background events. Also, according to the presence of high energetic two forward jets, the di-jet mass can reject the background events. Furthermore, since the signal event does not much produce jets in central region due to a nature of the color-connection in VBF production process, the central jet veto may be imposed. In Figure 9, the di-photon invariant mass distribution for $H \rightarrow \gamma \gamma+2$jets channel is shown at 30 $fb^{-1}$ data. This is a new result in ATLAS. The analysis with two photon events is categorized by the number of jets. The backgrounds are largely suppressed by the forward jets requirement.

While the VBF selection is applied to reject the backgrounds, the minimal supersymmetric
extension of the Standard Model (MSSM) predicts larger production cross sections for the Higgs bosons over a wide parameter range since they have large couplings with down-type fermions in large tan\(\beta\) region. The production associated with a bottom quark is a promising search channel in this case. The \(H \rightarrow \tau\tau\) and \(H \rightarrow \mu\mu\) channels are analyzed. Also, the b-jet tagging are valid tool to enhance the signal events over QCD backgrounds. The branching ratio of Higgs decaying to \(\tau\) is much larger than that to \(\mu\). On the other hand, the mass reconstruction by the muon pair has much better resolution than that by the \(\tau\)-pair. The sensitivity of this analysis is a trade-off of the signal statistics or high mass resolution.

Finally, we present the SM Higgs discovery potential expected in ATLAS and CMS experiments at 30 fb\(^{-1}\) data in Figure 10 (a) and (b), respectively. Note that at that time being in ATLAS, the \(H \rightarrow \gamma\gamma + 2\text{jets}\) analysis was not included in this sensitivity. Both experiments fully cover in almost full mass-range. Even within \(1-5\) fb\(^{-1}\) data, it can reach at \(5\sigma\) discovery for the high mass region.

![Graphs showing Higgs production cross sections](image)

Figure 8: Reconstructed invariant mass with \(m_H=150\) GeV for 4l channel at 10 fb\(^{-1}\) data (ATLAS).

Figure 9: Di-photon invariant mass distribution for \(H \rightarrow \gamma\gamma + 2\text{jets}\) channel at 30 fb\(^{-1}\) data (ATLAS).

![Graphs showing Higgs discovery potential](image)

Figure 10: The SM Higgs discovery potential expected in ATLAS and CMS experiments at 30 fb\(^{-1}\) data.
4 Summary

The LHC is so-called the discovery machine. With 14-TeV proton-proton collision and high design luminosity, the ATLAS and CMS experiments will discover a new physics and/or Higgs bosons even with a few fb$^{-1}$ data.

For SUSY searches, the inclusive search is the most promising search channel in that the observed signature does not depend on the SUSY models. Many of the parameter space in mSUGRA is expected to be covered for their discovery reach. In the analysis, the understanding of the tail structure of the missing transverse energy is very important to estimate the background contribution. The data-driven background estimation methods are desired. Fortunately, the LHC has an unique opportunity that even abundant SM physics processes themselves become a kind of calibration processes to evaluate the event topology, where the event statistics is not in the SM processes. Using similar event topology, one may obtain a shape of the backgrounds and the associated shape uncertainty. In the di-lepton edge analysis, it indicates the mass difference of the SUSY particles, and reveals the nature of the SUSY model. If SUSY particles exist, this is the most powerful method to determine the model. The search for the resonance peak does not also require the specific model. It will be useful to determine or constrain the parameter space of the model.

In Higgs searches, the $H \to ZZ \to 4l$ and $WW \to lvlv$ channels are the most promising discovery channels in the early data in LHC. In addition to the analysis based on the gluon fusion production process, the analysis with the VBF production processes further increases the signal sensitivity to the backgrounds. In low mass region, the $H \to gg$ channel has the largest sensitivity. Both inclusive and VBF analyses are performed. With 1~5 fb$^{-1}$ data, the SM Higgs boson will be discovered for a wide range of mass-spectrum. Note that within the framework in the MSSM, the MSSM Higgs sensitivity is further increased, according to the strength of the couplings with fermions.

The first collision is expected to be made with 10 TeV on the summer 2008. Then, the 14-TeV physics run will turn on after a short running of 10-TeV commissioning, so that the first physics results will be successively followed by the next year. The Physics Readiness Report from both experiments has been prepared. In ATLAS, the latest report will be available soon, where all physics analysis respect the analysis computing model in the real data analysis. The MC data are distributed and analyzed on the grid as the part of the computing system commissioning (CSC). Most of results presented in the last TDR will be replaced to the CSC results. In this paper, a few selected CSC results are presented. There are many attractive updated results as well as completely brand-new results in the CSC report. And this report will be the final before real-data come out. We, LHC experimentalist, are waiting for the real-data with the most sophisticated prediction with our best knowledge.

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References

THE FIRST FEW FB$^{-1}$; POTENTIAL FOR OBSERVATION OF PHYSICS BEYOND THE STANDARD MODEL

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An important part of the ATLAS and CMS program is to search for new physics beyond the Standard Model. Some of the main ongoing studies are presented by signature, with particular emphasis on those channels already observable with the first collected data. Here only non-super-symmetric models are presented, as susy models are discussed elsewhere\(^1\).

1 Introduction

LHC will offer the possibility to observe or set limits on new physics beyond the Standard Model (SM), using relatively few data. Many models can suddenly be investigated, although early data may not be enough to identify the model that describes the signal.

In the following sections the signatures of some of the main searches at LHC are discussed.

2 Jet Final States

Inclusive di-jet production ($pp \rightarrow 2 \text{ jets} + \text{anything}$) is the dominant LHC process. To lowest order it arises from the $2 \rightarrow 2$ QCD scattering of partons in which only coloured particles are involved in the initial, intermediate and final states.

Di-jet resonances and contact interactions are the two major signals of new physics with di-jets\(^2\). Di-jet resonances produce compelling signals of a new particle at a mass $M$, but require that the incoming parton-parton collision energy to be close to that mass (which must be kinematically accessible). Contact interactions produce more ambiguous signals but come from an energy scale of new physics, $\Lambda$, which can be significantly larger than the available collision energy, thus allowing to probe a wider energy spectra. In both cases the observables
used to study such processes are very simple. In the inclusive jets analysis the number of jets inside an $\eta$ window are counted as a function of the jet $p_T$. For the di-jets study, the two jets with highest $p_T$, and inside a given pseudorapidity region, are selected and counted as a function of the invariant mass of the di-jet system.

2.1 Sensitivity to Contact Interactions: Inclusive Jet-$p_T$ Study

New physics at a scale above the energy scale of the process can be effectively modelled as a contact interaction. The canonical contact interaction studied in hadron collisions arises from the following left-left isoscalar colour-singlet term which is added to the QCD Lagrangian $^3$:

$$L_{qq} = \frac{A g^2}{2 \Lambda^2} (\bar{q} L \gamma^\mu q L)(\bar{q} L \gamma_\mu q L)$$

(1)

where $A = \pm 1$ determines the sign of the interference with QCD, $\Lambda$ is the contact interaction scale and the square of the coupling $g^2$ is by convention set equal to $4\pi\alpha_s$. $\Lambda^\pm$ is a compact notation commonly used to include the choice $A = \pm 1$.

Contact interactions produce a rise in rate, relative to QCD, at high inclusive jet $p_T$ as shown in Figure 1. The figure shows the jet rates, expected for an integrated luminosity of 10 pb$^{-1}$, using a simulation of the CMS experiment. A contact interaction with a scale of $\Lambda^+ = 3$ TeV clearly produces a large rate compared to QCD expectation for jet $p_T > 1$ TeV, event taking into account a 10% energy scale uncertainty.

2.2 Sensitivity to di-Jets Resonances: Di-Jet Mass Spectrum Study

Many models predict narrow di-jet resonances $^4$. In Figure 2 the cross section for an excited quark di-jet resonance to the statistical uncertainties expected on the QCD di-jet background are compared for a luminosity scenario of 10 pb$^{-1}$. The normalisation of the excited quark signal come from the lowest order calculation. Figure 2 illustrates that the di-jets channel is sensitive to an excited quark signal up to several TeV. With only 10 pb$^{-1}$ a 2 TeV excited quark signal begins to emerge above the statistical error bars with a total significance of 4.1, neglecting systematic uncertainties.
In a hadron collider the final states with leptons are a clear signature for many processes, in particular di-lepton final states are naturally the best candidate for discovery of new physics beyond Standard Model. Many models predict either resonances or deviation from the SM differential cross section \((d\sigma/dm)\) of the process.\

The main characteristics of the signal are the high momenta of the leptons, which are also isolated, and the large invariant mass of the lepton pair. The most important background is the Drell-Yan process, although its cross section is vanishing at the energy scale at which new physics is expected. The same signature is shared by many models and this would make difficult the identification of the correct theory using only early data. From the experimental point of view, due to the characteristics of the particles in the final state, the detection, reconstruction and identification of such leptons can be at the limits of the performance of the apparatus, depending on the energy scale at which the new process arises. Therefore a high-level understanding of the alignment and calibration of the detector is fundamental for the discovery.

### 3.1 Resonances in Final States: New Neutral Gauge Bosons

Additional heavy neutral gauge bosons \((Z')\) are predicted in many superstring-inspired and grand unified theories, as well as in dynamical symmetry breaking and little Higgs models. However, there are no reliable theoretical predictions of the \(Z'\) mass scale. Current lower limits on the \(Z'\) mass are (depending on the model) of the order of \(600 - 900\) GeV/c^2.

The \(Z'\) most frequently discussed and whose properties are representative of a broad class of extra gauge bosons are:

- **\(Z_{SSM}\)** within the Sequential Standard Model (SSM), which has the same couplings as the Standard Model \(Z^0\).
- **\(Z_\psi, Z_\eta\)** and **\(Z_X\)**, arising in \(E_6\) and \(SO(10)\) GUT groups.
- **\(Z_{LRM}\)** and **\(Z_{ALRM}\)**, arising in the framework of the so-called left-right and alternative left-right models.

The LHC offers the opportunity to search for \(Z'\) bosons in a mass range significantly larger than \(1\) TeV/c^2, already with the first data. In Figure 3 the summary plot shows that already with 100 pb^{-1} a region not yet explored by Tevatron experiments can be studied.

Figure 2: (Left) Di-jets differential cross section expected from QCD for \(|\eta| < 1\) as a function of the di-jets invariant mass, for generated jets (points), jets (triangles), and corrected jets (open boxes). (Right) Fractional difference with respect to QCD as a function of di-jets invariant mass, in case of excited quark signal.

Figure 3: Summary plot showing the expected signal strengths for \(Z'\) bosons with 100 pb^{-1}.
The $Z'$ is not the only neutral vector boson that can be seen in leptonic channels. Randal-Sundrum (RS) models predicts massive Kaluza-Klein (KK) modes of the graviton ($G_{RS}$).

Most collider physics phenomenology done with warped extra dimensions so far is based upon one very specific model, the original simple scenario called RS1. In RS1, the Standard Model is replaced at TeV scale by a new effective theory in which gravity is still very weak, but there are exotic heavy spin-2 particles.

At LHC the KK gravitons of RS1 would be seen as di-fermion or di-boson resonances. In particular, with early data, only the first excitation of the RS graviton can be accessible.

In Figure 4-(left) the reach of the CMS experiment for RS1 graviton in muon channel, is shown as a function of the coupling parameter ($k/M_{PL}$, where $k$ is the curvature of the warped extra dimension and $M_{PL}$ is the Planck mass in 5 dimensions) and the graviton mass. The ranges of the expected variations due to the systematic uncertainties are also drown: 1 fb$^{-1}$ is enough to explore a wide part of the region allowed by the theory.

Early data could not be enough to perform detailed angular distribution studies (crucial in order to distinguish a spin-1 particle, like the $Z'$, with respect to a spin-2 one, like the $G_{RS}$), however some handle is given by looking at resonances in some final states, which are precluded.
in the other models due to the nature of the new particle. In fact, while the the Z’ cannot decay into a pair of vector bosons, the RS graviton can. In Figure 4-(right) the sensitivity of the analysis to the $G_{RS} \rightarrow \gamma\gamma$ channel is also drown and shown to be comparable with the leptonic final states, thus allowing a cross-check of the resonance, even with a few collected data. The $G_{RS}$ branching ratio to photons is roughly twice that of electrons or muons, however the reach for low coupling and graviton mass is comparable between di-leptons and di-photons due to the QCD and prompt photon backgrounds in the photon channel which are harder to efficiently suppress. For higher masses and coupling the di-photon is leading the reach due to the higher branching ratio. The di-muon channel is trailing the reach compared to the di-electrons merely due to resolution.

4 Two Leptons and Two Jets Final State

4.1 Heavy Majorana Neutrinos and right-handed bosons

The two leptons and two jets final states can be a clear signature of process described by left-right symmetric model $SU_c(3) \otimes SU_L(2) \otimes SU_R(2) \otimes U(1)^{14,15}$. The model embeds the SM at the scale of the order of 1 TeV and naturally explains the parity violation in weak interactions as a result of the spontaneously broken parity. It necessarily incorporates three additional gauge bosons $W_R$ and $Z'$ and the heavy right-handed Majorana neutrino states $N$. The $N$ particles ($N_l$) can be the partners of the light neutrino states $\nu_l (l = e, \mu, \tau)$ and can provide their non-zero masses through the seesaw mechanism$^{16}$. The crosssection of $pp \rightarrow W_R \rightarrow l + N_l + X$, where $N_l \rightarrow l + j_1 + j_2$, depends on the value of the coupling constant $g_R$, the parameters of the CKM mixing matrix for the right-handed sector, the $W_R$-$W_L$ and $Z^\prime$-$Z$ mixing strengths, and the masses of the partners $N_l$ of the light neutrino state. In the study presented here the mixing angles are assumed small, the right-handed CKM matrix identical to the left-handed one and $g_R = g_L$. Finally it is assumed that only the lightest $M_{N_e}$ is reachable at LHC. In the case of degenerated masses of $N_l$, the channels with $\mu$’s and $\tau$’s are open resulting in the increase of the cross section of the process studied here by a factor of 1.2. The two major backgrounds considered in this study are the inclusive production of $Z$
Figure 6: Integrated luminosity needed to discover (at 5 sigma level) a $W'$ boson, depending on its mass, at ATLAS (left) and CMS (right).

and $t\bar{t}$. In the event selection two isolated electrons and at least two jets are required.

The 5 sigma discovery contour in the $(M_{W_R}, M_{N_e})$ plane is shown in Figure 5 for an integrated luminosity of 1, 10 and 30 fb$^{-1}$ (CMS experiment simulation 5). With 1 fb$^{-1}$ a 5 sigma observation of $W_R$ and $N_e$, with masses up to 2 TeV/c$^2$ and 1 TeV/c$^2$ respectively can be achieved.

5 Leptons and Missing Energy Final States

As mentioned in Section 3.1 many models predict additional heavy gauge boson, including charged particle. Here are presented the detection capabilities for a hypothetical heavy partner of the Standard Model $W$, a charged spin-1 boson $W'$, with the properties from the Reference Model by Altarelli 18. In this model, the $W'$ is a much massive copy of the $W$, with the very same left-handed fermionic couplings (including CKM matrix elements), while there is no interaction with the Standard Model gauge bosons or with other heavy gauge bosons as a $Z'$. Thus the $W'$ decay modes and corresponding branching fractions are similar to those for the $W$. In hadron collisions $W'$ bosons can be created through $q\bar{q}$ annihilation, in analogy to $W$ production. Previous searches for the reference $W'$ at LEP and at the Tevatron give rise to lower bounds approaching 1 TeV$^{12}$.

Given that the $W'$ boson has a large mass, it is likely to be produced without transverse momentum. Due to a boost along the $z$-axis, the angle between the muon and the neutrino might be different from $\pi$ in the laboratory system. However, the angle in the transverse plane stays invariant under boosts along the $z$-axis. Therefore the signature of a $W'$ event is high energy isolated muon, together with a large amount of missing energy pointing to the opposite direction in the transverse detector plane. Due to the small transverse momentum of the $W'$ boson, the transverse momentum of the muon and the missing transverse energy are of similar magnitude.

In Figure 6 the discovery potential of both ATLAS$^{19}$ and CMS$^5$ is shown. Less than 1 fb$^{-1}$ is needed to find a signal with a significance at 5-sigma level.

6 Black Hole

One of the consequence of large extra dimension is the possibility to produce microscopic black hole at LHC energy$^{20,21}$. From a semi-classical calculation the cross section of the black hole
production can be written as
\[ \sigma(M_{BH}) = \pi r_s^2(4+n) \]
where \( r_s(4+n) \) is the Schwarzschild radius in “4+n” dimensions.

Considering a black hole mass much larger than the Planck mass in 4+n dimensions \( (M_{PL}) \), and assuming the latter to be of the order of the TeV scale, then \( \sigma(M_{BH}) \sim \text{pb}^n \).

The black holes have a very short life time, predicted to be of the order of \( 10^{-12} \) fs and are expected to evaporate democratically by emission of all particle types that exist in nature, independent of their quantum numbers or interaction properties. Therefore they can be a source of new particles. Black holes would also be able to provide the possibility of probing quantum gravity.

Requiring \( M_{BH} \) to be larger than \( M_{PL} \), potentially the black hole observation can be accomplished within an integrated luminosity of \( 1 \text{ fb}^{-1} \) in all cases of \( n \) and if \( M_{PL} \) is less than 5 TeV, as shown in Figures 7.

7 Summary

The Large Hadron Collider will give the possibility to shade light on new physics and models. Here some examples of the phenomena that could be discovered with few data collected at ATLAS and CMS have been reviewed. In Table 1 a summary of such discoveries are reported. Although the understanding of the detector will be crucial for a claim of an observation of a new signal, both experiment show a high discovery potential already with an integrated luminosity of few \( \text{fb}^{-1} \).

---

*Recent studies claims that due to the “Apparent Horizon” effect, the event horizon is not formed as fast as needed, thus a large fraction of the initial energy could escape before the black hole is formed. This implies that more partonic energy is needed to form the black hole than the one predicted by the naive semi-classical calculation. In such a scenario the black hole cross section is a few orders of magnitude smaller than the above calculation.*
### Table 1: Summary of the principal discoveries accessible with the first data taken at LHC.

<table>
<thead>
<tr>
<th>Model</th>
<th>Mass Reach (TeV)</th>
<th>L (pb⁻¹)</th>
<th>Early Systematic</th>
</tr>
</thead>
<tbody>
<tr>
<td>Contact interaction</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>ALRM</td>
<td>M ~ 1</td>
<td>10</td>
<td>Jet efficiency and energy scale</td>
</tr>
<tr>
<td>SSM</td>
<td>M ~ 1</td>
<td>20</td>
<td>Alignment</td>
</tr>
<tr>
<td>LRM</td>
<td>M ~ 1</td>
<td>30</td>
<td></td>
</tr>
<tr>
<td>E₆, SO(10)</td>
<td>M ~ 1</td>
<td>30-100</td>
<td></td>
</tr>
<tr>
<td>Technirho</td>
<td>M ~ [0.3]</td>
<td>100</td>
<td>Jet energy scale</td>
</tr>
<tr>
<td>Axigluon or Colouron</td>
<td>M ~ [0.7,3.5]</td>
<td>100</td>
<td>Jet energy scale</td>
</tr>
<tr>
<td>Excited quark</td>
<td>M ~ [0.7,3.6]</td>
<td>100</td>
<td>Jet energy scale</td>
</tr>
<tr>
<td>E₆ di-quarks</td>
<td>M ~ [0.7,4]</td>
<td>100</td>
<td>Jet energy scale</td>
</tr>
<tr>
<td>mUED</td>
<td>M ~ 0.3 - 0.6</td>
<td>10-1000</td>
<td>MET, jet/photon energy scale</td>
</tr>
<tr>
<td>ADD real GKK</td>
<td>M_D ~ 1.5 (n=3), 1 (n=6)</td>
<td>100</td>
<td>MET, jet/photon energy scale</td>
</tr>
<tr>
<td>ADD virtual GKK</td>
<td>M_D ~ 4.3 (n=3), 3 (n=6)</td>
<td>100</td>
<td>Alignment</td>
</tr>
<tr>
<td></td>
<td>M_D ~ 5 (n=3), 4 (n=6)</td>
<td>1000</td>
<td></td>
</tr>
<tr>
<td>RS1</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>di-jets</td>
<td>M_G ~ [0.7,0.8], c=0.1</td>
<td>100</td>
<td>Jet energy scale, alignment</td>
</tr>
<tr>
<td>di-muons</td>
<td>M ~ [0.8,2.3], c=[0.01,0.1]</td>
<td>1000</td>
<td></td>
</tr>
</tbody>
</table>

### References

1. S. Tsuno, this proceedings.
4. K. Gumus et al., CMS Note 2006/070 (2006), and references therein.
13. L. Randall, this proceedings.
SEARCHES IN LEPTON FINAL STATES

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Abstract

Searches for new physics in lepton final states at the Tevatron are summarized in this paper. I describe the searches for supersymmetric particles such as chargino/neutralino and tau-sneutrino. I also describe searches for excited gauge bosons ($W'$, $Z'$) and excited electrons.

1 Introduction

The standard model (SM) of physics gives a successful description of natural phenomena. It is applicable over a wide range of energy scales. However, the standard model has a few unanswered questions. It does not include gravity, or describe the origins of masses of the fundamental particles. It suffers from the hierarchy problem, and gives no hints about the nature of dark matter or dark energy which constitute 96% or our universe.

There are several proposed models of new physics beyond the standard model. Supersymmetry (SUSY) is a new proposed symmetry between fermions and bosons. It posits the existence of boson superpartners to all SM fermions and vice versa. It solves the hierarchy problem and in certain models suggests an attractive dark matter candidate particle. Other models such as left-right-symmetric models ($SU(2)_L \times SU(2)_R$) or Grand Unification Theories which unify the electroweak and strong forces predict additional gauge bosons. The Randall-Sundrum model of warped extra dimensions predicts the existence of a massive spin-2 particle.

The search for new physics in lepton final states has many advantages over other channels. Lepton final states are relatively clean and free from backgrounds. Since identifying leptons is well understood at the CDF and DØ detectors, the backgrounds in lepton final states from SM processes are straightforward to estimate. Many new models predict leptonic final states and thus these analyses are sensitive to a wide variety of models beyond the standard model.

Results shown here use from 1 to 2.5 fb$^{-1}$ of data from the Tevatron for each CDF and DØ.
Table 1: CDF results of a search for $\tilde{\chi}_1^\pm \tilde{\chi}_2^0 \rightarrow llE_T$. The signal quoted is for mSUGRA model with parameters $m_0 = 60 \text{ GeV}/c^2$, $m_{1/2} = 190 \text{ GeV}/c^2$, $\tan(\beta)=3$, $A_0 = 0$, $\mu > 0$.

<table>
<thead>
<tr>
<th>Channel</th>
<th>Expected Signal</th>
<th>Background</th>
<th>Observed</th>
</tr>
</thead>
<tbody>
<tr>
<td>3 tight leptons</td>
<td>2.3±0.3</td>
<td>0.5±0.1</td>
<td>1</td>
</tr>
<tr>
<td>2 tight + 1 loose</td>
<td>1.6±0.2</td>
<td>0.3±0.04</td>
<td>0</td>
</tr>
<tr>
<td>1 tight + 2 loose</td>
<td>0.7±0.1</td>
<td>0.1±0.03</td>
<td>0</td>
</tr>
<tr>
<td>2 tight + track</td>
<td>4.4±0.6</td>
<td>3.2±0.7</td>
<td>4</td>
</tr>
<tr>
<td>1 tight + 1 loose + 1 track</td>
<td>2.4±0.3</td>
<td>2.3±0.6</td>
<td>2</td>
</tr>
</tbody>
</table>

Table 2: DØ results of a search for $\tilde{\chi}_1^\pm \tilde{\chi}_2^0 \rightarrow llE_T$. The signal quoted is for mSUGRA-inspired model with parameters $m_0 = 88 - 121 \text{ GeV}/c^2$, $m_{1/2} = 182 - 221 \text{ GeV}/c^2$, $\tan(\beta)=3$, $A_0 = 0$, $\mu > 0$. The RunIIb analysis is a new analysis.

<table>
<thead>
<tr>
<th>Channel</th>
<th>Expected Signal</th>
<th>Background</th>
<th>Observed</th>
</tr>
</thead>
<tbody>
<tr>
<td>$ee$ + Track (RunIIb)</td>
<td>0.5-2.1</td>
<td>1.0±0.3</td>
<td>0</td>
</tr>
<tr>
<td>$ee$ + Track (RunIIa)</td>
<td>1.7-4.7</td>
<td>0.8±0.7</td>
<td>0</td>
</tr>
<tr>
<td>$\mu\mu$ + Track (RunIIa)</td>
<td>0.5-2.5</td>
<td>0.3±0.3</td>
<td>2</td>
</tr>
<tr>
<td>$e\mu$ + Track (RunIIa)</td>
<td>2.0-2.6</td>
<td>0.9±0.4</td>
<td>0</td>
</tr>
<tr>
<td>$\mu^+\mu^-$ (RunIIa)</td>
<td>0.6-3.8</td>
<td>1.1±0.4</td>
<td>1</td>
</tr>
</tbody>
</table>

2 Supersymmetry searches

2.1 Chargino-Neutralino

In R-parity\(^a\) conserving models of supersymmetry, the associated production of chargino and neutralino gives rise to a distinctive signature. The chargino ($\tilde{\chi}_1^\pm$) and neutralino ($\tilde{\chi}_2^0$) each decay to leptons along with invisible particles ($\tilde{\chi}_1^\pm \rightarrow l^\pm \nu \tilde{\chi}_1^0$, $\tilde{\chi}_2^0 \rightarrow l^\pm l'^\mp \tilde{\chi}_1^0$), giving a final state with three leptons and a momentum imbalance (missing $E_T$ or $E_T$) in the detector.

CDF conducted a analysis which looked in five final states defined by purity of leptons used for the final states. Tight leptons have stricter selections and lower backgrounds, looser leptons are less pure, and tracks are simply charged particles. The results are summarized in Table 1. In this table lepton refers to electrons or muons. The final cross-section $\times$ branching ratio limits are shown in Figure 1. This analysis was performed with 2 fb$^{-1}$ of data. It improves the CDF published limits\(^1\). Charginos with mass below 145 GeV/$c^2$ are excluded. These are first direct limits on mSUGRA chargino mass from the Tevatron.

DØ has added to their published result\(^2\) an analysis which looks in the final state with two electrons and a track. The results are summarized in Table 2 and the limit is shown for an ‘mSUGRA-inspired’ scenario in Figure 2. The analyses performed by DØ used between 0.9 and 1.7 fb$^{-1}$. Charginos with mass below 145 GeV/$c^2$ are ruled out in this scenario.

2.2 Tau-sneutrino

Imposing $R_P$ conservation is not necessary for supersymmetry. An analysis probing $R_P$ violating SUSY is the DØ search for the superpartner to the tau-neutrino, the tau-sneutrino ($\tilde{\nu}_\tau$). The supersymmetric Lagrangian can then be modified to include terms such as

$$W_{RPV} = \frac{1}{2} \epsilon_{ab} \lambda_{ijk} L^a_i L^b_k E_k + \epsilon_{ab} \lambda'_{ijk} L^a_i Q^b_j D_k$$

\(^a\)A multiplicative quantum number defined as $R_P = (-1)^{2S+3B+L}$. 

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where the LLE and LQD terms represent lepton flavor violating interactions. The analysis is a direct search for resonant production of sneutrinos decaying into an electron and a muon performed under the hypothesis that the third-generation sneutrino (ντ) is the lightest supersymmetric particle and dominant, namely by assuming that all couplings but λ_{311} and λ_{312} = λ_{321} are zero. This analysis used 1 fb^{-1} of data to set limits on the R_P violating couplings as a function of $\tilde{\nu}_\tau$ mass. Figure 3 shows the limits.

### 3 Non-SUSY searches

#### 3.1 Excited electrons

A possible way to explain the observed mass hierarchy of the three generations in the SM is compositeness. According to this approach, a quark or a lepton is a bound state of three fermions or of a fermion and a boson. DØ performs a search for an excited electron with 1 fb^{-1} of data. Single production of an excited electron ($e^*$) is considered in association with an electron via a four-fermion contact interactions, with the subsequent electroweak decay of the $e^*$ into an electron and a photon. The $ee\gamma$ final state is fully reconstructable and nearly background-free. The selection is optimized for the mass of the $e^*$ using two variables $\Delta R_{e\gamma}$ (separation between electron and photon) and $M_{e\gamma}$ (invariant mass of electron and photon). The limits on the cross section are shown in Figure 4. For the scale for contact interactions to be $\Lambda = 1$ TeV, excited electron masses below 756 GeV are excluded at the 95% C.L.

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**Figure 1:** Figure shows the cross section $\times$ branching ratio limits from the CDF $\tilde{\chi}_\pm^0$ search for the mSUGRA model. $\tilde{\chi}_\pm^0$ with mass below 145 GeV/c$^2$ is excluded.

**Figure 2:** Figure shows the cross section $\times$ branching ratio limits from the DØ $\tilde{\chi}_\pm^0$ search for an ‘mSUGRA-inspired model’. $\tilde{\chi}_\pm^0$ with mass below 145 GeV/c$^2$ are excluded.

**Figure 3:** Figure shows the limits on the $R_P$ violating couplings on the left side and cross section $\times$ branching ratio limits on the right from the DØ $\tilde{\nu}_\tau$ search as a function of the $\tilde{\nu}_\tau$ mass.
3.2 New gauge bosons

Search for $Z'$

The E6 model, which unifies the forces in the SM into a E6 gauge group, predicts the presence of additional neutral spin-1 bosons. These new bosons are referred to as $Z'$s and they can mix with some arbitrary angle. Changing the value of the mixing angle used to benchmark the model gives the following six states: $Z'_{\eta}$, $Z'_{\chi}$, $Z'_{\psi}$, $Z'_{I}$ and $Z'_{t}$. The search based on 2.5 fb$^{-1}$ and is carried out as a search for a narrow resonance in the dielectron ($e^+e^-$) final state with a mass range from 150 GeV/$c^2$ to 1050 GeV/$c^2$. The invariant mass distribution is shown in Figure 5 with the most significant excess (3.8$\sigma$) at 240 GeV/$c^2$ (inset).

Search for $W'$

Left-right-symmetric models, along with E6 models, also introduce additional gauge bosons. In the most general case, a new gauge group is comprised of a new mixing angle $\zeta$, new couplings
Figure 7: Figure shows the transverse mass ($m_T$) distribution in the $e\nu$ final state for the DØ $W'$ search. The signal from a 500 GeV/c², and 1.1 TeV/c² $W'$ boson is also shown.

Figure 8: Figure shows the cross section limits on the $W'$ production from a search at DØ. $W'$ with standard model like couplings is ruled out below 1 TeV.

to fermions and a new CKM matrix $U'$. This analysis, based on 1 fb⁻¹, assumes no mixing, couplings equal to the SM couplings and the same CKM matrix as the SM. The width of the $W'$ is assumed to scale with its mass, $m_{W'}$. The final state is an electron and a neutrino ($W' \rightarrow e\nu$). The transverse mass ($m_T$) of the electron and neutrino is constructed, and the tail of the $m_T$ distribution is searched for possible excesses. Figure 7 shows the $m_T$ distribution with signal shown for two possible masses of the $W'$. Figure 8 shows the cross section $\times$ branching ratio limits on $W'$ production. A $W'$ boson with mass below 1 TeV is ruled out at 95% C.L.

4 Conclusions

CDF and DØ have searched for new physics in lepton final states with more than 2 fb⁻¹. The leptonic final states provide a rich set for searches for physics beyond the standard model. No signs of new physics have been found yet, but new constraints have been set on new-physics models. With the improving Tevatron performance and optimal working of the two experiments, next year will push the boundaries of physics even further.

References

5. DØ collaboration, http://www-d0.fnal.gov/Run2Physics/WWW/results/np.htm
MODULI, ANOMALOUS U(1) AND LHC PHENOMENOLOGY

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In these proceedings we present a phenomenological model of moduli stabilization where the uplift of the cosmological constant to zero is provided by a Fayet-Iliopoulos sector. In the presence of an extra anomalous U(1) gauge symmetry, fields with the same features of the "messengers" in gauge-mediation scenarios are naturally introduced. The original phenomenology induced at low-energy in this kind of mixed gravity-gauge mediation presents a superpartners spectrum efficiently compressed and a good dark matter relic density compatible with WMAP bounds.

1 Introduction

A quite general result concerning the high-energy models in the presence of extra-dimensions, it’s that when one reduces to the four dimensional space time, new fields appear in the model, often parametrizing flat directions. This kind of fields are called moduli. For example, in the following, the modulus $T$ will be the superfield representing the fluctuations of the overall internal volume. Since the vev of the moduli are strictly related to physical parameters, it is compelling to find a mechanism to provide them a potential in order to properly define a minimum.

Recently, Kachru et al. 1 (KKLT) proposed a strategy to stabilize the moduli in the context of Type IIB string theory orientifold, following earlier work 2. The KKLT set-up involves different logical steps to achieve a supersymmetry breaking Minkowski vacuum, while stabilizing all moduli. All steps except the last one (uplifting the vacuum energy through the addition of anti D3-branes) can be understood within the context of an effective supergravity. Other works changed this point by insisting on the possibility of using F-terms or D-terms of matter fields in a decoupled sector to perform the uplift. In this proceedings we will show an alternative way to obtain de Sitter space with a TeV gravitino mass by using a Fayet-Iliopoulos (FI) model as uplift sector 3.
It is important to stress that this kind of approach is not an attempt to solve the problem of the cosmological constant, but instead it is meant to be a pragmatic program: The aim is to look at the low-energy phenomenology starting from a high-energy model, imposing some constraints due to consistency requirements and fixing some basic phenomenological inputs (like \( \Lambda_c = 0 \) and the value of the mass of the gravitino).

The peculiarity of our model is due to the presence of one extra \( U(1)_X \) gauge symmetry in the game. This kind of symmetry appears in a very natural way in many compactification of extra-dimensional models, and in the most general case, all the fields entering in the stabilization and uplifting procedure can be charged under it.

More precisely the \( U(1)_X \) transformations for the gauge superfield \( V_X \), the matter chiral superfields \( \Phi_i \) and the modulus \( T \) have the form:

\[
\delta V_X = \Lambda_X + \bar{\Lambda}_X , \quad \delta \Phi_i = -2q_i \Phi_i \Lambda_X , \quad \delta T = \delta_{GS} \Lambda_X ,
\]

where \( q_i \) are the charges of the fields \( \Phi_i \) and \( \delta_{GS} \) a suitable constant. Gauge invariance forces the Kahler potential for the modulus \( T \) to be of the form \( K(T + \bar{T} - \delta_{GS} V_X) \) and this leads in turn to the FI term

\[
\xi_{FI} = \frac{3\delta_{GS}}{2} \frac{1}{T + \bar{T}} .
\]

The presence of this \( T \)-dependent FI term is crucial, because the corresponding non-vanishing D-term, even if it doesn’t change directly the cosmological constant, at the same time induces the suitable F-terms performing the uplift and play an important role for the corresponding low-energy phenomenology.

2 Uplifting and Gravity mediation

2.1 The model

The supergravity model we focus on, is defined in terms of the modulus \( T \) and two scalar fields \( \Phi_{\pm} \) of opposite charges under \( U(1)_X \), by the superpotential:

\[
W = W_0 + m \phi_+ \phi_- + a \phi_-^q e^{-bT} .
\]

In the presence of a charged modulus \( T \), the last term in Eq. 3 is the right gauge invariant version of the KKLT gaugino condensation contribution to the superpotential. The usual negative \( W_0 \) constant, the presence of the charged fields \( \Phi_{\pm} \), their mass term and the interaction term between \( T \) and \( \Phi_- \), are motivated by stringy argument and can be microscopically defined in the type IIB orientifold setup, in terms of fluxes, intersecting branes and stringy instantons effects.

Using a conventional Kahler potential of the form \(^a\) \( K = |\phi_+|^2 + |\phi_-|^2 - 3 \ln(T + \bar{T}) \) and considering a region of the parameters space where

\[
\delta_{GS} \sim 1 \quad , \quad m \ll M_P \quad , \quad W_0 \ll M_P^3 \quad , \quad a e^{-bT} \ll W_0 \ll m \quad (4)
\]

hold, the minimization of the scalar potential given by standard supergravity formula in terms of the auxiliary fields \(^b\) \( F_i \) and \( D \)

\[
V(\phi_+, \phi_-, T) = F_T^T F_T + F_-^T F_- + F_+^T F_+ + \frac{\alpha_2}{2} D^2 - 3e^K |W|^2
\]

\(^a\)The Kahler metric of the charged fields \( \Phi_{\pm} \) can be more complicated and can also depend on \( T \). We checked explicitly that with the Kahler potential \( K = -3 \ln(T + \bar{T} - |\Phi_-|^2 - |\Phi_+|^2) \) we obtain very similar results.

\(^b\)Here we use the definitions \( F_i = e^{K/2} D_i W \) and \( D = K_+ \Phi_+ - K_- \Phi_- + \xi_{FI} \). As usual, the indices are raised and lowered by using the Kahler metric.
shows that the vacuum of the theory breaks supersymmetry. Moreover it can give a zero cosmological constant $\Lambda_c$ if the parameters $m$ and $W_0$ satisfy $|m\phi_-| \approx \sqrt{3} |W_0|$, considering that the vev of $\Phi_-$ is fixed proportional to $\xi_{FI}$ by the FI mechanism. In Fig. 1 we show the shape of the scalar potential in the two cases $m = 0$ and $m$ tuned in order to have $\Lambda_c = 0$. In our case the uplift of the AdS to a Minkowski vacuum is mainly provided by $F_+$, but the crucial point is that this is induced by the non-vanishing $D$-term. Moreover, since the superpotential for $T$ is not completely decoupled from the supersymmetry breaking sector, due to the interaction term with $\Phi_-$, $F_T$ is bigger than the values obtained in typical sequestered $F$-term uplifting models, even if the numerical value for $T$ is very close to its supersymmetric solution.

2.2 First phenomenological results

Even if the supersymmetry is broken in a hidden sector, (super)gravity interactions communicate this breaking to the observable sector, that we take for simplicity to be the Minimal Supersymmetric Standard Model (MSSM). In particular, irrespective on the string theory brane configuration giving our model as effective field theory, if magnetic fluxes are turned on, the coupling constants of the MSSM gauge fields contain a $T$-dependence. This implies that under very general assumptions, a mass for the gaugino fields is directly provided, and it as the form

$$(M_a)_{grav.} \approx \frac{F^T}{T}. \quad (6)$$

Concerning the scalars soft masses, the relevant quantity for computing the soft terms is the coupling of the matter fields metric $K_{ij}$ to the SUSY breaking fields. This can in turn be parameterized as

$$K_{ij} = (T + \bar{T})^{n_i} \left[ \delta_{ij} + (T + \bar{T})^{m_{ij}} |\phi_+|^2 Z_{ij}^{\prime} + (T + \bar{T})^{p_{ij}} |\phi_-|^2 Z_{ij}^{\prime\prime} + (T + \bar{T})^{q_{ij}} (\phi_+ \phi_- Z_{ij}^{\prime\prime} + \text{h.c}) + O(|\phi_i|^4) \right], \quad (7)$$

but the final results is quite simple:

$$\langle \tilde{m}_{0_{ij}}^2 \rangle_{grav.} = m_{3/2}^{\frac{3}{2}} \left[ \delta_{ij} + (\ldots) \right]. \quad (8)$$

Here $(\ldots)$ represents subleading terms if the weights $n_i$ and $m_{ij}$ satisfy the relation $r_{ij} = m_{ij} + (n_i - n_j)/2 \leq -1$. Actually this relation is strongly motivated from the string theory point of view, and the only dangerous case is $r_{ij} = 0$, where a flavour dependence and FCNC
effects could arise.
The first important result is that in our non completely decoupled model, $F^T$ is greater than the usual KKLT-like models and we obtain a splitting between the scalar and gaugino masses smaller by a factor of two. Nonetheless this implies that the one-loop contributions (AMSB) are less important here compared to the tree-level ones.

Finally, in the presence of gravity mediation, trilinear couplings are produced in a similar way and the $\mu$ and $B_\mu$ parameters for the Higgs sector can be generated at the TeV scale through a Giudice-Masiero mechanism.

3 Anomalies and Mixed mediation

3.1 Anomalies and messengers

As introduced in the previous section, the T-modulus transforms under the extra anomalous $U(1)_X$. Moreover, in a very generic way, it is related to the MSSM gauge coupling via a dependence on $T$ of the gauge kinetic functions. Therefore this implies that under $U(1)_X$ gauge transformation, mixed $U(1)_X \cdot G_a$ anomalous terms are produced (with $G_a$ subgroup of the SM gauge group) and a chiral spectrum is required. More precisely, there should be fields carrying Standard Model quantum numbers charged under the additional $U(1)_X$.

Since quarks and leptons carrying $U(1)_X$ charge should imply various phenomenological problems (related to very large soft masses), the most natural possibility is to keep uncharged under $U(1)_X$ the SM fields and to introduce additional heavy fields with the right quantum numbers. These fields have exactly the features of the ”messengers” fields in gauge-mediated scenario (GMSB).

Since the cancelation of the anomaly implies a positive $U(1)_X$ charge for the messengers $^c M$ and $\tilde{M}$, a natural gauge invariant superpotential is

$$W_{\text{mess}} = \lambda \phi_- M \tilde{M}, \quad (9)$$

which naturally pushes the messenger scale up to the GUT scale.

In the usual gauge mediated scenario, adding messengers to a supersymmetry breaking sector generates a new supersymmetric vacuum. However, in our case, the vacuum presented in the previous section is preserved by $U(1)_X$.

Nonetheless another very important new point with respect to the standard gauge-mediation concerns the contribution to the scalar masses. Indeed, as pointed out by Poppitz and Trivedi\(^6\), when the supertrace of the messenger mass matrix is non-vanishing, a new UV divergent term appear at the quantum level and play a very crucial role in what follows.

More precisely, in our case the supertrace is proportional to the $U(1)_X$ D-term:

$$\left(\text{Str}M^2\right)_{\text{mess}} \sim 2 g_X^2 \frac{m^2}{(T + \bar{T})^3} \neq 0. \quad (10)$$

While this does not affect the GMSB one-loop contribution for the gaugino masses,

$$M_{a\text{GMSB}}^G \simeq m \frac{g_a^2}{(T + \bar{T})^{3/2}} \frac{1}{16\pi^2} \left( \frac{\phi_+}{\phi_-} \right), \quad (11)$$

it changes significantly the two-loop soft masses for the scalar superpartners

$$\left(\tilde{m}^0_{\text{GMSB}}\right)^2 \simeq \frac{m^2}{(T + \bar{T})^3} \sum_a \frac{g_a^4}{128\pi^4} \frac{C_a}{C} \left[ 1 - \log \left( \frac{\Lambda_{\text{UV}}}{\lambda \phi_-} \right)^2 + \left( \frac{\phi_+}{\phi_-} \right)^2 \right], \quad (12)$$

\(^c\)The messengers fields are chosen in a complete vector-like SU(5) in order to preserve the perturbative gauge coupling unification.
Table 1: Low energy sample spectra for two different choices of the high-energy parameters, in both the cases of simple gravity and mixed gravity-gauge mediation. All superpartner masses are in GeV, whereas $W_0$, $m$ and $t$ are in Planck units. The last line correspond to the relic abundance, within WMAP bounds in each case.

<table>
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<th>(B) Gravity</th>
<th>(B) Mixed</th>
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where $C_a$ is the Casimir in the MSSM scalar fields representations. From a low energy - GMSB point of view, the logarithmic divergence shows the scale beyond which ”new physics” occurs, since there the scalars can become tachyonic. In our case, we can take $\Lambda_{\text{UV}}$ as the Planck scale, and for suitable values of $\lambda$ the GMSB contribution is actually negative.

### 3.2 Phenomenology

Once the messengers are introduced in the model, in the complete framework scalar and gaugino masses get contributions both from gravity and gauge mediation diagrams

$$ (\tilde{m}_0^2) = (\tilde{m}_0^2)_{\text{grav.}} + N_{\text{Mess}}(\tilde{m}_0^{\text{GMSB}})^2, $$

$$ M_a = (M_a)_{\text{grav.}} + N_{\text{Mess}}(M_a^{\text{GMSB}}). $$

The negative contribution to $\tilde{m}_0^2$ induced by the UV divergence has strong consequences on the mass spectrum and the phenomenology of the model. Indeed, first of all, the spectrum is generically ”compressed”, since the values of the gaugino masses are increased whereas the
scalars ones decreased. Moreover, since the GMSB negative contributions are proportional to the SM charges of the scalars, the squarks are more sensitive than sleptons to them, whereas the gravitational contribution is universal, as shown above. The result of this interplay is shown in Table 1, where two different point in the space of the high-energy parameters are chosen (together with the value of the coupling $\lambda$ and the number of messengers $N_{\text{Mess}}$) and the comparison between the simple gravity mediated model and the complete one is shown. The low-energy mass spectrum is calculated using the Fortran package SUSPECT$^7$.

In addition, also the nature of the neutralino is considerably altered. Indeed, decreasing the value of $m_{T/3}^2$ and $m_{Q3}^2$ affects the RG equation for $m_{H2}^2$ and consequently one can have a smaller value for $\mu^2$. In this case the lightest neutralino is generally higgsino-like or a mixed bino-higgsino state and a good value for the relic abundance, compatible with WMAP bounds, is obtained, whereas this is not possible in the simple gravity mediated models, as computed using the routines provided by the program micrOMEGAs2.0$^8$.

Acknowledgments

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See also the web page http://wwwlapp.in2p3.fr/lapth/micromegas.
MODEL-DEPENDENT SEARCHES FOR NEW PHYSICS AT HERA

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Some of the searches for phenomena and particles beyond the Standard Model performed at HERA relying on specific theoretical models and using almost all the collected luminosity are discussed here. They particularly concern leptoquarks, lepton flavour violation, excited fermions, the anomalous top coupling and contact interactions, with improved limits on the quark radius.

1 Introduction

The HERA collider delivered luminosity for 15 years colliding $e^+ (e^-)$ with $p$ at a centre-of-mass energy $\sqrt{s} \simeq 300-320$ GeV. The data taking started in 1992 and continued until 2000 (HERA I phase). Then a substantial upgrade program involved both the machine and the experiments and the data taking was resumed in 2003 and continued until 2007 (HERA II phase), during which longitudinally polarised $e^\pm$ were available in most of data. The two general purpose experiments H1 and ZEUS ended data taking in summer 2007, after collecting a total integrated luminosity of about 1 $fb^{-1}$. At HERA an extensive program of searches for new particles and phenomena beyond the Standard Model (SM) has been carried out in a unique $ep$ environment. The focus in this paper will be on recent results of searches inspired or driven by specific theoretical models.

2 Leptoquarks

Starting from the symmetry between the quark and the lepton sectors many extensions of the SM predict bosons with fractional electromagnetic charge and both lepton and baryon numbers. A widely used model for leptoquarks is the phenomenological model of Buchmüller-Rückl-Wyler (BRW)\(^1\) which assumes invariance under $SU(3)_C \times SU(2)_L \times U(1)_Y$, conservation of the lepton number $L$ and the baryon number $B$ and a set of 7 scalar and 7 vector leptoquarks (4 decaying into both $e\bar{q}$ and $\nu q$) classified according to the fermion number $F = 3B + L = 0, 2$ and coupling to either left handed or right handed leptons, but not to both, with fixed branching ratio into $e\nu$ (1, 1/2), $\nu q$ (0, 1/2). At HERA, leptoquarks can be resonantly produced in the $s$ channel or exchanged in the $u$ channel between the incoming lepton and the quark from the proton. The resonant production shows up as a peak in the mass spectrum or an enhancement in $x$ distribution at the value corresponding to the mass $M$ of the leptoquark: $x = M^2/s$. As a consequence of quark densities in the proton, $e^- p$ and $e^+ p$ collisions offer respectively best sensitivities to $F = 2$ and $F = 0$ leptoquarks. The availability of polarisation of both signs within the HERA II sample has the advantage of enhancing the sensitivity to individual leptoquarks species. H1 searched for leptoquarks...
studying the inclusive Neutral Current and Charged Current Deep Inelastic Scattering high $Q^2$ $e^\pm p$ samples from HERA I and HERA II and using an integrated luminosity of 482 pb$^{-1}$.

No excess was seen in the $e-jet$, $\nu-jet$ mass spectra (fig. 1) and limits were set on the couplings and masses of the different leptoquark types $^2$ (fig. 2).

3 Lepton Flavour Violation

Leptoquarks can couple to different fermion generations and mediate lepton flavour violation processes in family non diagonal models.

H1 searched for $F = 2$ leptoquarks coupling to $eq$ and $\mu q$ using $e^\pm p$ HERA II data and an integrated luminosity of 158 pb$^{-1}$. No evidence for leptoquarks mediating lepton flavour violation was obtained and limits were set on couplings and masses of leptoquarks coupling to 1st and 2nd generation fermions (fig. 2). For an electromagnetic type coupling masses below 291-433 GeV can be excluded depending on the leptoquark type $^3$.

4 Excited leptons

To try to explain the hierarchy problem, models of compositeness introduce substructures to SM fermions, implying the existence of fermion excited states. Couplings between excited fermions and SM fermions can be described with phenomenological gauge mediated models $^4$. Excited fermion states have spin and isospin $1/2$ with both left-handed ($F^*_L$) and right-handed ($F^*_R$) components in weak iso-doublets. They can decay into fermions and gauge bosons. Magnetic type transitions between SM fermions $F$ and excited states $F^*$ can take place. Weight factors $f$, $f'$ and $f_\Sigma$ are used to set the coupling strength to the three gauge groups (U(1), SU(2) and SU(3)). The branching ratios of excited lepton decays can be fixed by assuming a specific relation between $f$ and $f'$ and then the production cross section depends only on $f/\Lambda$ where $\Lambda$ is the compositeness scale.
H1 searched for $e^+ \rightarrow e\gamma$, $e^+ \rightarrow eZ$ with $Z \rightarrow q\bar{q}$, and $e^+ \rightarrow \nu W$ with $W \rightarrow qq'$ using $e^\pm p$ data and an integrated luminosity of 475 pb$^{-1}$. No evidence for $e^+$ production was observed. Improved limits with respect to LEP and Tevatron were set\(^7\).

Due to the helicity structure of electroweak interactions and the valence quark densities in the proton, signals for excited neutrinos are expected to be stronger in $e^- p$ rather than in $e^+ p$ data. H1 searched for $\nu^+ \rightarrow \nu\gamma$, $\nu^+ \rightarrow \nu Z$ with $Z \rightarrow q\bar{q}$, and $\nu^+ \rightarrow eW$ with $W \rightarrow qq'$ using $e^- p$ data and an integrated luminosity of 184 pb$^{-1}$. No evidence was found and new limits were set\(^8\)(fig. ??). Masses were excluded in the range up to 213 GeV ($f = -f'$) and 196 GeV ($f = f'$). The H1 analysis has entered regions of masses not previously explored.

5 Anomalous top coupling

At HERA top quarks can only be singly produced. SM single-top production proceeds via the Charged Current reaction $ep \rightarrow \nu tbX$. As the SM cross section at HERA is less than 1 fb any observed single-top event must come from physics beyond the SM. In a Flavour Changing Neutral Current reaction the incoming lepton exchanges a $\gamma$ or $Z$ with an up-type quark in the proton, yielding a top quark in the final state most sensitive to a coupling of the type $tq\gamma$. The $u$-quark dominates at large $x$ and therefore the production of single top quark is related to the coupling $tu\gamma$. H1 searched for single top events in a sample of isolated leptons with high $p_t$ using $e^\pm p$ data and an integrated luminosity of 482 pb$^{-1}$. The analysis searched for anomalous production of $t$ decaying into $b$ and $W$ with subsequent decay of $W$ into an electron or a muon. A multivariate discrimination, based on a phase space density estimator with a range searching algorithm was used to separate the signal from the SM background (mostly real W production). The upper limit on the cross section set by H1\(^9\) is $\sigma_{ep \rightarrow t\ell X} < 0.16$ pb, leading to the most stringent limit to date on $k_{tu\gamma} < 0.14$ at 95% C.L. (fig.4).
Figure 3: H1 exclusion limits at 95% C.L. on the coupling $f/\Lambda$ as a function of the mass of the $e^*$ for gauge mediated interactions, with the assumption $f = +f'$ (left). H1 exclusion limits at 95% C.L. on the coupling $f/\Lambda$ as a function of the mass of the $\nu^*$ assuming $f = -f'$ (right).

6 Contact interactions and quark radius

Four-fermion contact interactions describe effects from processes at much higher scales, which could alter the SM distributions at high $Q^2$ and interfere with the predictions at intermediate $Q^2$. These effects modify the tree level amplitude $eq \rightarrow eq$. Let us focus on vector terms (as scalar and tensor terms are already constrained by previous searches). The Lagrangian can be written as:

$$L_{CI} = \sum_{\alpha,\beta = L,R}^{q=u,d} \eta_{\alpha\beta} (\bar{q}_\alpha \gamma^\mu q_\alpha)(\bar{q}_\beta \gamma_\mu q_\beta)$$  \hspace{1cm} (1)

The equation:

$$\eta_{\alpha\beta} = \frac{g_{CI}^2}{\Lambda^2}$$  \hspace{1cm} (2)

where $g_{CI} = 4\pi \epsilon = \pm 1$ defines the structure of the model.

Contact interaction effects could come from the exchange of extra gauge bosons ($Z'$), the production or exchange of leptoquarks or squarks, compositeness, gravitational effects (extra-dimensions) or from a finite quark radius. ZEUS analysed inclusive Neutral Current Deep Inelastic $e^+p$ data from HERA I and HERA II corresponding to an integrated luminosity of 330 pb$^{-1}$, comparing the data to SM predictions and performing a QCD fit where experimental and theoretical uncertainties are taken into account.\textsuperscript{10} Besides general model independent limits on contact interactions (values of the scale $\Lambda_{eqcl}$ in the range 2.0-8.0 TeV) depending on the chiral structure, limits were also set on the heavy leptoquark (beyond the available CM energy) couplings to the first generation ($M_{LQ}/\lambda$ in the range 0.29-2.08 TeV).

In some $4+n$ dimensional string theories\textsuperscript{11,12,13} compactified extra dimensions have size $R \simeq 1$ mm. The effective Planck scale $M_S$ related to the Planck scale $M_P \simeq 10^{19}$ GeV: $M_P^n = M_S^{2+n} R^n$ can be as small as 1 TeV. Graviton can propagate into the extra dimension, visible in the or-
ordinary 4 dimensions as a Kaluza-Klein tower of excited states with spacing $\Delta m = \frac{1}{R}$. Such states can be summed up to $M_S$, give sizeable effects, equivalent to a contact interaction term $\eta_Q \simeq \frac{2\lambda}{M_S^2}$ where $\lambda \simeq 1$. The interference with the SM can be constructive or destructive. Constraints were derived by ZEUS for such extra dimension scales: $M_S > 0.9$ TeV for $\lambda = -1$ and $M_S > 0.88$ TeV for $\lambda = +1$.

As far as the finite size of the quark is concerned in a classical approach to the quark substructure a charge distribution of radius $R_q$ in the quark can be described using a form factor:

$$\frac{d\sigma}{dQ^2} = \frac{d\sigma^{SM}}{dQ^2} \cdot (1 - \frac{R_q^2}{6} \cdot Q^2)^2 \quad (3)$$

This effect leads to a decrease of cross sections at high $Q^2$. An upper limit on quark radius was extracted from the ZEUS analysis: $R_q < 0.62 \cdot 10^{-16} \text{ cm}$. A study of high $Q^2$ Neutral Currents single differential cross section by H1 using the complete HERA I and HERA II data and an integrated luminosity of 270 pb$^{-1}$ ($e^+p$) and 165 pb$^{-1}$ ($e^-p$) led to a limit: $R_q < 0.74 \cdot 10^{-16} \text{ cm at 95 \% C.L.}$ (fig. 5).

7 Conclusions

The complete statistics of 15 years of data taking is being exploited by H1 and ZEUS to improve the sensitivity of the searches for new physics in the unique HERA environment. H1 and ZEUS at HERA have performed a number of model dependent searches finding no evidence for leptoquarks or lepton flavor violation, for excited electrons or excited neutrinos. Looking for single top production new limits on the anomalous top coupling are set. Limits on the contact interaction scales and quark radius have been updated fitting the Deep Inelastic Scattering...
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\[ R_q^2 = (0.62 \times 10^{-16} \text{cm}^2) \]

\[ R_q^2 = -(0.8 \times 10^{-16} \text{cm}^2) \]

\[ \frac{d^2 \sigma}{dQ^2} \]

\[ Q^2 (\text{GeV}^2) \]

\[ \text{H1 Preliminary} \]

Figure 5: Ratio of inclusive neutral current deep inelastic scattering data obtained by ZEUS in $e^\pm$ (left) and single differential cross sections obtained by H1 in $e^+p$ (top right) $e^-p$ (bottom right) to SM expectations as a function of $Q^2$, compared with 95% C.L. limits on the effective mean square radius of the electroweak charge of the quark.

differential cross sections at high $Q^2$. For some of these searches the two collaborations are going to provide a combination of H1 and ZEUS data.

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Theoretical issues and methods
Is $N = 8$ Supergravity an Ultraviolet Finite Quantum Field Theory?

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Advances in the computation of quantum amplitudes in supergravity theories raise the question whether maximal supergravity in $D = 4$ spacetime dimensions might actually be free of ultraviolet divergences. On the other hand, supersymmetric non-renormalization theorems give no indication of cancellations for anything beyond half-BPS counterterm operators. The jury is still out, and bets are being taken on the outcome.

Formulating an acceptable quantum theory of gravity remains the prime challenge to fundamental theoretical physics. A basic problem in formulating such a theory was already recognized in the earliest approaches to the problem in the 1930’s: the dimensional character of Newton’s constant gives rise to ultraviolet divergent quantum correction integrals. In the 1970’s, this was confirmed explicitly in the first Feynman diagram calculations of the radiative corrections to systems containing gravity plus matter. The time lag between the general perception of the divergence problem and its first concrete demonstration was due to the complexity of Feynman diagram calculations involving gravity. The necessary techniques were an outgrowth of the struggle to Lorentz-covariantly control the quantization of non-abelian Yang-Mills theories, the Standard Model of weak and electromagnetic interactions, and in quantum chromodynamics.
With the advent of supergravity in the mid 1970's, hopes rose that the specific combinations of quantum fields in supergravity theories might possibly tame the gravitational UV divergence problem. Indeed, it turns out that all irreducible supergravity theories in four-dimensional spacetime, i.e. theories in which all fields are irreducibly linked to gravity by supersymmetry transformations, have remarkable cancellations in Feynman diagrams with one or two internal loops.

There is a sequence of such irreducible (or “pure”) supergravity models, characterized by the number $N$ of local (i.e. spacetime-dependent) spinor parameters. In four-dimensional spacetime, minimal, or $N = 1$, supergravity thus has 4 supersymmetries corresponding to the components of a single Majorana spinor transformation parameter. The maximal possible supergravity in four dimensional spacetime has $N = 8$ spinor parameters, i.e. 32 independent supersymmetries.

The hopes for “miraculous” UV divergence cancellations in supergravity were subsequently dampened by the realization that the divergence-killing powers of supersymmetry most likely do not extend beyond the two-loop order for generic pure supergravity theories. The three-loop anticipated invariant is quartic in curvatures, and has a purely gravitational part given by the square of the Bel-Robinson tensor.

The flowering of superstring theory in the 1980’s and 1990’s, in which the UV divergence problems of gravity are cured by a completely different mechanism replacing the basic field-theory point-particle states by extended relativistic object states, pushed the UV divergence properties of supergravity out of the limelight, leaving the supergravity UV problem in an unclear state.

Nonetheless, among some researchers a faint hope persisted that at least the maximal $N = 8$ supergravity might have special UV properties. This hope was bolstered by the fact that the fact that the maximal supersymmetric Yang-Mills theory, which has $N = 4$, i.e. 16-component supersymmetry, is completely free of ultraviolet divergences in four-dimensional spacetime. This was the first interacting UV-finite theory in four spacetime dimensions.

It is this possibility of “miraculous” UV divergence cancellations in maximal supergravity that has now been confirmed in a remarkable 3-loop calculation by Z. Bern et al. Performing such calculations at high loop orders requires a departure from textbook Feynman-diagram methods, because the standard approaches can produce astronomical numbers of terms. Instead of following the standard propagator & vertex methods for the supergravity calculations, Bern et al. used another technique which goes back to Feynman: loop calculations can be performed using the unitarity properties of the quantum S-matrix. These involve cutting rules that reduce higher-loop diagrams to sums of products of leading-order “tree” diagrams without internal loops. This use of unitarity is an outgrowth of the optical theorem in quantum mechanics for the imaginary part of the S-matrix.

In order to obtain information about the real part of the S-matrix, an additional necessary element in the unitarity-based technique is the use of dimensional regularization to render UV divergent diagrams finite. In dimensional regularization, the dimensionality of spacetime is changed from 4 to $4 - \epsilon$, where $\epsilon$ is a small adjustable parameter. Traditional Feynman diagram calculations also often use dimensional regularization, but normally one just focuses on the leading $1/\epsilon$ poles in order to carry out a renormalization program. In the unitarity-based approach, all orders in $\epsilon$ need to be retained. This gives rise to logarithms in which real and imaginary contributions are related.

In the maximal $N = 8$ supergravity theory, the complexity of the quantum amplitudes factorizes, with details involving the various field types occurring on the external legs of an amplitude multiplying a much simpler set of scalar-field Feynman diagrams. It is to the latter that the unitarity-based methods may be applied. Earlier applications of the cutting-rule unitarity methods based on iterations of two-particle cuts gave an expectation that one might have cancellations for $D < 10/L + 2$, where $D$ is the spacetime dimension and $L$ is the number.
of Feynman diagram loops (for \( L > 1 \)). Already, this gave an expectation that \( D = 4 \) maximal supergravity would have cancellations of the UV divergences at the \( L = 3 \) and \( L = 4 \) loop orders. This would leave the next significant test at \( L = 5 \) loops. In the ordinary Feynman-diagram approach, a full calculation at this level would involve something like \( 10^{30} \) terms. Even using the unitarity-based methods, such a calculation would be a daunting, but perhaps not impossible, task.

The impressive new elements in the 3-loop calculation of Bern et al.® are the completeness of their calculation and the unexpected further patterns of cancellations found. This could suggest a possibility of unexpected UV cancellations at yet higher loop orders. Although the various 3-loop diagram classes were already individually expected to be finite on the basis of the earlier work by Bern et al., the new results show that the remaining finite amplitudes display additional cancellations, rendering them “superfinite”. In particular, the earlier work employed iterated 2-particle cuts and did not consider all diagram types. The new complete calculation displays further cancellations between diagrams that can be analyzed using iterated 2-particle cuts and the additional diagrams that cannot be treated in this way. The set of three-loop diagrams is shown in Figure 1. The end result is that the sum of all diagram types is more convergent by two powers of external momentum than might otherwise have been anticipated.

Does such a mechanism cascade in higher-order diagrams, rendering the maximal \( N=8 \) theory completely free of ultraviolet divergences? No one knows at present. Such a scenario might pose puzzling questions for the superstring program, where it has been assumed that ordinary supergravity theories need string ultraviolet completions in order to form consistent quantum theories. On the other hand, there are hints from superstring theory that precisely such an all-orders divergence cancellation might take place in the \( N = 8 \) theory. On the other hand, it is not clear exactly what one can learn from superstring theory about purely perturbative field-theory divergences.

One thing that seems clear is that ordinary Feynman diagram techniques coupled with the “non-renormalization” theorems of supersymmetry are unlikely to be able to explain finiteness properties of \( N = 8 \) supergravity at arbitrary loop order. Earlier expectations\(^1,5,6,7\) were that the first loop order at which divergences that cannot be removed by field redefinitions would be three loops in all pure \( D = 4 \) supergravities. A key element in this anticipation was the expectation that the maximal amount of supersymmetry that can be linearly realized in Feynman diagram calculations (aka “off-shell supersymmetry”) is half the full supersymmetry of the theory, or 16
out of 32 supercharges for the maximal $N = 8$ theory.

Similarly to the way in which chiral integrals of $N = 1$, $D = 4$ supersymmetry achieve invariance from integrals over less than the theory’s full superspace, provided the integrand satisfies a corresponding BPS type constraint, there are analogous invariants involving integration over varying portions of an extended supersymmetric theory’s full superspace. “Half-BPS” operators require integration over just half the full set of fermionic variables. And if half the full supersymmetry were the maximal amount that can be linearly realized (so giving strong results from the corresponding Ward identities), such operators would be the first to be allowed as UV counterterms.

The results of Ref. 9 show that the half-BPS expectation for the first allowed counterterms is too conservative in the case the maximal theory. But more recent advances in the understanding of supersymmetric non-renormalization theorems push the divergence onset boundary out slightly for the maximal theory, so that half-BPS counterterms that require superspace integrals over half the 32 component superspace are now expected to be the last disallowed counterterms instead of the first allowed ones. The resulting current expectations for first divergences from a traditional Feynman diagram plus non-renormalization viewpoint are shown for various spacetime dimensions in Table 1.

<table>
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<tr>
<th>Dimension $D$</th>
<th>11</th>
<th>10</th>
<th>8</th>
<th>7</th>
<th>6</th>
<th>5</th>
<th>4</th>
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<tr>
<td>Loop order $L$</td>
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<td>1</td>
<td>2</td>
<td>3</td>
<td>4</td>
<td>5</td>
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<tr>
<td>Gen. form</td>
<td>$\partial^{12} R^4$</td>
<td>$\partial^{10} R^4$</td>
<td>$\partial^8 R^4$</td>
<td>$\partial^6 R^4$</td>
<td>$\partial^6 R^4$</td>
<td>$\partial^4 R^4$</td>
<td></td>
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</tbody>
</table>

Table 1: Current maximal supergravity divergence expectations from Feynman rules and non-renormalization theorems.

The behavior of maximal $N = 4$ supersymmetric Yang-Mills theory in dimensions $D > 4$ may be a model for what is happening. Contrary to earlier expectations of UV divergences at the 4 loop order in $D = 5$ spacetime, the unitarity-based methods indicate that this SYM onset should be postponed to the 6-loop order. But here, the standard Feynman diagram methods have a comeback through the realization that the 4-loop finiteness could be explained using more sophisticated “harmonic superspace” methods.

There are two new recent elements to the non-renormalization theorem perspective. One is the realization that maximal SYM can be formulated in a “1/2 SUSY + 1” formalism which is not however Lorentz covariant. Such a SYM formulation dimensionally reduces to (8,1) supersymmetry in $D = 2$. Although considerations of gauge invariance implications in various dimensions are still ongoing, this formulation should be just the minimum needed to rule out the half-BPS operators. Moreover, there is an analogous “1/2 SUSY + 1” formulation for maximal supergravity dimensionally reduced to $D = 2$, having (16,1) supersymmetry. Providing this can be successfully lifted to a viable quantization formalism in $D = 4$, it should be just enough to rule out the $D = 4$ 3-loop candidate counterterm, now known from Ref. 10 not to occur.

The second new approach to the derivation of non-renormalization theorems is via “algebraic renormalization”, which uses BRST cohomological techniques and has been used to give yet another demonstration of the finiteness of $D = 4, N = 4$ SYM. Similar techniques for maximal supergravity are anticipated also to kill the eligibility of the 1/2 BPS $D = 4$ 3-loop candidate counterterm.

The overall picture that emerges from the non-renormalization theorems and the currently known divergence results from calculation is that the half-BPS operators are ruled out as UV counterterms, but that operators with less than half BPS character (thus requiring superspace integrals with more than half of the theory’s full supersymmetry) are not. The most accessible test of this proposition will occur at 4 loops in $D = 5$. As is not uncommon in this subject, bets are being taken on the outcome, the payoff to be made in bottles of wine.
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HOLOGRAPHIC TECHNIQUES FOR ASYMPTOTICALLY-FREE GAUGE THEORIES

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Novel techniques based on holographic ideas are tried on the prototype strongly coupled gauge theory: QCD. The ideas are developed and a well motivated phenomenological model (Improved Holographic QCD) is presented and compared to various non-perturbative regimes, both at zero and finite temperature.

1 Introduction

Strongly coupled gauge theories are omnipresent in theoretical physics, and have been forced experimentally upon us with the realization that the strong interactions are best described by an asymptotically free SU(3) gauge theory. The gauge coupling is weak at large energies and perturbation theory is applicable. However it is strong at low energy and almost all realistic observables contain parts that are sensitive to strongly coupled physics.

Beyond QCD, theorists have argued that strongly coupled gauge theories can play an important role in the physics beyond the standard model. We will mention here two such incarnations. The first concerns a strongly coupled gauge theory that is responsible for producing a composite Higgs that will break the electroweak symmetry at lower energies, $^1$. Such classes of theories come under the name of "technicolor" and although their popularity had its ups and downs, they are reanalyzed currently due to the use of novel non-perturbative holographic tools.

The second example concerns strong coupling dynamics that can trigger supersymmetry
breaking in a hidden sector. This supersymmetry breaking is expected to be transferred to the Supersymmetric SM sector either via universal interactions (gravity) or via gauge gauge interactions (gauge mediation). We should also mention that in theories beyond the SM, and especially in string theory vacua, strongly coupled hidden sectors are generic and if their associated scales are in the TeV region they might produce signals at LHC.

Various techniques have been developed to deal with the strong coupling problem of gauge theories. The most straightforward one, is numerical evaluation of the quantities of interest on a computer. This is the lattice approach that has been applied mostly to QCD, with considerable success. The lattice approach after 30 years is a mature discipline that has however its limitations, that basically translate into limitations of computing power. Despite the success of computational approaches, several interesting QCD observables remain out of reach, or cannot be computed to the required accuracy (examples are transport coefficients at finite temperature, relevant for recent heavy ion data from RHIC, or the physics at finite baryon number density and chemical potential).

A different theoretical approach was postulated by ’t Hooft in 1974, in order to generate a different perturbative expansion of strongly-coupled gauge theories. This is known as the large-N expansion where $N$ is the number of colors. Although it turned out that it was not possible to calculate even the leading approximation in this expansion for 4d gauge theories, several important properties were uncovered: (a) The perturbative expansion in powers of $1/N$ has the structure of a string theory with string coupling constant $g_s \sim 1/N$. (b) The gauge invariant QCD bound states, namely glueballs and mesons are non-interacting with $O(1)$ masses, to leading order. Their widths vanish as $1/N^2$ for glueballs and $1/N$ for mesons. (c) Baryons are heavy, with masses $N$, and behave as solitonic objects. Theorists have attempted to construct this string theory in four dimensions, but existing experience with string theories did not suggest optimism in this direction.

A new twist to the quest of the string theory underlying a strongly coupled gauge theory came with the realization that for such a gauge theory string should propagate in more than 4 dimensions. In a much more symmetric relative of QCD, namely $N=4$ superconformal SU(N) gauge theory, the dual string theory turned-out to be a type IIB string propagating in a ten-dimensional spacetime of the form $AdS_5 \times S^5$. In particular the fifth (radial) dimension of $AdS_5$ provided the holographic dimension that somehow captured the RG scale of the four-dimensional gauge theory that was defined on the $AdS_5$ boundary. This duality turned out to be a weak-strong coupling duality in the following sense: The gauge theory, that is exactly conformally invariant, has two dimensionless parameters: the number of colors $N$ that we take large, and the ’t Hooft coupling $\lambda \equiv g^2_{YM}N$ that we keep fixed in the large-$N$ limit. When $\lambda \ll 1$ one can use perturbation theory, and the relevant leading order diagrams are the planar diagrams. When $\lambda \gg 1$ perturbative techniques are of no use even as $N \to \infty$. On the other hand, the dual string theory is propagating on a manifold with curvature scale $1/\ell^2$, whose relation to the string length $\ell_s$ involves the ’t Hooft coupling: $\ell^2 = \sqrt{\lambda} \ell_s^2$. Moreover as the YM gauge coupling constant is given by $g^2_{YM} \sim g_s$, the string coupling constant in the large-$N$ limit is given by $g_s \sim \frac{1}{\sqrt{\lambda}}$ and for fixed $\lambda$ it is $O(1/N)$. Therefore, at large-$N$ and large ’t Hooft coupling, $\lambda \to \infty$, the theory is described by a string that moves on a weakly curved ten-dimensional manifold, and can therefore approximated by the dynamics of its zero modes: the strongly coupled large-N sYM theory is equivalent to type IIB supergravity on the $AdS_5 \times S^5$ background.

Since there has been a flurry of attempts to devise such correspondences for gauge theories with less supersymmetry with the obvious final goal: QCD. Several interesting string duals with a QCD-like low-lying spectrum and confining IR physics were proposed. Although such theories reproduced the qualitative features of IR QCD dynamics, they contain Kaluza-Klein modes, not expected in QCD, with KK masses of the same order as the dynamical scale of the gauge theory. Above this scale, the theories deviate from QCD.
A different and more phenomenological approach was in the meantime developed and is now known as AdS/QCD. The original idea was formulated in and it was successfully applied to the meson sector in. The bulk gravitational background consists of a slice of AdS$_5$, and a constant dilaton. There is a UV and an IR cutoff. Moreover, the confining IR physics is imposed by boundary conditions at the IR boundary. This approach, although crude, has been partly successful in studying meson physics, despite the fact that the dynamics driving chiral symmetry breaking must be imposed by hand via IR boundary conditions. Its shortcomings however include a glueball spectrum that does not fit well the lattice data, the fact that magnetic quarks are confined instead of screened, and asymptotic Regge trajectories for glueballs and mesons are quadratic instead of linear.

2 Improved Holographic QCD

In an improved holographic phenomenological model for QCD was proposed. It reunited inputs from both gauge theory and string theory while keeping the simplicity of a two-derivative action. It could describe both the region of asymptotic freedom as well as the strong IR dynamics of QCD.

The basic fields of the pure gauge theory (the closed string sector) that are non-trivial in the vacuum solution and describe the pure gauged dynamics, are the 5d metric $g_{\mu \nu}$ (dual to the YM stress tensor), a scalar $\Phi$ (the dilaton, dual to $Tr[F^2]$ ) that controls the ‘t Hooft coupling $\lambda_t$ of QCD, and an axion $a$, that is dual to the QCD instanton density $Tr[F \wedge F]$ and its source represents the $\theta$ angle. Quarks can be added to the pure gauge theory by adding $D^-D^-$ brane pairs in the background gauge theory solution. The $D^-D^-$ tachyon condensation then induces chiral symmetry breaking, $\alpha \geq 1$.

The action for the 5D Einstein-dilaton theory reads,

$$S_5 = M_p^3 N_c^2 \left( - \int d^5x \sqrt{g} \left[ R - \frac{4}{3} \lambda^2 + V(\lambda) \right] + 2 \int \partial_M d^4x \sqrt{h} \, K \right)$$

where $M_p$ is the Planck mass. The second term in the action is the Gibbons-Hawking with $K$ being the extrinsic curvature on the boundary.

The only nontrivial input in the two-derivative action of the graviton and the dilaton is the dilaton potential $V(\lambda)$, where $\lambda = e^\Phi$. $\lambda$ is proportional to the ‘t Hooft coupling of the gauge theory, $\lambda = \kappa \lambda_t$. The constant of proportionality $\kappa$ is treated as a parameter to be fitted to data. The potential is directly related to the gauge theory $\beta$-function once a holographic definition of energy is chosen. Although the shape of $V(\lambda)$ is not fixed without knowledge of the exact gauge theory $\beta$-function, its UV and IR asymptotics can be determined.

In the UV, the input comes from perturbative QCD. We demand asymptotic freedom with logarithmic running. This implies in particular that the asymptotic UV geometry is that of AdS$_5$ with logarithmic corrections. It requires a (weak-coupling) expansion of $V(\lambda)$ of the form

$$V(\lambda) = 12/\ell^2 (1 + v_1 \lambda + v_2 \lambda^2 + \cdots).$$

Demanding confinement of the color charges restricts the large-$\lambda$ asymptotics of $V(\lambda)$. In we focused on potentials such that, as $\lambda \to \infty$, $V(\lambda) \sim \lambda^\frac{4}{3} (\log \lambda)^{(\alpha-1)/\alpha}$ where $\alpha$ is a positive parameter. The IR asymptotics of the solution in the Einstein frame are:

$$ds_0^2 \to e^{-C(\bar{r})} \left( d\bar{r}^2 + d\bar{x}_4^2 \right), \quad \lambda_0 \to e^{3C/2(\bar{r})} \left( \frac{r}{\bar{r}} \right)^{\frac{4}{3}(\alpha-1)}$$

where the constant $C$ is related to $\Lambda_{QCD}$. Confinement requires $\alpha \geq 1$. The parameter $\alpha$ characterizes the large excitation asymptotics of the glueball spectrum, $m_n \sim n^{\frac{1}{\alpha - 1}}$. For linear confinement, we choose $\alpha = 2$. 145
The parameters of the holographic model a priori are: the Planck mass $M_p$, which governs the scale of interactions between the glueballs in the theory, the parameters $v_i$ that specify the shape of the potential, the scale $\Lambda$ that plays the role of $\Lambda_{QCD}$ and the AdS scale $\ell$. The latter is not a physical parameter but only a choice of scale: only $\Lambda \ell$ enters into the computation of physical observables. Before choosing a potential, $\kappa$ that relates $\lambda$ and the ’t Hooft coupling, is not a parameter as the physics is independent of $\kappa$. This is characteristic of the leading order in the large-$N_c$ expansion. Once a potential has been chosen then it is not the case anymore as $\kappa$ can be calculated by comparing for example to the perturbative QCD $\beta$-function. A specific choice for $V(\lambda)$ was made in\textsuperscript{9} with the appropriate asymptotic properties, that only depended on the parameter $\kappa$, hence fixing all $v_i$. Finally, $\kappa$ and $\Lambda$ are fixed by matching to the lattice data for the first two $0^{++}$ glueball masses. Once $\Lambda$ is fixed, all other interesting scales like the effective QCD string tension $\sigma$ are also fixed.

Glueball masses can be obtained by computing the spectrum of normalizable fluctuations of the metric and dilaton around the background solution. In table 1 we give an overview of the glueball spectrum calculated here and its comparison to the best existing lattice data both for $N = 3$ and $N \to \infty$. In figure 1 we give the almost linear trajectories of the $0^{++}$ and the $2^{++}$ states as computed from our model.

### 3 Finite temperature and deconfinement

We will now turn to the finite temperature dynamics in the pure gauge sector derived from the setup of\textsuperscript{9}. We find that this setup describes very well the basic features of large-$N_c$ Yang Mills at finite temperature. It exhibits a first order deconfining phase transition. The equation of state and speed of sound of the high temperature phase are remarkably similar to the corresponding lattice results. Moreover, using the zero temperature potential and without adding any extra parameter, we obtain a value for the critical temperature in very good agreement with the one computed from the lattice,\textsuperscript{10}.

**The deconfinement transition.** At finite temperature there exist two distinct types of solutions to the action (1) with AdS asymptotics:

i. The thermal graviton gas, obtained by compactifying the Euclidean time in the zero temperature solution with $\tau \sim \tau + 1/T$:

$$ds^2 = b_0^2(r) \left( dr^2 + d\tau^2 + dx^2 \right), \quad \lambda = \lambda_0(r).$$

This solution exists for all $T \geq 0$ and it corresponds to the confined phase, if the gauge theory at zero $T$ confines.

ii. The black hole (BH) solutions (in Euclidean time) of the form:

$$ds^2 = b^2(r) \left( \frac{dr^2}{f(r)} + f(r) dr^2 + dx^2 \right), \quad \lambda = \lambda(r). \quad (3)$$

with $f(0) = 1$. There exists a singularity in the interior at $r = \infty$ that is now hidden by a regular horizon at $r = r_h$ where $f$ vanishes. Such solutions correspond to a deconfined phase.

As we discuss below, in confining theories the BHs exist only above a certain minimum temperature, $T > T_{min}$.

The thermal gas as well as BH solution has two parameters: $T$ and $\Lambda$. Near the horizon, $f \to f_h(r_h - r)$ with $4\pi T = f_h$. From Einstein’s equations,\textsuperscript{10}:

$$4\pi T = b^{-3}(r_h) \left( \int_0^{r_h} \frac{du}{b(u)^3} \right)^{-1}. \quad (4)$$
In the large-$N_c$ limit, the physics is dominated by the saddle point with minimum free energy. For a given temperature we must therefore compare the free energies of solutions i. and ii.

We introduce a cutoff boundary at $r/\ell = \epsilon$ in order to regulate the infinite volume. The difference of the two scale factors is given near the boundary as

$$ b(\epsilon) - b_0(\epsilon) = C(T)\epsilon^3 + \cdots $$

(5)

By the standard rules of AdS/CFT we can relate $C(T)$ to the difference of VEVs of the gluon condensate: $C(T) \propto \langle \text{Tr} F^2 \rangle_T - \langle \text{Tr} F^2 \rangle_0$.

The free energy difference is given by

$$ \frac{\mathcal{F}}{M_p^3 N_c^2 V_3} = 12 \frac{C(T)}{\ell} - \pi T b^3(r_h) = 12 \frac{C(T)}{\ell} - \frac{T S}{4M_p^3 N_c^2 V_3}, $$

(6)

where, in the last equality, we used the fact that the entropy is given by the area of the horizon. It is clear that the existence of a non-trivial deconfinement phase transition is driven by a non-zero value for the thermal gluon condensate $C(T)$.

For a general potential we can prove the following (under mild assumptions):

**i.** There exists a phase transition at finite $T$, if and only if the zero-$T$ theory confines.

**ii.** This transition is of the first order for all of the confining geometries, with a single exception described in iii:

**iii.** In the limit confining geometry $b_0(r) \to \exp(-Cr)$ (as $r \to \infty$), the phase transition is of the second order and happens at $T = 3C/4\pi$.

**iv.** All of the non-confining geometries at zero $T$ are always in the black hole phase at finite $T$. They exhibit a second order phase transition at $T = 0^+$.

We illustrate the function $T(r_h)$ schematically in figure 2. It follows that in the confining geometries $\alpha > 1$, for a given $T > T_{\text{min}}$, there always exist a big and a small black hole solution. The big BH has positive specific heat hence it is thermodynamically stable, whereas the small BH is unstable. In the borderline confining geometry $\alpha = 1$, there is a single BH solution.

Existence of a $T_c \geq T_{\text{min}}$ follows from the physical requirement of positive entropy. From the first law of thermodynamics, it follows that $dF/dr_h = -S dT/dr_h$. Since $S > 0$ for any physical system, extrema of $F(r_h)$ coincide with the extrema of $T(r_h)$. Using also the fact that $F(r_h) \to -\infty$ for $r_h \to 0$ and $F(r_h) \to 0$ near $r_h \to \infty$, we arrive at conclusion (ii) described above: There is a first order transition for all of the confining geometries (This becomes second order for the borderline case $\alpha = 1$).

The small $r_h$ asymptotics also allows us to fix the value of the Planck mass in (1). This geometry corresponds to an ideal gas of gluons with a free energy density (We use lowercase letters for the densities of the corresponding functions) $f \to (\pi^2/45)N_c^2 T^4$. As the geometry becomes AdS, eq. (6) implies that: $f \to \pi^4(M_p\ell)^3N_c^2 T^4$. We conclude that $M_p\ell = (45\pi^2)^{-\frac{3}{4}}$. Using the value of $\ell$ in $^9$, we obtain $M_p \approx 2.3$ GeV.

4 Numerical Results at finite temperature

In $^9$ an explicit form of the scalar potential with the correct asymptotics was proposed. The resulting background, that corresponds to the choice $\alpha = 2$ in (2), exhibits asymptotic freedom, linear confinement, and a glueball spectrum in very good quantitative agreement with the lattice data. Here we present a numerical computation of the relevant thermodynamic quantities in this same theory. Our general analysis shows that this theory has black hole solutions above a temperature $T_{\text{min}}$ and exhibits a first order phase transition at some $T_c > T_{\text{min}}$.

To analyze the behavior of the theory at finite temperature, we have solved numerically Einstein’s equations for the metric and dilaton. The integration constants were fixed as explained
earlier. We find a minimum temperature for the existence of black hole solutions, \( T_{\text{min}} = 210 \text{ MeV} \).

Next, we compute the free energy difference between the black hole and thermal gas solutions, as a function of temperature.

The resulting free energy as a function of the temperature is shown in the left of figure 3, which clearly shows the existence of a minimum temperature, and a first order phase transition at \( T = T_c \), where \( \mathcal{F}(T_c) = 0 \). For \( T < T_c \), the thermal gas dominates, and the system is in the confined phase. For \( T > T_c \), the (large) black hole dominates, corresponding to a deconfined phase. The entire small black hole branch is always thermodynamically disfavored.

The value we obtain for the critical temperature, \( T_c = 235 \pm 15 \text{ MeV} \), is close to the value obtained for large-N Yang-Mills\(^{12} \), which with our normalization of the lightest glueball would be \( 260 \pm 11 \text{ MeV} \) (combining the results in\(^{12} \) and\(^{13} \)).

From the free energy we can determine all other quantities by thermodynamic identities:

\[
p = -\mathcal{F}/V_3, \quad s = 4\pi M_p^2 N_c^2 b_T^3(r_h), \quad \epsilon = p + Ts. \tag{7}
\]

Next, we present some of the thermodynamic quantities that are compared with the lattice results.

**Latent Heat.** The latent heat per unit volume is defined as the jump in the energy at the phase transition, \( L_h = T_c \Delta s(T_c) \), and it is expected to scale as \( N_c^2 \) in the large \( N_c \) limit\(^ {12} \). From eq. (7) we note that this expectation is reproduced in our theory. Quantitatively, we find \( L_h^{1/4}/T_c \approx 0.65\sqrt{N_c} \). This is to be compared with the value 0.77 reported in\(^ {12} \).

**Equation of state and the trace anomaly.** A useful indication about the thermodynamics of a system is given by the relations between the quantities \( \epsilon/T^4, 3(p/T^4), 3/4(s/T^3) \). In the right of figure 3 we compare our results for these quantities with the corresponding lattice results, reported in\(^ {14} \) (for \( N_c = 3 \)). We find good qualitative agreement. In the low temperature phase, the thermodynamic functions vanish to the leading order in \( N_c^2 \) and the jump in \( \epsilon \) and \( s \) at \( T_c \) reflects the first order phase transition. The fact that our curves lay below the lattice curves may be traced back to the relative smallness of the latent heat in our model.

The trace anomaly, \( (\epsilon - 3p)/T^4 \), is plotted in the left of figure 4, together with the lattice result from\(^ {14} \). From eq. (6), \( \epsilon - 3p \propto C(T) \), consistent with our interpretation of \( C(T) \) as the gluon condensate.

**Speed of sound.** This quantity is defined as \( c_s^2 = (\partial p/\partial \epsilon)_s = s/c_v \). It is expected to be small at the phase transition, and to reach the conformal value \( c_s^2 = 1/3 \) at high temperatures. In the right of figure 4 we compare our results with the lattice data, finding good agreement.

**Shear viscosity.** In agreement with the general results of\(^ {15} \), the ratio between shear viscosity and entropy density is \( \eta/s = (4\pi)^{-1} \).

**Acknowledgments**

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**References**

Table 1: Comparison between the glueball spectra in Ref. 1 and in our model. The states we use as input in our fit are marked in bold. The parenthesis in the lattice data indicate the percent accuracy.

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<th>This model (MeV)</th>
<th>Mismatch</th>
<th>$N_c \to \infty$</th>
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<td>4%</td>
<td>2153 (10%)</td>
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<td>$0^{++***}$</td>
<td>3990 (5%)</td>
<td>6%</td>
<td>4253</td>
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Figure 1: *Left:* Linear pattern in the spectrum for the first $40 \ 0^{++}$ glueball states. $M^2$ is shown units of $0.015 \ell^{-2}$. *Right:* The first $8 \ 0^{++}$ (squares) and the $2^{++}$ (triangles) glueballs. We used $b_0 = 4.2, \lambda_0 = 0.05$.

Figure 2: Schematic behavior of temperature as a function of $r_h$ (left) and the free energy density as a function of $r_h$ (right) for the infinite-$r$ geometries of the type (2), for different values of $\alpha$.

Figure 3: Left: Black hole free energy. Right: Dimensionless thermodynamic functions. The dashed curves correspond to the lattice data of reference [14].

Figure 4: *Left:* The trace anomaly. *Right:* The speed of sound. The dashed curves are the lattice result of reference [14].
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Electroweak Interactions and Unified Theories

III - Flavour Physics
1 Introduction

The Higgs sector is the only part of the Standard Model (SM) that has not been unveiled by experimental searches yet. What puzzles theorists is not only the lack of direct evidence of a Higgs boson so far, but also a theoretical prejudice against a light fundamental scalar particle. In fact, quantum corrections would like to push the Higgs mass and the electroweak scale near the cutoff, that can be as high as the Planck scale. Recent efforts, however, have focused on a less severe problem that has more impact on the LHC experiments: the little hierarchy problem. The stability of the electroweak scale would require the presence of new particles below or around a TeV, however precision electroweak measurements generically push such scale above 5-10 TeV. This bound is severely worsened if flavour is also taken into account: measurements in the Kaon and B systems push the scale of new physics up to $10^4$ TeV, thus requiring a fine tuning of several orders of magnitude unless a protection mechanism is summoned.

In the early ‘90, it has been realized that extra space dimensions are a rich playground for models of new physics. L. Randall and R. Sundrum proposed an interesting metric in 5 dimensions that may account for large hierarchies in a natural way: such metric can be written as

$$ds^2 = \left( \frac{R}{z} \right)^2 (dx_\mu dx^\mu - dz^2).$$

In this parametrization the 5D metric is explicitly invariant if we rescale the 4D coordinates $x_\mu$ and $z$ by the same amount: this means that moving along the coordinate $z$ is equivalent, from the 4D point of view, to a rescaling of lengths and energies. The space is compactified by placing two branes at the boundaries. The brane at small $z$ (UV-brane) will feel a large fundamental scale and therefore acts as...
Figure 1: Portrait of a generic model of EWSB in warped space.

a fundamental cutoff of the theory, while the brane at large \( z \) will feel a smaller scale which could be identified with the electroweak scale. In this setup a large hierarchy is rephrased in terms of order one parameters thanks to the exponential nature of the metric: this is more evident if one uses the coordinate \( z = R \exp \frac{\mu}{\Lambda} \).

This idea has sprouted many interesting models. Among them, one can identify the Higgs as the 5th polarization of a gauge boson from a broken bulk symmetry \(^2\): gauge invariance itself will protect the Higgs potential and solve the little hierarchy problem, given that the Kaluza-Klein (KK) resonances that cut off the loop divergencies are light enough. However, precision electroweak tests (PEWTs) require the heavy bosons to be above 2 TeV. Flavour physics plays an important role in these models: in fact one can use warped geometry to generate the hierarchies in the fermion mass spectrum naturally \(^3\). Once the fermions propagate in the bulk of the extra dimension, there will be more sources of flavour than in the SM: schematically the relevant terms in the action can be written as

\[
S = \int d^4x \int_{z_{UV}}^{z_{IR}} dz \left( \frac{R}{z} \right) \frac{4}{z} \left[ \bar{\psi}_{Q,u,d} \psi_{Q,u,d} + Y_{u,d} \bar{\psi}_Q H \psi_{u,d} \delta(z - z_{IR}) \right] + \ldots ,
\]

where the dots represent eventual UV localized terms and higher order operators. The SM fermion masses are generated by the interactions with the Higgs which is localized on or near the IR brane: for instance in gauge-Higgs unification models the delta function is replaced by the Higgs profile, peaked at large \( z \). The bulk masses \( c \), matrices in flavour space, are not real masses: they control the fermion localization along the extra dimension, and therefore the overlap with the Higgs. The wave functions are in fact exponentially sensitive to the \( c \)'s. Generically, this flavour-dependence of the wave functions will induce flavour non-universal couplings with the gauge KK modes, in particular the KK gluons, which will generate flavour changing neutral currents (FCNCs) at tree level. Flavour therefore may constrain the KK masses well above the TeV scale! Moreover, one needs to worry about new CP violating phases and higher order operators which may be suppressed by the IR scale.

If the bounds were as tight as in 4 dimensions, it would be the death of such models: however this is not the case. In order to understand this statement, we need to understand better the structure of a generic model of EWSB in warped geometry. The key is the localization of the wave functions: in fact it will determine both the spectrum via the boundary conditions, and the strength of their couplings via their overlap with other fields. Therefore, a generic model of EWSB can be portrayed in Fig. 1: the gauge boson wave functions are flat due to gauge invariance; the light fermions are localized towards the UV brane in order to suppress their coupling to the Higgs, or any other source of EWSB; on the other hand the top is necessarily localized toward the IR brane due to its heaviness. Finally the KK modes of all the bulk fields are localized towards the IR brane: as a generic consequence, they will couple more to the heavy SM particles than to the light ones. Assuming anarchic Yukawa couplings, the spectrum and mixings are both determined by the values of the fermion wave functions on the IR brane. The couplings of light fermions to the KK modes are small due to the localizations, and universal up to corrections of order \( O(m_f^2/m_{KK}^2) \): the light fermion are localized away from the IR brane, where KK wave functions are small and approximately constant. The flavour non-universal contribution comes from the values of the fermion wave functions on the IR brane, which are proportional to the fermion masses: this is the origin of the so-called Randall-Sundrum-GIM (Glashow Iliopoulos Maiani) mechanism. The situation is
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different for the top, which is localized on the IR brane. Therefore, all the FCNCs are induced by the third generation, and they are proportional to the mixing angles to the top. This mechanism allows to lower considerably the flavour bounds on KK masses from thousands to 8 TeV. However, bounds from flavour are still generally more severe than EWPTs, and reopens the little hierarchy problem and a fine tuning in the Higgs potential.

In recent years a lot of work has been dedicated to weaken this flavour bounds and push them below the EWPTs bounds\textsuperscript{5,6,7}. In the rest of the paper we will review the two mechanisms involving flavour symmetries in the bulk\textsuperscript{5} and minimal flavour violation in the bulk\textsuperscript{6}.

2 Flavour Symmetries in the Bulk

The easier way to avoid flavour bounds is to introduce flavour symmetries in the bulk. The simplest choice is to impose an SU(3)\textsubscript{L} \times SU(3)\textsubscript{R} in the bulk for both quarks and leptons, where we impose a single flavour symmetry SU(3)\textsubscript{R} for the right handed fermions due to an eventual custodial symmetry in the bulk\textsuperscript{8} which will contain them in the same multiplet. The symmetry will be broken to the diagonal SU(3)\textsubscript{D} on the IR brane by the Yukawa couplings: in this way Yukawas, bulk masses and bulk operators are all flavour diagonal. The SU(3)\textsubscript{R} is broken on the UV brane where localized kinetic operators for the right-handed fermions will generate both the mass hierarchies and the mixings: therefore the number of flavour matrices in this model is the same as in the SM and no extra CP violating phases appear. Also, the symmetries forbid FCNCs: one can use two SU(3) rotations in the up and down sector to diagonalize the kinetic operators. The neutral sector of the gauge bosons will remain flavour universal, while flavour violation will only appear in the interactions with charged gauge bosons like the W. Finally the only flavour violating higher order operators will be localized on the UV brane and will be suppressed by the large UV cutoff of the theory, therefore they can be safely neglected.

We can look more in detail to the main features of this scenario: the only flavour structure appears in the UV boundary conditions for the right-handed fields:

\[ f_R(m, z_{UV}) \bar{A}_{u,d} = m g_R(m, z_{UV}) K_{u,d} \bar{A}_{u,d}, \]

where \( f \) and \( g \) are generic flavour-blind wave functions, \( \bar{A} \) is the normalization - a vector in flavour space, and \( K_{u,d} \) are the UV kinetic matrices. One can diagonalize the kinetic matrices

\[ K_{u,d} = U_{u,d}^{\dagger} K_{u,d}^{\text{diag}} U_{u,d}, \]

so the spectrum will be determined by the eigenvalues \( k_i \) while the mixing matrices will fix the normalization coefficients \( \bar{A} \). Now, the couplings of neutral gauge bosons are diagonal, because they are proportional either to \( U \cdot U^{\dagger} = 1 \) or \( U \cdot K \cdot U^{\dagger} = K^{\text{diag}} \); on the other hand, the charged boson couplings will be proportional to \( U_{u,d}^{\dagger} U_{u,d} \). Therefore

\[ V_{CKM} = U_{u}^{\dagger} U_{d} + O(m_t^2) \]

where the corrections are due to the mass dependence of the wave functions, and all the flavour violating contributions will be proportional to the Cabibbo-Kobayashi-Maskawa matrix.

This model can be realized easily for the leptons, however it has problems when applied to quarks. The reason is that the top is very heavy and, due to the flavour symmetries, all the quarks share the same Yukawa coupling on the IR brane. The large Yukawa coupling will modify the fermion wave functions and generate universal corrections to the couplings. The flavour bound is therefore projected into EWPTs: the latter will push the KK masses above 10 TeV.

In order to solve this issue one needs to separate the top Yukawa from the light quarks. One can use different representations for the up and down right-handed quarks: using a singlet for the up quarks, including the top, can also help in lowering the bound from the coupling of the bottom with the Z boson\textsuperscript{9}. Moreover, one can impose a looser U(1)\textsuperscript{3} flavour symmetry for the right-handed up-type quarks and leave it unbroken. In this way the up type quarks Yukawas are all different:
The down sector is as before, therefore all the flavour mixing is induced in the down sector. One can show that in this model FCNCs are still forbidden, and the strongest bound on the KK masses is again the 2 TeV from precision measurements\(^5\).

3 5D MFV: 5 Dimensional Minimal Flavour Violation

Another interesting approach is to impose minimal flavour violation on the 5 dimensional model: flavour violating effects are not protected by a symmetry, but by the assumption that all the flavour structure can only be determined by the Yukawa matrices. In this case, the bulk masses in Eq. 2 are

\[
c_{u,d} \sim Y_{u,d}^+ + \ldots \quad c_Q \sim r Y_{u}^+ Y_u + Y_{d}^+ Y_d + \ldots
\]

The advantage of this approach is that one can still use different bulk masses to explain the hierarchies in the spectrum and the mixing angles, and at the same time gain a factor of \(\sim 3\) suppression in the flavour bounds that makes them again as low as the precision tests. Assuming anarchic Yukawa matrices is still enough to generate the required hierarchies due to the exponential sensitivity to the \(c\) parameters. Moreover, in the limit when \(c_Q\) only depends on one Yukawa, for instance when \(r \to 0\), one can diagonalize the down sector and eliminate all the flavour violating effects involving down type quarks. This means that the processes that violates flavour by 2 units, like for example the neural Kaon mixing which gives the strongest bounds, are suppressed by small \(r\). Therefore a moderately small \(r\) can provide the required factor of 3 in the bound without any flavour symmetry. Those small values are also preferred by the fit of the masses and mixing angles. Moreover the CP problem is also removed, because there isn’t any additional phase besides the SM one: for instance one can check that electric dipole moments only arise at two loops and they do not pose any additional bound\(^6\).

4 Conclusion and Outlook

Flavour physics is an important component of models in warped extra dimension. In fact, flavour bounds generically apply to the KK masses of the gauge bosons which play an important role in the electroweak symmetry breaking sector and are required to be at or around a TeV in order for the model to be natural. The bounds are much lower than in a generic 4 dimensional model due to a Randall-Sundrum GIM mechanism, however they are still one order of magnitude tighter than bounds from precision electroweak tests. Moreover, the warped geometry offer the possibility to construct an elegant model of flavour where both the hierarchies in the masses and in the mixing angles are explained in terms of order one parameters. If we were not concerned by the two orders of magnitude still separating the scale of new physics and the electroweak scale, this would be one of the most appealing models of flavour.

However, trying to lower the bounds from flavour has inspired a dense activity in recent years. We reviewed two nice ideas. One involving the use of bulk flavour symmetry, and one proposal of a minimal flavour violation paradigm. In the former case, one can eliminate all the flavour changing neutral currents at the price of giving up the nice explanation of the hierarchies. The heaviness of the top quark still requires some massaging as the light quarks cannot share its large Yukawa, however it is still possible to construct models with a relaxed flavour symmetry where the flavour bound is as low as 2 TeV.

In the case of minimal flavour violation, no symmetry is needed and a relation between the Yukawa matrices and all the other sources of flavour violation is enough to solve the CP problem and to parametrically suppress the most dangerous flavour violating effects. It is important to notice that the required suppression is just a factor of 3, and that this suppression is also preferred by the fit of the fermion masses and mixing angles.

The precise bound from flavour physics is therefore very important as it can have severe consequences on the phenomenology and viability of such models. It can easily push the new physics above the reach of the LHC and the electroweak sector of the model un-natural. There cannot be a viable model unless its flavour structure is studied in detail.

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\( \Delta F = 1 \) CONSIDERATIONS ON MINIMAL FLAVOR VIOLATION

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We present an updated phenomenological analysis of the minimal flavor violating (MFV) effective theory, both at small and large tan \( \beta \), in the sector of \( \Delta F = 1 \) processes. We evaluate the bounds on the scale of new physics derived from recent measurements (in particular from \( B \to X_s \gamma \), \( B \to X_s \ell^+ \ell^- \), \( B_s \to \mu^+ \mu^- \) and \( K \to \pi \nu \bar{\nu} \)) and we use such bounds to derive a series of model-independent predictions within MFV for future experimental searches in the flavor sector.

1 Introduction

The Standard Model (SM) accurately describes high energy physical phenomena up to the electro-weak (EW) scale \( \mu_W \sim 100 \text{ GeV} \). It is however known to be incomplete due to the lack of description of gravity, proper unification of forces as well as neutrino masses. In view of these shortcomings, it can be regarded as a low-energy effective description of physics below a UV cut-off scale \( \Lambda \). But if it is an effective theory, at what scale \( \Lambda \) below the unification or the Planck scale does it break down? The only dimensionful parameter in the renormalizable part of the Lagrangian is the Higgs mass, which is known to be quadratically sensitive to the cut-off scale of the theory. Then the EW hierarchy problem suggests that new physics (NP) should appear around or below \( \Lambda \lesssim 1 \text{ TeV} \). The non-renormalizable higher dimensional terms, formally suppressed by the increasing powers of the cut-off scale on the other hand mediate flavor changing neutral currents (FCNCs), may contain additional sources of CP violation and can violate baryon and lepton numbers. Even in absence of the later, precision measurements of low energy experiments put severe constraints on the scale of flavor and CP violating NP. Excellent agreement between SM predictions and experiment on \( \epsilon_K \) (constraining \( s \to d \) sector), \( A_{CP}(B_d \to \Psi K_s) \) and \( \Delta m_d \) (in the \( b \to d \) sector) and \( B \to X_s \gamma \) (for \( b \to s \) transitions) constrains
a general flavor violating NP to appear above $\Lambda \gtrsim 2 \times 10^5$ TeV, $2 \times 10^3$ TeV and 40 TeV respectively. The resulting tension between the two estimates of the NP scale illustrates what is often called the new physics flavor problem.

The Minimal Flavor Violation (MFV) hypothesis\textsuperscript{1,2} aims to solve the issue by demanding that all flavor symmetry breaking in and also beyond the SM is proportional to the SM Yukawas. A few direct consequences follow from this assumption: Firstly the Cabibbo-Kobayashi-Maskawa (CKM) matrix is the only source of flavor mixing and CP violation even beyond the SM. Thus, all (non-helicity suppressed) tree level and CP violating processes are constrained to their SM values. Finally, CKM unitarity is maintained and a (universal) unitarity triangle (UUT) can be determined from a constrained set of observables\textsuperscript{3}. Other details of phenomenology depend on the form of the EW Higgs sector of the theory. In case of a SM-like single Higgs doublet, the FCNCs in the down quark sector are all driven by the large top Yukawa ($\lambda_t$). At the same time, when performing the operator product expansion (OPE) at the EW scale, the SM basis of operators contributing to the effective weak Hamiltonian is complete also in presence of NP, making the MFV effective theory approach predictive. The same holds true at low $\tan \beta \equiv v_u/v_d$ if the Higgs sector is described in terms of an effective two Higgs doublet model of type II with the vacuum expectation values of the Higgses coupling to up(down) quarks denoted by $v_u(d)$. However, bottom Yukawa ($\lambda_b$) contributions become important at large $\tan \beta$ as $\lambda_b(\sim m_b \tan \beta/v_u) \sim \lambda_t$. Accompanied by the partial lifting of helicity suppression in the down sector, contributions due to new density operators have to be taken into account in the effective weak Hamiltonian. Still, the predictivity of the MFV effective theory approach is maintained by the small number of additional operators which need to be considered.

The symmetry principles underlying the MFV hypothesis establish solid links among different flavor observables at low energy and allow to probe and constrain the scale of MFV NP. Since (non-helicity suppressed) charged current interactions are not affected, bounds can be derived from $\Delta F = 2$ and $\Delta F = 1$ FCNC phenomenology. The $\Delta F = 2$ processes are box loop mediated in the SM, and only a few operators contribute to the effective weak Hamiltonian. The main observables here are the $K, B_q$ oscillation parameters to which MFV NP at low $\tan \beta$ contributes universally\textsuperscript{2}. A recent analysis\textsuperscript{4} was able to constrain this contribution and put a lower bound on the effective NP scale $\Lambda > 5.5$ TeV at 95% probability. The $\lambda_b \tan \beta$ contributions break the universality among kaon and $B$ meson sectors at large $\tan \beta$, resulting in a slightly weaker bounds of $\Lambda > 5.1$ TeV. New operators due to Higgs exchange in the loop start contributing only at very large values of $\tan \beta$, resulting in a bound on a certain combination of charged Higgs parameters. $\Delta F = 1$ processes on the other hand are penguin loop mediated in the SM, with many operators contributing. In concrete MFV models, they are often related to the $\Delta F = 2$ as well as flavour conserving phenomenology\textsuperscript{5}. On the other hand in our effective theory bottom-up approach they have to be considered completely orthogonal. An analysis of bounds coming from radiative, and (semi)leptonic decays of $K$ and $B$ mesons was performed a while ago\textsuperscript{2}, however limited experimental information at the time barred from exploring in particular the interesting role of the large $\tan \beta$ scenario. In the meantime, the situation has drastically improved and the new updated experimental and theoretical results on $\Delta F = 1$ FCNC mediated processes further motivate the revisiting and updating of this analysis. In the following we present a selection of results from such a study, the details of which will be presented elsewhere\textsuperscript{6}.

\section{Updating Analysis of $\Delta F = 1$ Constraints}

In the SM the effective weak Hamiltonian describing $\Delta F = 1$ FCNC processes among down-type quark flavors $q_i - q_j$ can be written as\textsuperscript{2}

\begin{equation}
\mathcal{H}^{\Delta F = 1}_{\text{eff}} = \frac{G_F \alpha_{\text{em}}}{2\sqrt{2} \pi \sin^2 \theta_W} V_{ti}^* V_{tj} \sum_n C_n Q_n + \text{h.c.},
\end{equation}

\begin{align}
\mathcal{H}^{\Delta F = 1}_{\text{eff}} = & \frac{G_F \alpha_{\text{em}}}{2\sqrt{2} \pi \sin^2 \theta_W} V_{ti}^* V_{tj} \sum_n C_n Q_n + \text{h.c.},
\end{align}
where $G_F$ is the Fermi constant, $\alpha_{em}$ is the fine structure constant, $\theta_W$ is the Weinberg angle and $V_{ij}$ are the CKM matrix elements. The short distance SM contributions are encoded in the Wilson coefficients $C_n$, computed via perturbative matching procedure at the EW scale. MFV NP manifests itself in the shifts of the individual Wilson coefficients in respect to the SM values $C_n(\mu_W) = C_n^{SM} + \delta C_n$. These shifts can be translated in terms of the tested NP energy scale $\Lambda$ as $\delta C_n = 2a\Lambda^2/\Lambda^2$, where $\Lambda_0 = \lambda_i \sin^2(\theta_W)m_W/\alpha_{em} \sim 2.4$ TeV is the corresponding typical SM effective energy scale. The value of the free variable $a$ depends on the details of a particular MFV NP model. In general $a \sim 1$ for tree level NP contributions, while $a \sim 1/16\pi^2$ for loop suppressed NP contributions. In our numerical results we put $a$ to unity.

In order to address low energy phenomenology, one needs to evaluate the appropriate matrix elements of the corresponding effective dimension 6 operators $Q_n$. At low tan $\beta$ we consider the EM and QCD dipole operators

$$Q_{7\gamma} = \frac{2}{g^2} m_j \bar{d}_{iL} \gamma_{\mu} d_{jR} (e_{\mu\nu}) , \quad Q_{8G} = \frac{2}{g^2} m_j \bar{d}_{iL} \gamma_{\mu} T^a_{\nu} d_{jR} (g_s G^a_{\mu\nu}) ,$$

where $g$ is the EW $SU(2)_L$ coupling, $e$ is the EM coupling, $g_s$ is the QCD coupling, $T^a$ are the $SU(3)_c$ generator matrices, while $F^\mu_{\nu}$ and $G^a_{\mu\nu}$ are the EM and QCD field tensors. They contribute to $B \to X_s \gamma$ decay as well as to the $B \to X_s \ell^+ \ell^-$ phenomenology, where in addition we get contributions from the EW-penguin operators

$$Q_{9V} = 2 \bar{d}_{iL} \gamma_{\mu} d_{jL} \ell_{\mu} \ell , \quad Q_{10A} = 2 \bar{d}_{iL} \gamma_{\mu} d_{jL} \ell_{\mu} \gamma_{5} \ell .$$

Here $\ell = e, \mu, \tau$ denotes the charged leptons. $Q_{10A}$ also mediates $B_q \to \ell^+ \ell^-$. Finally the Z-penguin operator

$$Q_{\nu\bar{\nu}} = 4 \bar{d}_{iL} \gamma_{\mu} d_{jL} \ell_{\mu} \gamma_{\nu}^\nu L$$

enters solely in $B \to X_s \nu\bar{\nu}$ and $K \to \pi \nu\bar{\nu}$ decays and can thus be constrained independently of the others. We do not consider NP contributions to QCD penguin operators as their impact on phenomenology is subdominant compared to long distance effects. At large tan $\beta$, one needs to take into account an additional density operator

$$Q_{S-P} = 4 \bar{d}_{iL} \gamma_{\mu} d_{jL} (\bar{\ell}_{\mu} \ell)$$

contributing to $B \to X_s \ell^+ \ell^-$ and $B_q \to \ell^+ \ell^-$. On the other hand, contributions from additional four quark density operators $^7^a$ which are also $\tan \beta$ enhanced and enter $B \to X_s \gamma$ and $B \to X_s \ell^+ \ell^-$ through one loop mixing with $Q_{7\gamma,8G}$ are $\alpha_{em}/4\pi \sim 0.001$ suppressed relative to those of $Q_{S-P}$ and thus turn out to be negligible after imposing the bounds on $Q_{S-P}$.

In our analysis we consider the most theoretically clean observables in order to derive reliable bounds on possible NP contributions. In particular, we use the inclusive branching ratio of the radiative $B \to X_s \gamma$ decay, measured with a lower cut on the photon energy. The latest HFAG value averaged over different measurements$^{10}$ is $Br(B \to X_s \gamma)_{E_{\gamma}>1.6 \text{ GeV}}^{\text{exp}} = 3.52(23)(9) \times 10^{-4}$, where the first error is statistical and the second systematic. Theoretically, the SM value is known to better than 8% and the expansion in terms of $\delta C_n$ evaluated at the weak scale is$^8$

$$Br(B \to X_s \gamma)_{E_{\gamma}>1.6 \text{ GeV}}^{\text{th}} = 3.16(23) \left( 1 - 2.28\delta C_{7\gamma} - 0.71\delta C_{8G} ight. \\
+ 1.51\delta C_{7\gamma}^2 + 0.78\delta C_{8G}\delta C_{7\gamma} + 0.25\delta C_{8G}^2 \right) \times 10^{-4} ,$$

where the central value and its error have been adjusted to take into account the CKM matrix element determination from the UUT analysis$^4$. Since $\delta C_7$ and $\delta C_8$ in absence of four quark density operator contributions enter in the same fixed combination to all relevant observables (any differences being artifacts of the truncated perturbative expansion) one can always eliminate

$^a$We thank Ulrich Haisch for pointing out these potential contributions.
one of them (e.g. $\delta C_{3G}$) from the analysis and then reconstruct the bound on both from the quadratic combination in eq (6).

A completely different combination of operators contributes to the helicity suppressed decay $B_s \rightarrow \mu^+\mu^-$. Experimentally the best upper bound on the branching ratio was recently put by the CDF collaboration $^{9}$ $Br(B_s \rightarrow \mu^+\mu^-)^{\exp} < 4.7 \times 10^{-8}$ at 90% C.L., which is only an order of magnitude above the SM prediction. The theoretical error of which is around 23% and is dominated by the lattice QCD determination of the $B_s$ decay constant. Again using UUT CKM inputs, the expansion in terms of $\delta C_i$ reads

$$Br(B_s \rightarrow \mu^+\mu^-)^{\exp} = 3.8(9) \left( 1 - 2.1\delta C_{10A} - 2.3\delta C_{S-P}$
$$+ 1.1\delta C_{10A}^2 + 2.4\delta C_{S-P}\delta C_{10A} + 2.7\delta C_{S-P}^2 \right) \times 10^{-9}. \quad (7)$$

Analysis of $B \rightarrow X_s\ell^+\ell^-$ is more involved since, not only do almost all of the above mentioned operators $(\mathcal{Q}_7,\mathcal{Q}_8,\mathcal{Q}_{9V},\mathcal{Q}_{10A},\mathcal{Q}_{S-P})$ contribute here, experimentally there are already a number of inclusive as well as exclusive measurements available, constraining different combinations of NP parameters. On the inclusive side, only the branching ratio $Br(B \rightarrow X_s\ell^+\ell^-)$, where $\ell = e, \mu$ is measured by the B factories $^{11}$ in several bins of di-lepton invariant mass squared ($q^2$). The errors vary from almost 90% in the first bin where only Belle has obtained a relevant signal, to around 30% in the other bins. The latest calculations estimate the theoretical error at around 7% for the bins below the charmonium region and around 10% for the high $q^2$ bin $^{12}$. The relevant formulae including NP contributions are rather lengthy and can be found in ref. $^{6,12}$.

Much more experimental information is available for exclusive channels where the $B \rightarrow K^{(*)}\ell^+\ell^-$ branching ratios as well as several angular distributions have already been measured $^{13}$. Theoretically however, despite considerable theoretical progress on the evaluation of the non-perturbative matrix elements of $\mathcal{Q}_n$ entering exclusive channels in the recent years $^{14}$, a reliable determination can only be expected from fundamentally non-perturbative methods, such as lattice QCD. In the meantime, any phenomenological implications based on existing form factor estimates should be treated with care. We will present an analysis of the impact of the exclusive modes on the MFV NP bounds elsewhere $^6$.

Finally MFV NP contributions to the Z-penguin operators can be constrained using the first experimental hints $^{15}$ of the $K^+ \rightarrow \pi^+\nu\bar{\nu}$ decay $Br(K^+ \rightarrow \pi^+\nu\bar{\nu})^{\exp} = 147(120) \times 10^{-12}$ and comparing them to the theoretical predictions, which are brought under control by the use of experimental data on $K\ell\nu$ decays $^{16}$ resulting in only 11% theoretical error. In presence of MFV NP the corresponding expression reads

$$Br(K^+ \rightarrow \pi^+\nu\bar{\nu}(\gamma))^{\th} = 7.53(82)(1 + 0.93\delta C_{\nu\bar{\nu}} + 0.22\delta C_{\nu\bar{\nu}}^2) \times 10^{-11}. \quad (8)$$

Common parametric inputs in our analysis are the particle masses and lifetimes from PDG $^{17}$ as well as the parameters of the CKM matrix, which, as already mentioned, we take from the UUT analysis $^4$. We perform a correlated fit of subsets of observables turning on NP contributions and extract probability bounds on the shifts of the Wilson coefficients away from their SM values.

3 Results

The compilation of bounds on the MFV NP scale in respect to all the probed operators is summarized in table 1. We present two sets of bounds. In the conservative estimate we take into account all the possible fine-tunings and cancellations among the various operator contributions, including discrete ambiguities in cases where the NP contributions might flip the sign of the SM pieces. For the second, more natural bounds, we consider each $\delta C_n$ individually and also discard flipped-sign fine-tuned solutions. The strongest bounds come naturally from the $B \rightarrow X_s\gamma$ decay rate and affect $\mathcal{Q}_{7,\gamma,8G}$. As can be seen, the effect of the discrete ambiguity is large and
Table 1: Summary of bounds on the MFV NP scales related to the probed effective operators. All the numerical values are the lower bounds at 95% probability on the MFV NP scale $\Lambda$ as explained in the text.

<table>
<thead>
<tr>
<th>Operator</th>
<th>Conservative bound [TeV]</th>
<th>Natural bound [TeV]</th>
</tr>
</thead>
<tbody>
<tr>
<td>$Q_{7\gamma}$</td>
<td>1.6</td>
<td>5.3</td>
</tr>
<tr>
<td>$Q_{8G}$</td>
<td>1.2</td>
<td>3.1</td>
</tr>
<tr>
<td>$Q_{9V}$</td>
<td>1.4</td>
<td>1.6</td>
</tr>
<tr>
<td>$Q_{10A}$</td>
<td>1.5</td>
<td>1.5</td>
</tr>
<tr>
<td>$Q_{S-P}$</td>
<td>1.2</td>
<td>/</td>
</tr>
<tr>
<td>$Q_{S-P}$</td>
<td>1.5</td>
<td>/</td>
</tr>
</tbody>
</table>

Figure 1: Correlation plots showing the most pronounced correlations among the bounds on the various NP Wilson coefficient shifts. The 68% (95%) probability regions are shown in green (red).

only the natural bounds on $\Lambda > 5.2(3.1)$ for $Q_{7\gamma}(8G)$ are competitive with the ones on $\Delta F = 2$ operators. The discrete ambiguity (also seen on utmost left plot in figure 1) could however be completely removed in the future once the experimental situation concerning the lowest $q^2$ region in $B \to X_s \ell^+\ell^-$ rate and especially the forward-backward asymmetry (FBA) improves. As expected, $Q_{S-P,\nu\bar{\nu}}$ operators are mainly bounded from single observables ($B_s \to \mu^+\mu^-$ and $K^+ \to \pi^+\nu\bar{\nu}$ respectively) leading to robust bounds around 1.2 TeV and 1.5 TeV respectively. Finally $\delta C_{9V,10A}$ are mainly bounded by $B \to X_s \ell^+\ell^-$ and using only presently available inclusive information the bounds are around 1.5 TeV. In all of the considered observables except $B \to X_s\gamma$ the experimental uncertainties strongly dominate and at present do not allow to discern discrete ambiguities or strong correlations as can be also deduced from figure 1 showing the most interesting pairwise correlation plots of the 68% and 95% allowed parameter regions.

4 Discussion and Outlook

In summary, immense experimental and theoretical progress in the area of flavor physics in the last decade has made it possible to constrain in a model independent way the complete set of possible beyond SM contributions to $\Delta F = 1$ and $\Delta F = 2$ processes due to possible MFV NP both at small and large $\tan \beta$. Bounds coming from $\Delta F = 2$ phenomenology are already very constraining, pushing the effective MFV NP scale beyond 5 TeV. In $\Delta F = 1$ sector, at present only the bounds coming from $B \to X_s\gamma$ are of comparable strength. However most uncertainties are dominated by experiments and one can look forward for the results of full dataset analyses by the B factories.

Using the derived bounds on the MFV NP contributions in $\Delta F = 1$ processes we are able to
make predictions for other potentially interesting observables to be probed at LHCb or a future Super Flavor Factory. As already mentioned, angular distributions like the FBA probe different combinations of the operators and would provide complimentary bounds. At the moment, considering bounds from inclusive measurements alone, no firm constraints on the FBA or its zero can be imposed within MFV models. This conclusion reinforces the importance of these observables and their potentiality of discovering relevant deviations.

Another set of observables displays interesting sensitivity to the tan $\beta$ enhanced $C_{S-P}$ contributions. Such are lepton flavor universality ratios $\Gamma(B \to K(e^+e^-))/\Gamma(B \to K(\mu^+\mu^-))$, which are very close to 1 with the SM as well as MFV models with low tan $\beta$. However even at tan $\beta$ present constraints already disallow deviations larger then 10% from unity for such ratios.

Finally the derived bounds allow to construct tests able to potentially rule out of MFV. Besides the interesting CP violation signals already emerging in the $B_s$ sector, in $\Delta F = 1$ sector first there are the firm relations among the different flavor transitions $[|V_{tb}V_{ts}/|V_{tb}V_{td}|/|V_{ts}V_{td}|]$ which might be probed with $K \to \pi\ell^+\ell^-$, $B \to X_\nu\bar{\nu}$ or $B_d \to \mu^+\mu^-$ processes. Also interesting in this respect is the FBA in $B \to K\ell^+\ell^-$ which is already restricted to be below 1% within MFV models regardless of tan $\beta$.

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References

Flavour Permutation Symmetry and Fermion Mixing

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We discuss our recently proposed $S_3 \downarrow \times S_3 \uparrow$ flavour-permutation-symmetric mixing observables, giving expressions for them in terms of (moduli-squared) of the mixing matrix elements. We outline their successful use in providing flavour-symmetric descriptions of (non-flavour-symmetric) lepton mixing schemes. We develop our partially unified flavour-symmetric description of both quark and lepton mixings, providing testable predictions for $CP$-violating phases in both $B$ decays and neutrino oscillations.

1 Introduction

Flavour observables, namely quark and lepton masses and mixings are neither predicted nor predictable in the Standard Model. Neither are they correlated with each other in any way. However, their experimentally determined values display striking structure: viewed on a logarithmic scale, the fermion masses of any given non-zero charge are approximately equi-spaced; the spectrum of quark mixing angles is described by the Wolfenstein form,\textsuperscript{1} suggestive of correlations between mixing angles and quark masses, and the lepton mixing matrix is well-approximated by the tri-bimaximal form.\textsuperscript{2} These striking patterns are the modern-day equivalents of the regularities observed around a century ago in hydrogen emission spectra, which were mathematically well-described by the Rydberg formula, but nevertheless had no theoretical basis before the advent of quantum mechanics. While consistent with the Standard Model, they lie completely outside its predictive scope, and are surely evidence for some new physics beyond it.

\textsuperscript{a}Talk given at the 43rd Rencontres de Moriond, La Thuile, Italy, March 2008.

\textsuperscript{b}Speaker.
In this talk, we report on our recent attempts\(^3\) to find a new description of fermion mixing which builds on the Standard Model and allows constraints on the mixing observables which make no reference to individual flavours, while describing mixing structures which are manifestly not flavour-symmetric, as observed experimentally. This approach does not in itself constitute a complete theory of flavour mixing beyond the Standard Model, but we hope that it might help stimulate new developments in that direction.

2 The Jarlskogian and Plaquette Invariance

Jarlskog’s celebrated \(CP\)-violating invariant,\(^4\) \(J\), is important in the phenomenology of both quarks and leptons. As well as parameterising the violation of a specific symmetry, it has two other properties which set it apart from most other mixing observables. First, its value (up to its sign) is independent of any flavour labels;\(^5\) Mixing observables are in general dependent on flavour labels, eg. the moduli-squared of mixing matrix elements, \(|U_{\alpha i}|^2\), certainly depend on \(\alpha\) and \(i\). Indeed, \(J\) itself is often calculated in terms of a subset of four mixing matrix elements, namely those forming a given plaquette\(^5\) (whose elements are defined by deleting the \(\gamma\)-row and the \(k\)-column\(^6\) to leave a rectangle of four elements):

\[
J = \text{Im}(\Pi_{\gamma k}) = \text{Im}(U_{\alpha i}U_{\alpha j}^* U_{\beta i} U_{\beta j}).
\]

However, it is well-known\(^4\) that the value of \(J\) does not depend on the choice of plaquette (ie. on its flavour labels, \(\gamma \) and \(k\) above) - it is “plaquette-invariant”. This special feature originates in the fact that \(J\) is \textit{flavour-symmetric}, carrying information sampled evenly across the whole mixing matrix. We recently pointed-out\(^3\) that in fact, \textit{any} observable function of the mixing matrix elements, flavour-symmetrised (eg. by summing over both rows and columns), and written in terms of the elements of a single plaquette (eg. using unitarity constraints), will be similarly plaquette-invariant. Both its expression in terms of mixing matrix elements, as well as its value, will be independent of the particular choice of plaquette.

The second exceptional property of \(J\) is that it may be particularly simply related to the fermion mass (or Yukawa) matrices:

\[
J = -i \frac{\text{Det}[L, N]}{2L \Delta N \Delta}
\]

where for leptons, \(L\) and \(N\) are the charged-lepton and neutrino mass matrices respectively\(^6\) (in an arbitrary weak basis) and \(L \Delta = (m_e - m_\mu)(m_\mu - m_\tau)(m_\tau - m_e)\) (with an analogous definition for \(N \Delta\) in terms of neutrino masses and likewise for the quarks). This is useful, as, despite \(J\) being defined purely in terms of mixing observables via Eq. (1), by contrast, Eq. (2) relates it to the mass matrices, which appear in the Standard Model Lagrangian.

We will discuss our recently proposed\(^3\) plaquette-invariant (ie. flavour-symmetric mixing) observables, which, in common with \(J\), are independent of flavour labels and can be simply related to the mass matrices. Again like \(J\), we find that our observables parameterise the violation of certain phenomenological symmetries which have already been considered significant\(^6\)\(^7\)\(^8\)\(^9\) in leptonic mixing. In the next section, we define more precisely what we mean by flavour symmetry.

---

\(^3\)We focus first on the leptons, although many of our considerations may be applied equally well to the quarks. In the leptonic case, neutrino mass eigenstate labels \(i = 1...3\) take the analogous role to the charge \(-\frac{1}{2}\) quark flavour labels in the quark case. In this sense, we will often use the term “flavour” to include neutrino mass eigenstate labels, as well as charged lepton flavour labels.

\(^4\)We use a cyclic labelling convention such that \(\beta = \alpha + 1, \gamma = \beta + 1, j = i + 1, k = j + 1\), all indices evaluated mod 3.

\(^5\)Throughout this paper, \(L\) and \(N\) are taken to be Hermitian, either by appropriate choice of the flavour basis for the right-handed fields, or as the Hermitian squares, \(MM^\dagger\), of the relevant mass or Yukawa coupling matrices. The symbols \(m_\alpha, m_i\) generically refer to their eigenvalues in either case.
3 The $S3_\uparrow \times S3_\uparrow$ Flavour Permutation Group

The $S3_\uparrow$ group is the group of the six possible permutations of the charged lepton flavours and/or of the charge $-\frac{1}{3}$ quark flavours, while the $S3_\uparrow$ group is the group of the six possible permutations of the neutrino flavours (ie. mass eigenstates) or of the charge $\frac{2}{3}$ quark flavours (the arrow subscript corresponds to the direction of the z-component of weak isospin of the corresponding left-handed fields). We consider all possible such permutations, which together constitute the direct product $S3_\uparrow \times S3_\uparrow$ flavour permutation group (FPG) with 36 elements.

We next consider the $P$ matrix (for “probability”\(^{10}\)) of moduli-squared of the mixing matrix elements, eg. for leptons:

$$P = \begin{pmatrix}
|U_{e1}|^2 & |U_{e2}|^2 & |U_{e3}|^2 \\
|U_{\mu1}|^2 & |U_{\mu2}|^2 & |U_{\mu3}|^2 \\
|U_{\tau1}|^2 & |U_{\tau2}|^2 & |U_{\tau3}|^2
\end{pmatrix}. \quad (3)$$

It should be familiar: for quarks, semileptonic weak decay rates of hadrons are proportional to its elements, while for leptons, the magnitudes of neutrino oscillation probabilities may be written in terms of its elements.\(^{10}\) Moreover, the $P$ matrix may easily be related to the fermion mass matrices, as we will see in Section 5 below. The $P$ matrix manifestly transforms as the natural representation of $S3_\uparrow \times S3_\uparrow$, the transformations being effected by pre- and/or post-multiplying by $3 \times 3$ real permutation matrices.\(^{7}\)

Jarlskog’s invariant $J$ is a pseudoscalar under the FPG; under even permutations, it is invariant, while under odd permutations (eg. single swaps of rows or columns of the mixing matrix, or odd numbers of them), it simply changes sign. This is our prototype Flavour Symmetric Mixing Observable (FSMO). As we commented in the previous section, it is easy to find other similar such quantities, which, surprisingly had not appeared in the literature until recently.\(^{3}\) There are two types of singlets under the S3 group: even (1) which remain invariant under all permutations, and odd ($\bar{1}$) which flip sign under odd permutations. So, under the FPG, there are four types of singlet: 1×1, $\bar{1}$×$\bar{1}$ (like $J$), 1×$\bar{1}$ and $\bar{1}$×1. By Flavour Symmetric Observables (FSOs), we mean observables with any of these transformation properties under the FPG. They may be functions of mixing matrix elements alone (FSMOs), or functions of mass eigenvalues alone, or functions of both.

Starting with elements of $P$ and combining and (anti-)symmetrising them over flavour labels in various ways, we find that, apart from their (trivial) overall normalisation, and possibly scalar offsets, there are a finite number of independent FSMOs at any given order in $P$. Enumerating them, we found that there are no non-trivial ones linear in $P$, while at 2nd order in $P$, there is only one each of 1×1, $\bar{1}$×$\bar{1}$. At third order, there is exactly one each of the four types of singlet, while at higher orders in $P$, there are multiple instances of each. Recognising that we need only four independent variables to specify the mixing, it is clearly enough to stop at third order, up to which, the singlets are essentially uniquely defined by their order in $P$ and their transformation property under the FPG.

4 Flavour-Symmetric Mixing Observables

We introduce four FSMOs,\(^{3}\) uniquely defined as outlined above:

<table>
<thead>
<tr>
<th>2nd Order in $P$:</th>
<th>3rd Order in $P$:</th>
</tr>
</thead>
<tbody>
<tr>
<td>$G = \frac{1}{2} \sum_{\alpha i} (P_{\alpha i})^2 - 1$</td>
<td>$C = \frac{3}{2} \sum_{\alpha i} (P_{\alpha i})^3 - (P_{\alpha i})^2 + 1$</td>
</tr>
</tbody>
</table>

\(^{7}\)Less obviously, any given plaquette of $P$ transforms as a 2-dimensional (real) irreducible representation of $S3_\uparrow \times S3_\uparrow$. 

\(^{10}\)
where \( L_{\gamma} = (P_{\alpha i} + P_{\beta j} - P_{\beta i} - P_{\alpha j}) \). Alternative, but equivalent definitions in terms of the elements of a single plaquette are given elsewhere. 3 Note that \( F \) is only quadratic in \( P \), because of the constraints of unitarity. We comment briefly on the normalisations and offsets we have given them. \( F \) and \( A \), being anti-symmetric, need no offset, as they are already centred on zero, which they reach for threefold maximal mixing (uniquely defined by all 9 elements of the mixing matrix having magnitude \( \frac{1}{\sqrt{3}} \)). \( G \) and \( C \) are defined with offsets such that they likewise vanish for threefold maximal mixing. All four variables are normalised so that their maximum value is unity, which they attain for no mixing. In Ref. 3, we also give the \( 1 \times 1 \) and the \( 1 \times \bar{1} \) FSMOs at 3rd order (called \( B \) and \( D \) respectively), but they will not concern us here.

The four FSMOs introduced in Eq. 4 are the simplest ones in terms of \( P \) and are sufficient to completely specify the mixing, up to a number of discrete ambiguities associated with the built-in flavour symmetry. \( J \) is of course not independent, and is given by 18.\( J = 1/6 - G + (4/3)C - (1/2)F^2 \). In Table 1, we summarise their properties and values (estimated at 90% CL from compilations of current experimental results) for both quarks \( ^{12} \) and leptons \( ^{13} \).

Table 1: Properties and values of flavour-symmetric mixing observables for quarks and leptons. The experimentally allowed ranges are estimated (90% CL) from compilations of current experimental results, neglecting any correlations between the input quantities.

<table>
<thead>
<tr>
<th>Observable Name</th>
<th>Order in ( P )</th>
<th>Symmetry: ( S_3 \times S_3^\dagger )</th>
<th>Theoretical Range</th>
<th>Experimental Range for Quarks</th>
<th>Experimental Range for Leptons</th>
</tr>
</thead>
<tbody>
<tr>
<td>( F )</td>
<td>2</td>
<td>( 1 \times \bar{1} )</td>
<td>((-1, 1))</td>
<td>((-0.14, 0.12))</td>
<td>((0.893, 0.896))</td>
</tr>
<tr>
<td>( G )</td>
<td>2</td>
<td>( 1 \times 1 )</td>
<td>((0, 1))</td>
<td>((0.15, 0.23))</td>
<td>((0.898, 0.901))</td>
</tr>
<tr>
<td>( A )</td>
<td>3</td>
<td>( 1 \times 1 )</td>
<td>((-1, 1))</td>
<td>((-0.065, 0.052))</td>
<td>((0.848, 0.852))</td>
</tr>
<tr>
<td>( C )</td>
<td>3</td>
<td>( 1 \times 1 )</td>
<td>((-\frac{1}{3}, 1))</td>
<td>((-0.005, 0.057))</td>
<td>((0.848, 0.852))</td>
</tr>
</tbody>
</table>

5  
Flavour-Symmetric Mixing Observables in Terms of Mass Matrices

Equation (2) gives \( J \), our prototype FSMO, in terms of the fermion mass matrices, which in turn are proportional to the matrices of Yukawa couplings which appear in the Standard Model Lagrangian. In this section, we show how to write the FSMOs of Section 4 above also in terms of the mass matrices. It is useful to define a reduced \( P \) matrix:

\[
\tilde{P} = P - D
\]

where \( D \) is the \( 3 \times 3 \) democratic matrix with all 9 elements equal to \( \frac{1}{3} \). We also define the reduced (ie. traceless) powers of the fermion mass matrices: \( L^m := L^m - \frac{1}{3} \operatorname{Tr}(L^m) \) (similarly for \( \tilde{N}^m \)), in terms of which, we can define the \( 2 \times 2 \) matrix of weak basis-invariants:

\[
\tilde{T}_{mn} := \operatorname{Tr}(\tilde{L}^m \tilde{N}^n), \quad m, n = 1, 2.
\]

For known lepton masses, \( \tilde{T} \) is completely equivalent to \( P \). In fact, it is straightforward to show that \( \tilde{P} \) is a mass-moment transform of \( \tilde{T} \):

\[
\tilde{P} = \tilde{M}_\ell^T \cdot \tilde{T} \cdot \tilde{M}_\ell
\]

where

\[
\tilde{M}_\ell = \frac{1}{L^\Delta} \left( \begin{array}{ccc}
\frac{m_\mu - m_\tau}{m_\mu - m_e} & \frac{m_\tau - m_e}{m_\mu - m_e} & \frac{m_e - m_\mu}{m_\mu - m_e} \\
\frac{m_\mu - m_\tau}{m_\mu - m_e} & m_\tau - m_e & m_e - m_\mu \\
\frac{m_\mu - m_e}{m_\mu - m_e} & m_e - m_\mu & m_\mu - m_e
\end{array} \right),
\]

\(^9\)They also treat the two weak-isospin sectors symmetrically, though this is not an essential feature.
with an analogous definition for $\tilde{M}_\nu$ (the inverse transform is easily obtained).

Starting from Eq. (4) and substituting for $P$ from Eqs. (5) and (7), we find that:

$$\mathcal{F} \equiv \det P = \frac{3 \det \tilde{T}}{L_\Delta N_\Delta}; \quad \text{cf. Eq. (2)}: \quad J = -i \frac{\det[L, N]}{2L_\Delta N_\Delta} \quad (9)$$

$$\mathcal{G} = \frac{\tilde{T}_{mn} \tilde{T}_{pq} L_{mp} N_{nq}}{(L_\Delta N_\Delta)^2}; \quad \mathcal{C}, \mathcal{A} = \frac{\tilde{T}_{mn} \tilde{T}_{pq} \tilde{T}_{rs} L_{cm,A} \ell_{c,A}^{(mpq)}}{(L_\Delta N_\Delta)^{n_c,A}}, \quad (10)$$

where the $L(N)$ are simple functions of traces of $\tilde{L}(\tilde{N})$, given in Ref. 3, and $n_C(n_A) = 2(3)$.

6 Application 1: Flavour-Symmetric Descriptions of Leptonic Mixing

The tribimaximal mixing\textsuperscript{2} ansatz for the MNS lepton mixing matrix:

$$U \simeq \begin{pmatrix} -2/\sqrt{6} & 1/\sqrt{3} & 0 \\ 1/\sqrt{6} & 1/\sqrt{3} & 1/\sqrt{2} \\ 1/\sqrt{6} & 1/\sqrt{3} & -1/\sqrt{2} \end{pmatrix} \quad (11)$$

is compatible with all confirmed leptonic mixing measurements from neutrino oscillation experiments, and may be considered a useful leading-order approximation to the data. It is defined by three phenomenological symmetries:\textsuperscript{6} $CP$ symmetry, $\mu$-$\tau$-reflection symmetry and Democracy, which may each be expressed (flavour-symmetrically) in terms of our FSMOs. For example, as is well known, the zero in the $U_{e3}$ position, if exact, ensures that no $CP$ violation can arise from the mixing matrix. $CP$ symmetry is thus represented simply by $J = 0$ (which is a necessary, but not sufficient condition for a single zero in the mixing matrix, see Section 7 below).

$\mu$-$\tau$-reflection symmetry\textsuperscript{7} means that corresponding elements in the $\mu$ and $\tau$ rows have equal moduli: $|U_{\mu i}| = |U_{\tau i}|$, $\forall i$, and this implies the two flavour-symmetric constraints:

$$\mathcal{F} = \mathcal{A} = 0 \quad (12)$$

(flavour symmetry means that although these two constraints imply just such a set of equalities, they do not define which pair of rows or columns are constrained). Democracy\textsuperscript{8 9} ensures that one row or column is trimaximally mixed, i.e. has the form $\frac{1}{2\sqrt{3}}(1, 1, 1)^T$, as is the case for the $\nu_2$ column in tribimaximal mixing. Democracy is ensured flavour-symmetrically by the two constraints:

$$\mathcal{F} = \mathcal{C} = 0. \quad (13)$$

Taking all three symmetries, tribimaximal mixing (or one of its trivial permutations) is ensured by the complete set of constraints $\mathcal{F} = \mathcal{C} = \mathcal{A} = J = 0$, which may be written as the single flavour-symmetric condition:

$$\mathcal{F}^2 + \mathcal{C}^2 + \mathcal{A}^2 + J^2 = 0. \quad (14)$$

Tribimaximal mixing is manifestly not flavour symmetric. The flavour-symmetry of our constraint, Eq. (14), is spontaneously broken by its tribimaximal solutions. The symmetry is manifested by the existence of a complete set of solutions of the generalised tribimaximal form, each related to the other by a member of the flavour permutation group.

Of course, generalisations of the tribimaximal form\textsuperscript{6} possessing subsets of its three symmetries may be similarly defined, and their corresponding flavour-symmetric constraints may be obtained by analogy to the above. These, and those of other special mixing forms\textsuperscript{14 15} are tabulated in Ref. 3.
7 Application 2: A Partially Unified, Flavour-Symmetric Description of Quark and Lepton Mixings

A unified understanding of quark and lepton mixings is highly desirable. This is difficult because their mixing matrices have starkly different forms: the quark mixing matrix is characterised by small mixing angles,\[12\] while the lepton mixing matrix is characterised mostly by large ones.\[13\] Many authors have ascribed this difference to the effect of the heavy majorana mass matrix in the leptonic case, via the see-saw mechanism.\[17\] Notwithstanding the attractiveness of this explanation, it is clearly still worthwhile to ask if there are any features of the respective mixings which the quark and lepton sectors have in common.

Neutrino oscillation data\[13\] require that \(|U_{e3}|^2 \lesssim 0.05\), significantly less than the other MNS matrix elements-squared. At least one small mixing element is hence a common feature of both quark and lepton mixing matrices. We are thus led first to ask the question: “what is the flavour-symmetric condition for at least one zero element in the mixing matrix?” We should perhaps anticipate two constraints, as the condition implies that both real and imaginary parts vanish. A zero mixing element implies \(CP\) conservation, so that \(J = 0\). A clue to the second constraint is that with \(\mu-\tau\)-reflection symmetry, \(J = 0\) ensures a zero somewhere in the \(\nu_e\) row of the MNS matrix. However, \(\mu-\tau\)-reflection symmetry implies two more constraints, Eq. (12).

In order to find a single additional constraint we consider the \(K\) matrix\[16\] with elements:

\[K_{\gamma k} = \text{Re}(U_{\alpha i}^* U_{\alpha j}^* U_{\beta k}^* U_{\beta j}),\]  

(15) which is the \(CP\)-conserving analogue of \(J\) (cf. the definition of \(J\), Eq. (1)). \(K\) should be familiar: in the leptonic case, its elements are often used to write the magnitudes of the oscillatory terms in neutrino appearance probabilities;\[10\] in the quark case, its elements are just the CKM factors of the \(CP\)-conserving parts of the interference terms in penguin-dominated decay rates. A single zero in the mixing matrix leads to four zeroes in a plaquette of \(K\) and this clearly implies:

\[\text{Det} K = 0,\]  

(16) which is our sufficient second condition, along with \(J = 0\).\[h\] We note that Eq. (16) can easily be cast in terms of our complete set of FSMOs, since \(54 \text{Det} K \equiv 2A + \mathcal{F}(\mathcal{F}^2 - 2\mathcal{C} - 1)\). Hence, \(\mu-\tau\)-reflection symmetry, Eq. (12), is a special case of Eq. (16).

Experimentally, there is no exactly zero element in the CKM matrix, so that \(\text{Det} K = 0\) and \(J = 0\) cannot both be exact for quarks. Moreover, for leptons, despite there being no experimental lower limit for \(|U_{e3}|\), there is no reason to suppose that the MNS matrix has an exact zero either. In order to ensure a small, but non-zero element in the mixing matrices, we need to consider a modest relaxation of either condition, or of both. For quarks, we know from experiment that \(CP\) is slightly violated, with\[12\] \(|J_q/\text{max}| \simeq 3 \times 10^{-4}\), while\[i\] for leptons, fits to oscillation data\[13\] imply a fairly loose upper bound on their \(CP\) violation: \(|J_\ell/\text{max}| \lesssim 0.33\). Turning to \(\text{Det} K\), we find that for quarks, \(|\text{Det} K_q/\text{(Det} K)_{\text{max}}| \lesssim 3 \times 10^{-7}\), while for leptons, \(|\text{Det} K_\ell/\text{(Det} K)_{\text{max}}| \lesssim 0.6\) (the precision of lepton mixing data does not yet allow a strong constraint). However, there is no experimental lower limit for \(|\text{Det} K|\) for quarks or for leptons, each being compatible with zero, so that it is sufficient to relax only the condition on \(J\).

We are thus led to conjecture that for both quarks and leptons:

\[\text{Det} K = 0; \quad |J/\text{max}| = \text{small}\]  

(17) (it is not implied that the small quantity necessarily has the same value in both sectors). Equation (17) is a unified and flavour-symmetric, partial description of both lepton and quark mixing

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\[h\]The two conditions may even be expressed as one, noting that the product of all nine elements of \(P\) is given by \(\prod_s \prod_{ij} P_{sij} = (\text{Det} K)^2 + J^2(2J^2 + R)^2\), which is zero iff \(\text{Det} K = 0\) and \(J = 0\) (as \(R > 0\), as long as \(J \neq 0\)).

\[i\]We note that \(J_{\text{max}} = \frac{1}{6\sqrt{3}} \simeq 0.1\) and \((\text{Det} K)_{\text{max}} = \frac{e^6}{3^5} \simeq 0.0033\).
matrices, being associated with the existence of at least one small element in each mixing matrix, \( U_{e3} \) and \( V_{ub} \) respectively (it is partial in the sense that only two degrees of freedom are constrained for each matrix). However, in the case that \( J \) is not exactly zero, the condition \( \text{Det} \ K = 0 \) also implies that in the limit, as \( J \to 0 \), there is at least one unitarity triangle angle which \( \to 90^\circ \). This is rather obvious in the \( \mu-\tau \)-symmetry case, but is less obvious more generally. While the flavour symmetry prevents an a priori prediction of which angle is \( \simeq 90^\circ \), we know from experiment\(^{12} \) that for quarks, \( \alpha \simeq 90^\circ \). A detailed calculation shows that our conjecture, Eq. (17), predicts, in terms of Wolfenstein parameters:\(^1 \)

\[
(90^\circ - \alpha) = \eta \lambda^2 = 1^\circ \pm 0.2^\circ
\]  

(18)

at leading order in small quantities, to be compared with its current experimental determination:\(^{12} \)

\[
(90^\circ - \alpha) = 0^\circ +3^\circ -7^\circ.
\]  

(19)

It will be interesting to test Eq. (18) more precisely in future experiments with \( B \) mesons, in particular, at LHCb and at a possible future Super Flavour Factory. For leptons, experiment tells us not only that it is the \( U_{e3} \) MNS matrix element which is small but also that only the unitarity triangle angles\(^j \phi_\mu_1 \) or \( \phi_\tau_1 \) can be close to \( 90^\circ \). Then Eq. (17) implies that:

\[
|90^\circ - \delta| = 2\sqrt{2} \sin \theta_{13} \sin (\theta_{23} - \frac{\pi}{4}) \lesssim 4^\circ
\]  

(20)

at leading order in small quantities (we use the PDG convention here). It thus requires a large \( CP \)-violating phase in the MNS matrix, which is promising for the discovery of leptonic \( CP \) violation at eg. a future Neutrino Factory.

8 Discussion and Conclusions

Given that our flavour-symmetric variables are defined (essentially) uniquely by their flavour symmetry properties and by their order in \( P \), it is remarkable that the leptonic data may be described simply by the constraints \( \mathcal{F} = \mathcal{A} = \mathcal{C} = J = 0 \). This is suggestive that these variables may be fundamental in some way. It is furthermore tantalising that the smallness of one element in each mixing matrix, the approximate \( \mu-\tau \)-symmetry in lepton mixing and the existence of a right unitarity triangle may all be related to each other, through our simple partially-unified constraint, Eq. (17). The precision of the resulting prediction, Eq. (18), motivates more sensitive tests at future \( B \) physics facilities, while the synergy with tests at a neutrino factory is manifest.

All elements of the Standard Model, apart from the Yukawa couplings of the fermions to the Higgs, treat each fermion of any given charge on an equal footing - they are already flavour-symmetric. The Yukawa couplings, on the other hand, depend on flavour in such a way that each flavour has unique mass and mixing matrix elements. Using our flavour-symmetric observables, or combinations of them appropriately chosen, we have shown how it is also possible to specify the flavour-dependent mixings in a flavour-independent way.\(^k \) This recovers flavour symmetry at the level of the mixing description, the symmetry being broken only spontaneously by its solutions, which define and differentiate the flavours in terms of their mixings.

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\(^{1}\)We use the nomenclature of unitarity triangle angles we defined in reference \([46]\) of Ref.\(^9 \).

\(^{k}\)We illustrated another variant of this in Ref.\(^{18} \).
References

Lepton flavour violation in constrained MSSM-seesaw models

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We calculate the predictions for lepton flavour violating (LFV) tau and muon decays, \( l_j \rightarrow l_i \gamma \), \( l_j \rightarrow 3l_i \), \( \mu \rightarrow e \) conversion in nuclei and LFV semileptonic tau decays \( \tau \rightarrow \mu PP \) with \( PP = \pi^0 \pi^-, \pi^0 \pi^0, K^0 K^-, K^0 \bar{K}^0 \) \( \tau \rightarrow \mu P \) with \( P = \pi^0, \eta, \eta' \) and \( \tau \rightarrow \mu V \) with \( V = \rho^0, \phi \), performing the hadronisation of quark bilinears within the chiral framework. We work within a SUSY-seesaw context where the particle content of the Minimal Supersymmetric Standard Model is extended by three right-handed neutrinos plus their corresponding SUSY partners, and where a seesaw mechanism for neutrino mass generation is implemented. Two different scenarios with either universal or non-universal soft supersymmetry breaking Higgs masses at the gauge coupling unification scale are considered. After comparing the predictions with present experimental bounds and future sensitivities, the most promising processes are particularly emphasised.

1 LFV within SUSY-seesaw models

The current knowledge of neutrino mass differences and mixing angles clearly indicates that lepton flavour number is not a conserved quantum number in Nature. However, the lepton flavour violation (LFV) has so far been observed only in the neutrino sector. One challenging task for the present and future experiments will then be to test if there is or there is not LFV in the charged lepton sector as well.

Here we focus in the Minimal Supersymmetric Standard Model (MSSM) enlarged by three right-handed neutrinos and their SUSY partners where potentially observable LFV effects in the charged lepton sector are expected to occur. We further assume a seesaw mechanism for neutrino mass generation and use, in particular, the parameterisation proposed in \(^1\) where the solution to the seesaw equation is written as \( m_D = Y_\nu v_2 = \sqrt{m_N^\text{diag} R} \sqrt{m_\nu^\text{diag} U^\dagger_{\text{MNS}}} \). Here, \( R \) is defined by \( \theta_i (i = 1, 2, 3) \); \( v_{1(2)} = v \cos(\sin(\beta), v = 174 \) GeV; \( m_\nu^\text{diag} = \text{diag}(m_{\nu_1}, m_{\nu_2}, m_{\nu_3}) \) denotes the three light neutrino masses, and \( m_N^\text{diag} = \text{diag}(m_{N_1}, m_{N_2}, m_{N_3}) \) the three heavy ones. \( U_{\text{MNS}} \) is given by the three (light) neutrino mixing angles \( \theta_{12}, \theta_{23} \) and \( \theta_{13} \), and three phases, \( \delta, \phi_1 \) and \( \phi_2 \). With this parameterisation is easy to accommodate the neutrino data, while leaving room for extra neutrino mixings (from the right-handed sector). It further allows for large Yukawa couplings \( Y_\nu \sim \mathcal{O}(1) \) by choosing large entries in \( m_N^\text{diag} \) and/or \( \theta_i \).

The predictions in the following are for two different constrained MSSM-seesaw scenarios, with universal and non-universal Higgs soft masses and with respective parameters (in addition to the previous neutrino sector parameters): 1) CMSSM-seesaw: \( M_0, M_{1/2}, A_0 \tan \beta \), and \( \text{sign}(\mu) \), and 2) NUHM-seesaw: \( M_0, M_{1/2}, A_0 \tan \beta \), \( \text{sign}(\mu) \), \( M_{H_1} = M_0(1 + \delta_1)^{1/2} \) and \( M_{H_2} = M_0(1 + \delta_2)^{1/2} \). All the predictions presented here include the full set of SUSY one-loop contributing diagrams, mediated by \( \gamma, Z \), and Higgs bosons, as well as boxes, and do not use the Leading Logarithmic (LLog) nor the mass insertion approximations. The hadronisation of quark bilinears is performed within the chiral framework. This is a very short summary of several publications \(^2,3,4,5\) to which we refer the reader for more details.

2 Results and Discussion

We focus on the dependence on the most relevant parameters which, for the case of hierarchical (degenerate) heavy neutrinos, are: the neutrino mass \( m_{N_3} (m_N) \), \( \tan \beta \), \( \theta_1 \) and \( \theta_2 \). We also
study the sensitivity of the BRs to θ_{13}. The other input seesaw parameters $m_{N_1}$, $m_{N_2}$ and $θ_3$, play a secondary role since the BRs do not strongly depend on them. The light neutrino parameters are fixed to: $m_{ν_2}^2 = Δm_{ν_1}^2 + m_{ν_1}^2$, $m_{ν_3}^2 = Δm_{ν_1}^2 + m_{ν_1}^2$, $Δm_{atm}^2 = 8 \times 10^{-5}$ eV$^2$, $Δm_{sol}^2 = 2.5 \times 10^{-3}$ eV$^2$, $m_{ν_1} = 10^{-3}$ eV, $θ_{12} = 30^0$, $θ_{23} = 45^0$, $θ_{13} ≲ 10^0$ and $δ = φ_1 = φ_2 = 0$.

The results for the CMSSM-seesaw scenario are collected in Figs. 1 through 5. In Fig. 1, we display the predictions of $BR(τ → µγ)$ and $CR(µ → e, Ti)$ as a function of the heaviest neutrino mass $m_{N_i}$, for the various SPS points, and for the particular choice $θ_i = 0$ ($i = 1, 2, 3$) and $θ_{13} = 5^0$. We have also considered the case of degenerate heavy neutrino spectra (not shown here). In both scenarios for degenerate and hierarchical heavy neutrinos, we find a strong dependence on the heaviest neutrino masses, with the expected behaviour $|m_{N_i} log m_{N_i}|^2$ of the LLog approximation, except for SPS 5 point, which fails by a factor of $∼ 10^4$. The rates for the various SPS points exhibit the following hierarchy, $BR_4 > BR_{1b} > BR_{1a} > BR_3 > BR_2 > BR_5$. This behaviour can be understood in terms of the growth of the BRs with tan β, and from the different mass spectra associated with each point. Most of the studied processes reach their experimental limit at $m_{N_3} ∈ [10^{13}, 10^{15}]$ which corresponds to $Y_{13}^{33,32} ∼ 0.1 - 1$. At present, the most restrictive one is $µ → eγ$ (which sets bounds for SPS 1a of $m_{N_3} < 10^{13} - 10^{14}$ GeV), although $µ → e$ conversion will be the best one in future, with a sensitivity to $m_{N_3} > 10^{12}$ GeV.

Figure 2: Dependence of LFV $τ$ and $µ$ decays with $|θ_1|$ for SPS 4 case with arg($θ_2$) = 0, $π/10$, $π/8$, $π/6$, $π/4$ in radians (lower to upper lines), $(m_{N_1}, m_{N_2}, m_{N_3}) = (10^{8}, 2 \times 10^{8}, 10^{14})$ GeV, $θ_2 = θ_3 = 0$, $θ_{13} = 0$ and $m_{ν_1} = 0$.

The horizontal lines are the present experimental bounds.
Fig. 2 shows the behaviour of the six considered LFV $\tau$ and $\mu$ decays, for SPS 4 point, as a function of $|\theta_1|$, for various values of $\arg\theta_1$. We see clearly that the BRs for $0 < |\theta_1| < \pi$ and $0 < \arg\theta_1 < \pi/2$ can increase up to a factor $10^2 - 10^4$ with respect to $\theta_1 = 0$. Similar results have been found for $\theta_2$, while BRs are nearly constant with $\theta_3$ in the case of hierarchical neutrinos. The behaviour of $\text{CR}(\mu - e, \text{Ti})$ with $\theta_i$ is very similar to that of $\text{BR}(\mu \rightarrow e\gamma)$ and $\text{BR}(\mu \rightarrow 3e)$. For instance, Fig. 3 shows the dependence of $\text{CR}(\mu - e, \text{Ti})$ with $\theta_2$, and illustrates that for large $\theta_2$, rates up to a factor $\sim 10^4$ larger than in the $\theta_1 = 0$ case can be obtained.

In Fig. 4 we show the dependence of $\mu \rightarrow e\gamma$, $\mu \rightarrow 3e$ and $\mu - e$ conversion on the light neutrino mixing angle $\theta_{13}$. These figures clearly manifest the very strong sensitivity of their rates to the $\theta_{13}$ mixing angle for hierarchical heavy neutrinos. Indeed, varying $\theta_{13}$ from 0 to $10^5$ leads to an increase in the rates by as much as five orders of magnitude.

On the other hand, since $\mu \rightarrow e\gamma$ is very sensitive to $\theta_{13}$, but $\text{BR}(\tau \rightarrow \mu \gamma)$ is clearly not, and since both BRs display the same approximate behaviour with $m_{N_3}$ and $\tan \beta$, one can study the impact that a potential future measurement of $\theta_{13}$ and these two rates can have on the knowledge of the otherwise unreachable heavy neutrino parameters. The correlation of these two observables as a function of $m_{N_3}$, is shown in Fig. 5 for SPS 1a. Comparing these predictions for the shaded areas along the expected diagonal “corridor”, with the allowed experimental region, allows to conclude about the impact of a $\theta_{13}$ measurement on the allowed/excluded $m_{N_3}$ values. The most important conclusion from Fig. 5 is that for SPS 1a, and for the parameter space defined in the caption, an hypothetical $\theta_{13}$ measurement larger than $1^\circ$, together with the present experimental bound on the $\text{BR}(\mu \rightarrow e\gamma)$, will have the impact of excluding values.
of $m_{N_3} \gtrsim 10^{14}$ GeV. Moreover, with the planned MEG sensitivity, the same $\theta_{13}$ measurement could further exclude $m_{N_3} \gtrsim 3 \times 10^{12}$ GeV.

The numerical results for the NUHM-seesaw scenario as a function of $M_0 = M_{1/2} = M_{\text{SUSY}}$ are collected in Figs. 6 and 7. The behaviour of the predicted $m_{H^0}$ as a function of $M_{\text{SUSY}}$ is shown in Fig. 6 (left panel). The most interesting solutions with important phenomenological implications are found for negative $\delta_1$ and positive $\delta_2$. Notice that, for all the explored $\delta_{1,2}$ values, we find a value of $m_{H^0}$ that is significantly smaller than in the universal case ($\delta_{1,2} = 0$).

In Fig. 6 (right panel) the various contributions from the $\gamma$, $Z$, Higgs mediated penguins and box diagrams as a function of $M_{\text{SUSY}}$ are shown. Here, we choose $\delta_1 = -1.8$ and $\delta_2 = 0$. We observe a very distinct behaviour with $M_{\text{SUSY}}$ of the Higgs-mediated contributions compared to those of the CMSSM case. In fact, the Higgs-mediated contribution can equal, or even exceed that of the photon, dominating the total conversion rate in the large $M_0 = M_{1/2}$ region. These larger Higgs contributions are the consequence of their exclusive SUSY non-decoupling behaviour for large $M_{\text{SUSY}}$, and of the lighter Higgs boson mass values encountered in this region, as previously illustrated in Fig. 6.

In Fig. 7 we display the predicted $\mu - e$ conversion rates for other nuclei, concretely Al, Ti,
we have also found $M = \tau \rightarrow \theta e = 0$. From top to bottom, the $\mu \rightarrow \pi^-\eta$ decays $\tau \rightarrow \mu$ mixing driven by $\delta_1 = 2$ and $\delta_2 = 0$. From top to bottom, the horizontal dashed lines denote the present experimental bounds for $\text{CR}(\mu - e, \text{Ti})$ and $\text{CR}(\mu - e, \text{Au})$.

Sr, Sb, Au and Pb, as a function of $M_{\text{SUSY}}$. We clearly see that $\text{CR}(\mu - e, \text{Sh}) > \text{CR}(\mu - e, \text{Sb}) > \text{CR}(\mu - e, \text{Ti}) > \text{CR}(\mu - e, \text{Au}) > \text{CR}(\mu - e, \text{Pb}) > \text{CR}(\mu - e, \text{Al})$. The most important conclusion from Fig. 7 is that we have found predictions for Gold nuclei which, for the input parameters in this plot, are above its present experimental bound throughout the explored $M_{\text{SUSY}}$ interval. Finally, although not shown here for shortness, we have also found an interesting loss of correlation between the predicted $\text{CR}(\mu - e, \text{Ti})$ and $\text{BR}(\mu \rightarrow e\gamma)$ in the NUHM-seesaw scenario compared to the universal case where the Higgs-contributions dominate the photon-contributions and could be tested if the announced future sensitivities in these quantities are reached.

The corresponding predictions for $\theta_2 = 2.9e^{i\pi/4}$ of the nine LFV semileptonic $\tau$ decays studied in this work as a function of $M_{\text{SUSY}}$ are shown in Fig. 8. In this case, we work with $\delta_1 = -2.4$ and $\delta_2 = 0.2$, that drive us to Higgs boson masses around 150 GeV even for heavy SUSY spectra. In this Fig. 8 we can see that, the choice of $\theta_2$ increase all the rates about two orders of magnitude respect to the case $\theta_1 = 0$, not shown here for brevity. $\text{BR}(\tau \rightarrow \mu \pi^+\pi^-)$ and $\text{BR}(\tau \rightarrow \mu \rho)$ the largest rates and, indeed, the predictions of these two latter channels reach their present experimental sensitivities at the low $M_{\text{SUSY}}$ region, below 200 GeV and 250 GeV respectively, for this particular choice of input parameters.

In Fig. 9 we plot finally the predictions for $\text{BR}(\tau \rightarrow \mu K^+K^-)$ and $\text{BR}(\tau \rightarrow \mu\eta)$ as a function of one the most relevant parameters for these Higgs-mediated processes which is the corresponding Higgs boson mass.
Firstly, we see that the approximate and exact results of the Higgs contribution agree within a factor of two for both channels, but the agreement of the full result with respect to the Higgs contribution is clearly worse in the case of $\tau \rightarrow \mu K^- K^-$ than in $\tau \rightarrow \mu \eta$. In the latter, the agreement is quite good because the $Z$-mediated contribution is negligible, and this holds for all $M_{\text{Susy}}$ values in the studied interval, $250$ GeV $< M_{\text{Susy}} < 750$ GeV. In the first, it is only for large $M_{\text{Susy}}$ that the $H$-mediated contribution competes with the $\gamma$-mediated one and the Higgs rates approach the total rates. For instance, the predictions for $\text{BR}(\tau \rightarrow \mu K^- K^-)$ shows that for $M_{\text{Susy}} = 750$ GeV and $m_H = 160$ GeV the total rate is about a factor 2 above the Higgs rate, but for $m_H = 240$ GeV it is already more than a factor 5 above.

In this figure we have also explored larger values of $m_{N_3}$ and $\tan \beta$, by using in those cases the approximate formula, and in order to conclude about the values that predict rates comparable with the present experimental sensitivity. We can conclude then that, at present, it is certainly $\tau \rightarrow \mu \eta$ the most competitive LFV semileptonic tau decay channel. The parameter values that provide rates being comparable to the present sensitivities in this channel are $\tan \beta = 60$ and $m_{N_3} = 10^{15}$ GeV which correspond to $|\delta_{12}| \approx 2$.

Interestingly, the most competitive channels to explore simultaneously LFV $\tau - \mu$ transitions and the Higgs sector are $\tau \rightarrow \mu \eta$, $\tau \rightarrow \mu \eta'$ and also $\tau \rightarrow \mu K^+ K^-$. Otherwise, the golden channels to tackle the Higgs sector are undoubtly $\tau \rightarrow \mu \eta$ and $\tau \rightarrow \mu \eta'$. On the other hand, the rest of the studied semileptonic channels, $\tau \rightarrow \mu \pi^+ \pi^-$, etc., will not provide additional information on LFV with respect to that provided by $\tau \rightarrow \mu \gamma$.

In conclusion, we believe that a joint measurement of the LFV branching ratios, the $\mu - e$ conversion rates, $\theta_{13}$ and the SUSY spectrum will be a powerful tool for shedding some light on the otherwise unreachable heavy neutrino parameters. Futhermore, in the case of a NUHM scenario, it may also provide interesting information on the Higgs sector. It is clear from this study that the connection between LFV and neutrino physics will play a relevant role for the searches of new physics beyond the SM.

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Top physics
TOP PAIR PRODUCTION

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We review the most recent results on top quark pair production at the Tevatron including production mechanism, cross sections, forward-backward asymmetry measurements and searches for resonances decaying to top quarks which were available at the time of the 2008 electroweak session of the Rencontres de Moriond conference.

1 Introduction

Since its discovery by the CDF and D0 experiments the Tevatron is the only place where the top quark (t) can be studied. The year 2008 will certainly be the last possible year for such a statement due to the advent of the LHC.

Doing top quark physics means covering a wide spectrum of different subjects including studies of the t (single and pair) production, decay and properties. The present mini-review focuses on top quark pair (tt) production and the emphasis is put on recent results concerning the tt production mechanism, cross section measurements and top quark mass (mt) measurements from cross sections measurements, forward backward measurement and finally searches of resonances decaying into t quarks. Recent results concerning mt direct measurements as well as other properties (W helicity, t charge) and the single t production can be found in these proceedings. The Tevatron is performing well and the results reported here correspond to Tevatron integrated luminosities ranging from 1 to 2 fb⁻¹.

At the Tevatron, within the Standard Model (SM), tt production is expected to occur via strong interactions namely through q̅q annihilation (85%) and gluon gluon (gg) fusion (15%). Typical next-to-leading (NLO) predictions range from \( \sigma(p\bar{p}) \rightarrow tt \approx 6.7 \pm 0.4 \text{ pb} \) for \( m_t = 175 \text{ GeV} \) to \( \sigma(p\bar{p}) \rightarrow tt \approx 7.8 \pm 0.5 \text{ pb} \) for \( m_t = 170 \text{ GeV} \). The t is expected to decay before it hadronizes. The t decays into a b quark and an on-shell W gauge boson (t → Wb) with a branching ratio close to 1. The final states corresponding to tt production are classified...
according to the decay of the W gauge boson from the t. The results reported here concentrate on the lepton+jets channels where one W gauge boson decays leptonically i.e. $W \rightarrow l\nu$ where $l = \mu$ or e (amounting to 30% of all the $t\bar{t}$ channels) and dilepton channels where both W decay leptonically (5%). Other channels include all hadronic channels where both W gauge bosons decay hadronically ($W \rightarrow q\bar{q}'$, 45%) as well as tauonic channels where both W decay leptonically into tau leptons (20%).

The main physics backgrounds from SM processes are also decay channel dependent. The main SM physics backgrounds for $t\bar{t}$ signals in the lepton+jets channel come from $W$+jets production as well as multijets production where one jet fakes an electron or a muon. In the dilepton channel the main physics backgrounds come from $Z$ gauge boson production decaying into a lepton pair, Drell Yan processes and gauge boson pair production. Typical event selections require high $p_T$ lepton (> 15 to 20 GeV), large missing transverse energy (> 15 to 20 GeV) and jets with large transverse energies (> 15 to 20 GeV). They also include cuts on several kinematical variables. In several analyses the selection uses of $b$-quark jets identification based either on displaced vertices (with efficiencies ranging from 50 to 60 %) or soft lepton taggers.

2 Production mechanism

The evaluation of the $gg$ fusion process in $t\bar{t}$ production suffers from theoretical uncertainties and can vary up to a factor of 2. This motivates the CDF experiment to perform a measurement of the relative fraction $C_f$ of gluon gluon fusion ($t\bar{t}gg$) versus quark antiquark annihilation ($t\bar{t}q\bar{q}$) by combining two complementary methods which were already reported elsewhere before this conference. The new CDF result presented at this conference concerns the combination.

The first method is a data driven method based on low $p_T$ tracks. Because gluons can radiate other partons and gluons, $(t\bar{t}gg)$ should have more low $p_T$ tracks. The shapes of the track $p_T$ distributions of the two components ($t\bar{t}gg$ and $t\bar{t}q\bar{q}$) are derived from inclusive dijets and $W$ + n jets (where n = 0, 1, 2) data samples. The background shape is then constructed as a combination of the $t\bar{t}gg$ and $t\bar{t}q\bar{q}$ shapes. The three shapes are then fit to the data sample selected for the $t$ signal in the lepton+jets channel. The second method uses the kinematics of the production and the decay of the $t\bar{t}$ to differentiate the two production mechanisms. The kinematic variables are used to train a neural network (NN) to increase the sensitivity of the method. This analysis relies on Monte Carlo (MC). A large ensemble of pseudo-experiments (PSE) are generated to calculate the statistical and systematical uncertainties and the Feldman- Cousins (FC) method is used to make the measurement. The track method and the NN method are then combined into a combined PSE method. Using a total integrated luminosity of 995 pb$^{-1}$ the CDF experiment finds: $C_f = 0.07^{+0.15}_{-0.07}$ at 68 % confidence level.

3 Top quark pair production cross section and top quark mass from cross sections measurements

3.1 Top quark pair production cross section measurements

The D0 experiment performed a new measurement of the $t\bar{t}$ production cross section in the lepton+jets channel employing two complementary methods for discrimination between signal and background namely using a likelihood discriminant and $b$-tagging. This new measurement is based on about 0.9 fb$^{-1}$ of data.

Events with $t\bar{t}$ decays differ in their event kinematics from background events. However no single kinematic quantity can separate signal and background very well. This motivates the development of the likelihood discriminant method which uses up to 6 kinematical quantities in each channel to discriminate the $t\bar{t}$ signal from the backgrounds. Four channels are defined
Figure 1: Summary of $t\bar{t}$ production cross-section measurements from the D0 experiment available at the time of the 2008 EW session of the Rencontres de Moriond conference.

by lepton flavor ($e, \mu$) and jet multiplicity ($3, \geq 4$).

The probability density functions of the likelihood discriminant is determined from MC for $t\bar{t}$ signal and prompt lepton backgrounds and from a control data sample for multijets backgrounds (backgrounds without prompt leptons) both using the TMVA method. A maximum likelihood fit to the distribution of the likelihood discriminant from the data is then performed in all four channels simultaneously with the $t\bar{t}$ production cross section as a free parameter.

The $b$-tagging method exploits the fact that every $t\bar{t}$ decay produces two $b$ quark to distinguish them from the backgrounds. The signal over background ratio is enhanced by requiring at least one $b$-tagged jet. The $t\bar{t}$ signal and prompt lepton backgrounds are modeled with the MC and the backgrounds from multijets events are determined from the data. The cross section is calculated using a maximum likelihood fit to the number of events in eight different channels defined by the lepton flavor ($e, \mu$), jet multiplicity ($3, \geq 4$) and $b$-tag multiplicity ($1, \geq 2$).

Combining the likelihood discriminant and the $b$-tagging methods with the help of the method described in, the D0 experiment measures the $t\bar{t}$ production cross section in the lepton+jets channel using a total integrated luminosity of 910 pb$^{-1}$: $\sigma(pp \rightarrow t\bar{t}) = 7.77 \pm 0.54({\text{stat}}.) \pm 0.47({\text{syst}}.) \pm 0.47({\text{lumi}}.)$ pb for $m_t = 170$ GeV and $\sigma(pp \rightarrow t\bar{t}) = 7.42 \pm 0.53({\text{stat}}.) \pm 0.46({\text{syst}}.) \pm 0.45({\text{lumi}}.)$ pb for $m_t = 175$ GeV.

Figure 1 shows that the measurements are consistent with each other and consistent with the SM predictions. New physics would show up as inconsistencies.

3.2 Top quark mass from cross sections measurements

The value of $m_t$ can vary significantly depending on its different possible (and related) definitions from the running $m_t$ definition in the (for example) $\hat{MS}$ scheme (from the 1-loop up to the 3-loop level) to the $m_{\text{pole}}$ which is itself defined up to some ambiguities such as the known renormalon ambiguity.

At the Tevatron, the $m_t$ measurements are performed by using template, ideogram, neutrino weighting or matrix element ‘direct’ methods. They rely on the detailed description of the $t\bar{t}$ production in the MC simulations which currently contain only matrix elements at the leading order (LO) of quantum chromodynamics (QCD) and higher orders are simulated by applying
parton showers thus leaving in principle the $m_t$ convention unknown. Therefore the world average of $m_t$ is extracted in a not very well-defined scheme. The $t$ quark mass can also be measured from the $t\bar{t}$ production cross section measurements. These ‘indirect’ measurements will thus provide valuable complementary informations on the value of $m_t$. Although efforts are put in improving their accuracy they are not meant to compete in precision with the ‘direct’ $m_t$ measurements. The measurement of $m_t$ from the most recent $t\bar{t}$ production cross section measurement from the D0 experiment reported in subsection 3.1 was not available at the time of the 2008 EW session of this conference but can be found in its QCD session 15. Therefore we will only mention the results obtained with the previous set of cross-section measurements of summer 2007 16 corresponding to an integrated luminosity of 910 pb$^{-1}$ for the lepton+jets channel and 1.05 fb$^{-1}$ for the dilepton channel. Comparing the cross section measurements in the lepton+jets channel with the predictions of 3 and 4 respectively, leads to $m_t = 166.9^{+5.9}_{-5.2}(\text{stat.} + \text{syst.})^{+3.7}_{-3.8}(\text{theory})$ GeV, and $m_t = 166.1^{+6.1}_{-5.3}(\text{stat.} + \text{syst.})^{+4.9}_{-6.7}(\text{theory})$ GeV respectively. This can be compared with the direct measurement from the D0 experiment with the matrix element method 19 $m_t = 170.5 \pm 2.4(\text{stat.} + \text{JES}) \pm 1.2(\text{syst.})$ GeV and with the 2007 world average $m_{top} = 170.9 \pm 1.1(\text{stat.}) \pm 1.5(\text{syst.})$ GeV.

Comparing the measurements in the dilepton channel 18 and predictions leads to $m_t = 174.5^{+10.5}_{-8.2}(\text{stat.} + \text{syst.})^{+3.7}_{-3.6}(\text{theory})$ GeV and $m_t = 174.1^{+9.8}_{-8.4}(\text{stat.} + \text{syst.})^{+4.2}_{-6.0}(\text{theory})$ GeV respectively. This can be compared with the direct measurement from the D0 experiment with the neutrino weighting method 20 $m_t = 172.5 \pm 5.8(\text{stat.}) \pm 5.5(\text{syst.})$ GeV.

The CDF experiment performed a new $m_t$ measurement using the $t\bar{t}$ production cross section measurement in the dilepton channel, with an integrated luminosity of 1.2 fb$^{-1}$, as a constraint. Since the number of $t\bar{t}$ signal events depends on $m_t$, the observed number of events can therefore be used to measure $m_t$.

The kinematics of the $t\bar{t}$ system in the dilepton channel data sample is solved using the information on the momentum z-component of the $t\bar{t}$ system taken from the $t\bar{t}$ data sample in the lepton+jets channel. Solving the kinematics of the $t\bar{t}$ system in the dilepton channel allows to reconstruct $m_t$. The CDF experiment then uses a likelihood fit to get the final $m_t$ measurement. The reconstructed $m_t$ distribution from data is compared to MC signal and backgrounds templates and the number of events is compared to the expected number of events. The result of the likelihood fit gives: $m_t = 170.7^{+4.2}_{-3.9}(\text{stat.}) \pm 2.6(\text{syst.}) \pm 2.4(\text{theory})$ GeV.

4 Forward backward asymmetry

At the Tevatron the $t\bar{t}$ production is predicted to be charge symmetric at LO in QCD. However NLO calculations predicts asymmetries in the 5%-10% range 23 and next-to-next-to-leading order (NNLO) calculations predict significant corrections for $t\bar{t}$ production in association with a jet 24. The charge asymmetry arises from the interferences between symmetric and antisymmetric contributions under the exchange $t \leftrightarrow \bar{t}$. The charge asymmetry depends on the region of phase space being and, in particular, on the production of an additional jet. The small asymmetries expected in the SM makes this a sensitive probe for new physics 25.

Using a data sample corresponding to an integrated luminosity of about 0.9 fb$^{-1}$, the D0 experiment performed the first measurement of the forward-backward charge asymmetry in $t\bar{t}$ production in the lepton+jets channel 26. The forward-backward charge asymmetry can be obtained from the signed difference between the rapidities of the $t$ and $\bar{t}$, $\Delta y = y_t - y_{\bar{t}}$ where the rapidity $y$ is defined as function of the polar angle $\theta$ and the ratio of the particle’s momentum to its energy $\beta$ as $y(\theta, \beta) = \frac{1}{2} \ln [(1 + \beta \cos \theta)/(1 - \beta \cos \theta)]$. The asymmetry is defined as:

$$A_{fb} = \frac{N_{\Delta y>0} - N_{\Delta y<0}}{N_{\Delta y>0} + N_{\Delta y<0}},$$

(1)
where \( N^{\Delta y>0} \) (\( N^{\Delta y<0} \)) is the number of event with positive (negative) \( \Delta y \).

Using a data sample with one lepton+ \( n \) jets, where \( n \geq 4 \) one jet at least being b-tagged in order to enhance the signal, the kinematics of the \( t\bar{t} \) is reconstructed with the help a kinematic fitter which varies the 4-momenta of the detected objects within their resolutions and minimizes a \( \chi^2 \) statistics, constraining both the known W gauge boson mass (\( M_W \)) and \( m_t \).

The sample composition, including \( t\bar{t} \) signal and W+jets from MC simulations and multijet background from data samples that fail lepton identification, as well as \( A_{fb} \) are then extracted from a simultaneous maximum-likelihood fit to data.

The observed asymmetry, uncorrected for acceptance and reconstruction effects, are \( A_{fb}^{obs} = 0.12 \pm 0.08(stat.) \pm 0.01(syst.) \) for \( n_{jets} \geq 4 \), \( A_{fb}^{obs} = 0.19 \pm 0.09(stat.) \pm 0.02(syst.) \) for \( n_{jets} = 4 \) and \( A_{fb} = -0.16_{-0.17}^{+0.15}(stat.) \pm 0.03(syst.) \) for \( n_{jets} \geq 5 \).

Using a lepton+(at least 4) jets sample, where at least one jet is b-tagged, corresponding to an integrated luminosity of 1.9 fb\(^{-1} \) and containing 484 candidates events, the CDF experiment performed a forward-backward asymmetry defined by 27:

\[
A_{fb} = \frac{N_{-Q_t \cos \Theta > 0} - N_{-Q_t \cos \Theta < 0}}{N_{-Q_t \cos \Theta > 0} + N_{-Q_t \cos \Theta < 0}},
\]

where \( \Theta \) is the production angle of the \( t \) i.e. the angle between the \( t \) and the proton beam, and \( Q_t \) is the charge of the lepton.

The \( t \) production angle in the lepton+jets final state is reconstructed by using a kinematic fitter. In order to compare to the theoretical prediction any bias and smear of the \( t\bar{t} \) asymmetry due to backgrounds, acceptance, and reconstruction have to be taken into account. The CDF experiment uses MC simulations to simulate these effects.

Including the reconstruction and acceptance corrections the forward backward asymmetry is measured to be \( A_{fb} = 0.17 \pm 0.07(stat.) \pm 0.04(syst.) \).

The measured is consistent (at the 2\( \sigma \)) level with the prediction 0.04 from the NLO MC generator MC\@NLO 28.

The CDF experiment performed a cross-check to the measurement by reweighting the \( t\bar{t} \) MC signal distribution to have a ‘true’ \( A_{fb} = 0.17 \). A Kolmogorov-Smirnov test has been performed to compare the shape of the reweighted distribution with backgrounds and data resulting into a probability of 45.6% showing a good agreement.

Due to different \( A_{fb} \) definitions and due to the usage (CDF) or not (D0) of acceptance and reconstruction corrections, the D0 and CDF results on \( A_{fb} \) are not to be compared.

5 Searches for resonances

The \( t \) is known so far as being the heaviest elementary particle. The production of \( t\bar{t} \) can be sensitive to physics beyond standard model in particular top-color and unknown heavy resonances 29, heavy Higgs boson decaying to \( t \) 30, \( t\bar{t} \) condensation 31, massive Z gauge boson in extended gauge theories 32, Kaluza-Klein states of the Z gauge boson or gluon 33 and axigluons 34. Such new effects may produce resonances in the \( t\bar{t} \) invariant mass distribution or may interfere with SM processes and cause distortion to the shape of this invariant mass distribution.

Using the same data sample in the lepton+jets channel as described in section 4 allowing also for a second b-tagged jet, the CDF experiment performed a measurement of the \( t\bar{t} \) differential cross section with respect to the invariant \( t\bar{t} \) mass \( d\sigma/dM_{t\bar{t}} \). The \( t\bar{t} \) invariant mass is reconstructed by combining the 4-momenta of the 4 leading jets, lepton and missing transverse energy. The neutrino momentum is taken from the missing transverse energy, the longitudinal component \( p_z \) of the neutrino being obtained by constraining the lepton and the neutrino invariant mass to be equal to \( M_W \).
The reconstructed $M_{t\ell}$ distribution is distorted from the true distribution by detector effects, resolutions and acceptances. These effects are corrected by using a regularized unfolding technique i.e. Singular Value Decomposition (SVD) \cite{SVD}. The CDF experiment then uses an Anderson-Darling \cite{AD} statistic to test for discrepancies from the standard model expectation. No evidence of inconsistencies with the Standard Model is seen, with an observed p-value of 0.45.

Using the same data sample as above the CDF experiment also searched for massive gluons decaying into $t\bar{t}$ \cite{CDFGluons}. In this search $M_{t\ell}$ is reconstructed event-by-event using the Dynamical Likelihood Method (DLM) \cite{DLM} also used for one of the CDF experiment $m_t$ measurement \cite{CDFmt}. After reconstructing $M_{t\ell}$, an unbinned likelihood fit is performed to extract the coupling strength. The fitted coupling strengths are consistent with the SM prediction within 1.7$\sigma$ in the width over coupling ratio range from 0.05 to 0.5 for a massive gluon mass range from 400 to 800 GeV.

The D0 experiment searched for a narrow-width heavy resonance $X$ decaying into $t\bar{t}$ using a lepton+jets sample with at least one $b$-tagged jet corresponding to an integrated luminosity of 0.9 pb$^{-1}$ \cite{D0Resonance}. The $t\bar{t}$ invariant mass is reconstructed in the same way as described above for the CDF $d\sigma/dM_{t\ell}$ measurement. Model independent upper limits on $\sigma_XBR(X \to t\bar{t})$ have been obtained using a bayesian method \cite{D0Bayesian}. Within a top-color-assisted technicolor model, the existence of a leptophobic $Z'$ boson with $M_{Z'} < 690$ GeV and width $\Gamma_{Z'} = 0.012M_{Z'}$ is excluded at 95\% confidence level.

An updated result was just available for the 2008 QCD session of the Rencontres de Moriond conference \cite{D0Resonance2}. With a data sample corresponding to an integrated luminosity of 2.1 pb$^{-1}$, the the existence of a leptophobic $Z'$ boson with $M_{Z'} < 690$ GeV and width $\Gamma_{Z'} = 0.012M_{Z'}$ is excluded at 95\% confidence level.

6 Conclusions

We review the most recent results on $t\bar{t}$ production at the Tevatron which were available at the time of the 2008 electroweak session of the Rencontres de Moriond conference and corresponding to about 1 to 2 fb$^{-1}$ of integrated luminosity for each of the CDF and D0 experiments. These results include production mechanism, cross sections and forward-backward asymmetry measurements which are found to be consistent with the SM expectations. The $t\bar{t}$ production cross section measurements allow for complementary $m_t$ measurements which can be compared to direct measurements. There are no evidence so far for resonances decaying into $t$ and model independent limits on masses as well as parameters of the different possible theoretical frameworks have been set. More data and results are expected to come after the winter 2008 as the Tevatron is continuing to perform very well.

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TOP MASS AND PROPERTIES

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The top quark was discovered in 1995. The top quark mass is now well measured at the Tevatron, with uncertainty getting below 1% of the top mass. The world average from last year was 176.9 ± 1.8 GeV/c². The new CDF measurement is 172 ± 1.2 (stat) ± 1.5 (sys) GeV/c², and D0 will soon present a new measurement. The top quark mass is an important parameter in the Standard Model, and should be measured as precisely as possible. To learn more about the top quark observed and study possible new physics, other properties also should be measured. At the Tevatron, the charge of the top quark can be measured directly.

Examples of other properties studied and reported in this presentation are W helicity, top decay branching ratio to b (Rb), searches for t → Hb and for flavor changing neutral current (FCNC). The results are all consistent with the Standard Model within current statistics. With significantly more data being collected at the Tevatron, precision measurements of the top properties are just starting.

1 Introduction

Top quarks are produced at the Tevatron mainly in top anti-top pairs, $P \bar{P} \rightarrow t \bar{t}$, through quark anti-quark annihilation and gluon gluon fusion. The t (\bar{t}) quark subsequently decays into a $W^+(W^-)$ boson and a b(\bar{b}) quark, $t \rightarrow Wb$, with a branching ratio close to 1. From the b's and the final products of the W decays, the mass and other properties of the top quark can be measured.

The $P \bar{P} \rightarrow t \bar{t} \rightarrow W^+bW^-\bar{b}$ production cross section and event selection have been reported in a previous talk at this Conference. Based on how the W's decay, three analysis channels are identified: the di-lepton channel (DIL) for both W's decaying leptonically, the lepton plus jet channel (LJ) for only one W decaying leptonically, and the all hadron channel for both W's decaying hadronically. (In this article we consider the final leptons being electron or muon only. The case of $W \rightarrow \tau \nu$ has to be handled differently due to the nature of tau decay.) Each channel
has its own challenges and strengths. Some common methods are developed and used when applicable.

2 top mass measurement

Mass is a fundamental property of a particle. While the top has been discovered for more than ten years, we have many interesting questions about the top quark. Are we seeing the same particle in all three analysis channels (DIL, LJ and all hadron)? Precise measurement of the top mass in these three channels could provide some insight into this question. If this particle seen as top quark is truly the one of the SM, then since its mass strongly correlates with the mass of Higgs particle, precise measurement of the top mass can help the search of the Higgs particle, and will also enable stringent constraints for electroweak tests and new physics.

The three channels, DIL, LJ and all hadron, have their own challenges and methods. In the DIL channel, each event has two neutrinos that are not measured directly, with only the missing transverse energy providing partial information for these two neutrinos. The system is under-constrained. Additional assumption is used to further constrain the system to be able to reconstruct the top mass. Also, various top mass could be used as input to obtain a probability density function to determine the most probable value for top mass.

A general issue with all three channels is which lepton and jet(s) in each $t\bar{t}$ event are decay products of the top quark and which are from the anti-top? One could try all possible combinations and select one based on reconstruction probability or simple kinematic information, such as the invariant mass of the top and anti-top system. Alternatively, one could use all possible solutions and assign weights based on the some relevant quantities, such as a weight defined by comparing missing Et from the reconstruction to that from what is measured in each event. These techniques are also generally applied in studying properties of the top quark.

2.1 The methods

One common issue with all channels is the jet energy calibration. In prior analysis jet energy calibration was based on predefined cone sizes. In an event where at least one $W$ decays hadronically, the known $W$ boson mass can be used as input to further fine tune the two jets associated with this $W$. This is called the in-situ jet energy calibration. In top mass measurement, this $W$ mass constraint is applied to the final events selected to find the average shift to the nominal jet energy calibration. The shift is applied to all jets including the $b$ jets. This procedure significantly improves the determination of the uncertainty in the top mass measurement.

The Template Method is one of the main methods used to obtain the top mass. In this method, top mass is reconstructed from the kinematic information available in the event. Templates are formed based on Monte Carlo with different top mass input. Comparing these templates with the observed events reconstructed in the same way reveals the top mass. In each DIL channel event, the system is under-constrained. Additional reasonable assumption has to be made, such as taking $P_{z}$ of top anti-top system from observed events¹, or weighting on the $\phi$ angle of the neutrino², etc. In the LJ channel, the assumption is made that the missing energy is due to the neutrino being not detected. A top mass fitter is used to find the most probable top mass, taking into account the resolution of $P_{t}$ and jet energy measured. In-situ jet energy calibration is commonly applied to improve the uncertainty. In each all hadronic event, there are two $W$'s decaying hadronically. In-situ jet energy calibration is generally applied to the jets which form the $W$'s. In the Template Method with 2-dimensional fit (MT2D)³ analysis in CDF, all possible jet pairing combinations are tried but only the one with best $\chi^{2}$ is kept.

The Matrix Element Method is based on theory. This takes into account all the kinematics information contained in an event, which are top mass dependent. A conditional probability
Table 1: top mass measurement at the Tevatron

<table>
<thead>
<tr>
<th>Analysis</th>
<th>Samples</th>
<th>Result</th>
</tr>
</thead>
<tbody>
<tr>
<td>ME+NN (CDF)</td>
<td>DIL, 2 fb$^{-1}$</td>
<td>171.2 ± 2.7(stat) ± 2.9(sys) GeV/c$^2$</td>
</tr>
<tr>
<td>TMF+NN (CDF)</td>
<td>Had., 1.9 fb$^{-1}$</td>
<td>177.0 ± 3.7(stat + JES) ± 1.6(sys) GeV/c$^2$</td>
</tr>
<tr>
<td>ME+NN (CDF)</td>
<td>L3, 1.9 fb$^{-1}$</td>
<td>172.7 ± 1.2(stat) ± 1.3(JES) ± 1.2(sys) GeV/c$^2$</td>
</tr>
<tr>
<td>MW (D0)</td>
<td>DIL, 1 fb$^{-1}$</td>
<td>175.2 ± 6.1(stat) ± 3.4(sys) GeV/c$^2$</td>
</tr>
<tr>
<td>NW (D0)</td>
<td>DIL, 1 fb$^{-1}$</td>
<td>172.5 ± 5.8(stat) ± 3.5(sys) GeV/c$^2$</td>
</tr>
</tbody>
</table>

can be formed for a given top mass. In DIL, this probability can be expressed as

$$P(x, M_t) = \frac{1}{N} \int d\Phi_s |M_{\tilde{t}}(p; M_t)|^2 \prod_{jets} W(p, j) f_{PDF}(q_1) f_{PDF}(q_2),$$

where $M_t$ is the top mass, $x$ contains the lepton and jet energy measurement, $M_{\tilde{t}}(p; M_t)$ is the $\tilde{t}$ production matrix, $q$ is the vector of incoming parton-level quantities, $p$ is the vector of resulting parton-level quantities: lepton and quark momenta, $W(p, j)$ is the transfer function which gives the probability to observe a jet with energy $j$ given a parton energy $p$, and finally, $f_{PDF}$ the parton distribution functions of the two quarks from the proton and anti-proton. The integral is over the entire six-particle phase space. Scanning through the top mass, the most probable point reveals the mass of the top quark. An example of applying such method for top mass measurement is performed at CDF using DIL samples.

This method “Matrix Weighting” is different from the “Matrix Element” method described previously. This method is applied to DIL samples, where the system is under-constrained due to missing neutrinos. For a given top mass, one could try to resolve for $\tilde{t}$ momentum. A weight is calculated for each solution found by comparing the missing energy calculated with the one observed in observed events. The top mass is determined from a scan through a range of top mass to find a maximum weight and the extremum of likelihood. This is described in D0’s public conference note.

“Neutrino weighting” is a method applied in D0. Using DIL samples, for a given top mass $\eta$ was thrown based on Monte Carlo simulation for each $\nu$. Then the set of energy-momentum conservation equations can be resolved for $\nu$ momentums. For each event a weight template was derived based on missing energy expected and observed at each given top mass. A maximum likelihood is formed, combining all events, and the extremum of this distribution reveals the top mass. This is described in D0’s public conference note.

At the Tevatron, many techniques have been developed to measure the top mass. Progress has been made to improve the uncertainty. Some of the methods have not been mentioned in this presentation. A single variable that has a distribution being sensitive to the top mass can be used to do the measurement. One such variable is the Lxy, which is the closest distance of the secondary vertex to the primary vertex in the transverse plan of the detector. may have a distribution which is sensitive to the top mass. The top mass measurement from the top production cross section is discussed by Marc Besancon at this Conference. All of the methods provide additional info, and could help in improving uncertainty of the combined top mass.

2.2 The results

The results on the top mass measurement at Tevatron given at this conference are listed in Table 1. All individual top mass measurement from all three channels show consistent results. There is no indication of seeing different particles in different channels.
At the moment of this presentation, CDF has already a combined result using various results from all hadronic, DIL and LJ channels. This yields $172.9 \pm 1.2^{(\text{stat})} \pm 1.5^{(\text{sys})}$ GeV/c² and is shown in Figure 1. This result is approaching an uncertainty of 1% of the top mass, which is similar to CDF and D0 combined result for the year 2007. Together with the updated D0 measurement, the new combined result would have an uncertainty below 1%. (This happened right after the Moriond EW 2008 conference.) CDF and D0 are working together on common systematic issues to improve uncertainty at the Tevatron for the high precision era of top mass measurement.

3 Top property studies

The SM top quark has spin (1/2), charge (+2/3), and other definite properties which should be measured. In contrast, the top mass is a free parameter in SM. Any significant deviation of the top quark properties from SM would indicate new physics. The top charge is among the fundamental properties of the top quark most accessible at Tevatron. Other properties, such as top spin, lifetime, decay width, either need significantly more data or are far beyond our capability to measure with our given detector resolution. Studies from the top decay include the helicity of W boson from top decay, measurement of branching ratio, search for charged Higgs, search for flavor changing neutral current, etc.

3.1 The charge of top quark

In the SM, the charge of top quark is $+2/3$. An alternative possibility suggested by an exotic model (XM) is $-4/3$. In this model, it is claimed that the particle seen at Tevatron may be an exotic top of charge $-4/3$, which decays into $W^-$ and $b$, unlike in the SM where the top quark
decays into $W^+$ and $b$. The two key elements in the study are then to identify the source of a jet being $b$ or $\bar{b}$, and how the $b$ and $\bar{b}$ jets are paired with the two $W$'s.

The identification of a jet being from $b$ or $\bar{b}$ is done via calculating the jet charge, which is sum of jet-track charges weighted by the track momentum amplitude and how close the track is to the jet axis. For true $b$ jets this method has 60% probability of identifying $b$ or $\bar{b}$ correctly.

The pairing can be done by taking the measured top mass as input and check which pairing is more probable. In DIL channel, events can be selected based on the square of invariant mass of the paired lepton and jet $m_{b2}$ to improve the purity. In each event there are two possible ways of pairing and four possible $m_{b2}$ values. The pairing having the maximum $m_{b2}$ does not always provide the correct pairing. In the events with the maximum $m_{b2}$ greater than certain value, this method can be almost 100% correct. However cutting too tight would lose too much in statistics. The best point for making such cut is 21000, assuming that SM is true and top mass is 175 GeV/c$^2$. With this selection, 94% of pairing purity can be reached with efficiency of 39%.

The charge of the top quark was first studied by D0. With 0.37 fb$^{-1}$ of data, the result prefers the SM instead of XM. In CDF, the study has been done with data up to 1.5 fb$^{-1}$. The result up to date support SM over XM, and the XM is rejected at 87% confidence level (CL). Combining DIL and LJ, among 225 top or anti-top quark decays 124 decays support SM and 101 support XM. Correcting for purity of the analysis, the measured true fraction of SM over total is 0.87, which based on our sensitivity gives a p value of 0.31. An additional way of showing this is the Bayes Factor (BF), which is defined as $P(N_+|SM)/P(N_+|XM)$, i.e. the probability of observed events happening assuming SM is true over the one of XM. A common way to utilize BF is the quantity $L = 2 * \ln(BF)$. For $L$ in the ranges (0-2), (2-6), (6-10), (>10), the result is uncertain, positive, strongly supporting SM, or very strongly supporting SM, respectively. Our result from CDF data is 12.01, thus very strongly support SM over XM. With more data, we will determine more precisely the top charge.

3.2 $W$ helicity

In the SM, V-A rules the weak decay. The $W$ boson from top quark decay is thus polarized. The SM predicts that the $W$ helicity in this case should have 70% longitudinal ($f_0$) and 30% left-handed ($f_-$). The component of right-handed ($f_+$) is very small, $3.6 \times 10^{-4}$. Significant deviation of $f_+$ would indicate new physics.

The study of $W$ helicity can be performed via looking at the $\cos \theta^*$ distribution, where $\theta^*$ is the angle of the electron or muon in the $W$ rest frame with respect to the anti-direction of top quark in this frame. The analysis can be performed in LJ and DIL channels. In case of LJ the missing energy is assumed to be due to the missing neutrino. Events can be reconstructed using top mass as input and lepton angle in $W$ rest frame can be calculated. The top mass used is generally 175 GeV/c$^2$. In case of DIL there are two missing neutrinos. Using top mass as input one can figure out which jet is paired with lepton and resolve for the neutrino momenta. Lepton angle in the $W$ rest frame can be obtained in this way. CDF does this analysis, using 1.9 fb$^{-1}$ data in the LJ channel. In a 2 dimensional fit where both $f_0$ and $f_+$ are fitted at the same time the result shows $f_0 = 0.65\pm0.19(stat)\pm0.03(sys)$ and $f_+ = -0.03\pm0.07(stat)\pm0.03(sys)$. If $f_0$ is fixed to the SM value CDF obtains $f_+ = -0.04\pm0.04(stat)\pm0.03(sys)$ and sets upper limit for $f_+$ at 0.07 at 95% CL. D0 collaboration does the analysis in both LJ and DIL channels. A 2-D fit of $f_0$ and $f_+$ reveals 0.425$\pm0.16(stat)\pm0.102(sys)$ and 0.119$\pm0.009(stat)\pm0.053(sys)$ respectively. Fixing $f_0$ to the SM value gives $f_+ = -0.002\pm0.047(stat)\pm0.047(sys)$. An upper limit of 0.13 at 95% CL is set.
3.3 Study of $R_b$

A study on the $R_b = Br(t \to Wb)/Br(t \to Wq)$, where $q$ represents all possible quarks allowed in the decay, is performed at D0. $R_b$ is correlated with the top pair production. Noting that D0 does simultaneous fit to both values, using LJ channel from 0.9 fb$^{-1}$ data\textsuperscript{15}. The result is $R_b = 0.97^{+0.07}_{-0.08}(\text{stat + sys})$. A lower limit of $R_b$ is set at 0.79 at 95\% CL. From this a lower limit on $[V_{tb}]$ is set at 0.89 at 95\% CL. From the same fit the resulted production cross section is $\sigma_{tt} = 8.1^{+0.9}_{-0.8}(\text{stat + sys}) \pm 0.50(\text{lumi})$ pb, which is consistent with the direct measurement.

3.4 Search for $t \to Hb$

It is interesting to search for charged Higgs in the top quark decay. D0 collaboration did this analysis by comparing the production cross section of top pair from the LJ channel against the one from the DIL channel. If there were charged Higgs in the top decay, it would mostly contribute to the LJ channel but much less in the DIL channel. The ratio of the two production cross sections is $R = 1.21^{+0.12}_{-0.10}(\text{stat + sys})$, based on the assumption that $R_b = 1$. Extracting the branching ratio of $t \to Hb$ from this cross section ratio, D0 obtains $Br = 0.13^{+0.12}_{-0.11}(\text{stat + sys})$. An upper limit is set at 0.35 at 95\% CL\textsuperscript{16}.

3.5 Search for FCNC

At CDF an analysis to study FCNC is to search for $t \to Zq$ in the top quark decay. The SM predicts a branching ratio at the order of $O(10^{-14})$. However beyond SM up to $O(10^{-3})$ is possible. At CDF events having two high $p_t$ leptons with at least 4 jets were selected with constraint on masses of top, $Z$, and W. Comparing the data (1.9 fb$^{-1}$) with expectation, no excess is seen. An upper limit is set at 3.7\% at 95\% CL\textsuperscript{17}.

4 Summary and Future Prospects

The top quark mass has been well measured at the Tevatron, with uncertainty getting below 1\% of the top mass. The top quark mass is an important parameter in the Standard Model, and should be measured as precisely as possible. Other properties of the top quark also should be measured, to learn more about the top quark and study possible new physics. Examples of other top studies at the Tevatron are the charge of the top quark, W helicity, top decay branching ratio to b ($R_b$), searches for $t \to Hb$ and for flavor changing neutral current (FCNC). The results are all consistent with the Standard Model within current statistics. With significantly more data being collected at the Tevatron, precision measurements of the top properties are just starting.

Acknowledgments

Many thanks to the people working on Tevatron. With their dedicated effort to improve the accelerator discoveries are made possible. Many thanks to the funding agencies to CDF and D0 collaboration. Many thanks to the authors and conveners of both collaboration who provided critical input to this presentation. At the end I want to thank the organizers of the Moriond EW 2008 conference for their warm hospitality and well organized program.

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Single top quark production at the Tevatron

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\textit{Michigan State University, 3234BPS, East Lansing, MI 48824, USA}

The Tevatron experiments D0 and CDF have found evidence for single top quark production, based on datasets between 0.9 fb\(^{-1}\) and 2.2 fb\(^{-1}\). Several different multivariate techniques are used to extract the single top quark signal out of the large backgrounds. The cross section measurements are also used to provide the first direct measurement of the CKM matrix element \(|V_{tb}|\).

1 Introduction

Evidence for single top quark production at the Tevatron and a first direct measurement of the CKM matrix element \(|V_{tb}|\) was first reported by the D0 collaboration\textsuperscript{1}. In contrast to top quark pair production through the strong interaction, which was observed in 1995\textsuperscript{2,3}, single top quarks are produced via the weak interaction. The Feynman diagrams for standard model (SM) s-channel (\(tb\)) and t-channel (\(tqb\)) single top quark production are shown in Fig. 1. There is third production mode, associated production of a top quark and a \(W\) boson, but its cross section is so small that it will not be considered further. The SM cross section for the s-channel process \(p\bar{p}\rightarrow t\bar{b} + X, \bar{t}b + X\) is 0.88 ± 0.14 pb at NLO for \(m_{t_{\text{top}}}=175\) GeV\textsuperscript{4,5}. At the same order and mass, the cross section for the t-channel process \(p\bar{p}\rightarrow tq\bar{b} + X, \bar{t}qb + X\) is 1.98 ± 0.30 pb\textsuperscript{4,6}.

Measuring the single top quark production cross section provides a direct measurement of the CKM matrix element \(|V_{tb}|\). The single top quark final state also allows for studies of the top quark polarization, and it is sensitive to many models of new physics, for example flavor changing neutral currents via the gluon\textsuperscript{7} or heavy new bosons \(W^\prime\) that only couple to quarks\textsuperscript{8}. The s-channel process is also an important background to Higgs searches in the associated production mode, and the advanced analysis techniques used in the single top searches will be applicable to Higgs searches as well.

\textsuperscript{a}On behalf of the D0 and CDF collaborations.
The D0 collaboration has updated two of their analysis methods using a dataset of 0.9 fb$^{-1}$. The updated results, including a combination of different methods are presented below. The CDF collaboration has analyzed a dataset of 2.2 fb$^{-1}$ and significantly improved the sensitivity to single top quark production. These new results are presented below.

2 D0 results

2.1 Event selection

The D0 analysis selects electron+jets and muon+jets events in 0.9 fb$^{-1}$ of data with the following requirements: One high-$p_T$ lepton (electron ($p_T > 15 \text{ GeV}$) or muon ($p_T > 18 \text{ GeV}$)), missing transverse energy $E_T > 15 \text{ GeV}$, and between two and four jets with jet $p_T > 15 \text{ GeV}$ and jet 1 $p_T > 25 \text{ GeV}$, at least one is tagged with a neural-network based b-tagging algorithm. Additional cuts remove fake-lepton background events. Events are collected by lepton+jets trigger requirements.

The number of events observed in data and expected from the background model and SM signal is shown in Table 1. The largest sources of systematic uncertainty are the background normalization, jet energy scale, as well as b-tag and trigger modelling.

Table 1: Numbers of events expected by D0 in 0.9 fb$^{-1}$ for electron and muon, 1 b-tag and 2 b-tag channels combined.

<table>
<thead>
<tr>
<th></th>
<th>2 jets</th>
<th>3 jets</th>
<th>4 jets</th>
</tr>
</thead>
<tbody>
<tr>
<td>s-channel</td>
<td>16±3</td>
<td>7±2</td>
<td>2±1</td>
</tr>
<tr>
<td>t-channel</td>
<td>20±4</td>
<td>12±3</td>
<td>4±1</td>
</tr>
<tr>
<td>tt</td>
<td>59±14</td>
<td>134±32</td>
<td>155±36</td>
</tr>
<tr>
<td>W+jets</td>
<td>531±129</td>
<td>248±64</td>
<td>70±20</td>
</tr>
<tr>
<td>Multijets</td>
<td>96±19</td>
<td>77±15</td>
<td>29±6</td>
</tr>
<tr>
<td>Total background</td>
<td>686±131</td>
<td>460±75</td>
<td>253±42</td>
</tr>
<tr>
<td>Data</td>
<td>697</td>
<td>455</td>
<td>246</td>
</tr>
</tbody>
</table>

Table 1 shows that after selection cuts, the expected SM single top signal is small compared to the background sum, and in fact the signal is significantly smaller than the background uncertainty. Thus, more advanced techniques are required to extract the signal.

2.2 Multivariate techniques

The D0 analysis employs three different multivariate techniques to extract the single top quark signal out of the large backgrounds. The boosted decision tree (BDT) analysis has not changed since the publication of evidence for single top quark production\cite{7}. Here we focus on the Bayesian neutral network analysis and the matrix element analysis, both of which have been re-optimized.
In a conventional neural network, the network parameters and weights are determined in an optimization (training) procedure. Rather than optimizing for these network parameters once and then fixing them, the optimal network configuration can be obtained as an average over many different values for the network parameters. In this Bayesian procedure, an integration over all of the possible network parameter space is performed. The network architecture is fixed, and the weight of each set of parameters is obtained through a Bayesian integration. The final network discriminant is then the weighted average over all the individual networks. Fig. 2 shows the output of the BNN for the D0 data.

The Matrix element analysis starts from the Feynman diagrams for the single top quark processes and uses transfer functions to relate the parton level quark-level information to the reconstructed jet and other information. Matrix elements for the single top quark signal as well as the $W$+jets backgrounds are included. For 3-jet events, a top pair matrix element is also included. For each event, an integration over the phase space is performed, employing the transfer functions to compute the probability for this particular event to arise from a specific matrix element. A likelihood function is then formed as the ratio of the signal and signal plus background probabilities.

2.3 D0 summary

The cross section is measured as the peak of the Bayesian posterior probability density, shown in Fig. 3 for the ME analysis. The three different methods measure the following cross sections for the sum of s- and t-channel:

$$\sigma^{\text{obs}}(pp \rightarrow tb + X, tqb + X) = 4.9^{+1.4}_{-1.0} \text{ pb} \text{ (DT)}$$
$$= 4.4^{+1.4}_{-1.0} \text{ pb} \text{ (BNN)}$$
$$= 4.8^{+1.4}_{-1.0} \text{ pb} \text{ (ME)}.$$  

The measured cross sections are consistent with each other and above the SM expectation.

The decision tree analysis has also measured the s- and t-channel cross sections separately,

$$\sigma^{\text{obs}}(pp \rightarrow tb + X) = 1.0 \pm 0.9 \text{ pb}$$
$$\sigma^{\text{obs}}(pp \rightarrow tqb + X) = 4.2^{+1.8}_{-1.4} \text{ pb},$$

where the standard model cross section is used for the single top process not being measured.

Removing the constraint of the standard model ratio allows to form the posterior probability density as a function of both the $tb$ and $tqb$ cross sections. This model-independent posterior is shown in Fig. 3 (right) for the DT analysis, using the $tb+tqb$ discriminant. The most probable value corresponds to cross sections of $\sigma(tb) = 0.9 \text{ pb}$ and $\sigma(tqb) = 3.8 \text{ pb}$. Also shown are...
the one, two, and three standard deviation contours. While this result favors a higher value for the \( t \)-channel contribution than the SM expectation, the difference is not statistically significant. Several updated models of new physics that are also consistent with this result are shown in Ref.\(^9\). These updated results have recently been published\(^{10}\).

## 3 CDF results

### 3.1 Event selection

The CDF analysis selects electron+jets and muon+jets events in 2.2 fb\(^{-1}\) of data with the following requirements: One high-\( p_T \) lepton \((p_T > 20 \text{ GeV})\), \( E_T > 25 \text{ GeV} \), and two or three jets with jet \( p_T > 20 \text{ GeV} \), at least one of which is tagged by a displaced vertex tagging algorithm. Additional cuts remove fake-lepton background events. Events are collected by single-lepton trigger requirements. The matrix element analysis uses additional triggers in the muon channel to increase the acceptance.

The number of events observed in data and expected from the background model and SM signal is shown in Table 2. The largest sources of systematic uncertainty are the background normalization, jet energy scale, and b-tag modelling. Again, it is clear that a advanced analysis techniques are required to extract the signal.

### Table 2: Numbers of events expected by CDF in 2.2 fb\(^{-1}\) for electron and muon, 1 b-tag and 2 b-tag channels combined.

<table>
<thead>
<tr>
<th>Source</th>
<th>2 jets</th>
<th>3 jets</th>
</tr>
</thead>
<tbody>
<tr>
<td>s-channel</td>
<td>41±6</td>
<td>14±2</td>
</tr>
<tr>
<td>t-channel</td>
<td>62±9</td>
<td>18±3</td>
</tr>
<tr>
<td>( \bar{t}t )</td>
<td>146±21</td>
<td>339±48</td>
</tr>
<tr>
<td>W+bottom</td>
<td>462±139</td>
<td>141±43</td>
</tr>
<tr>
<td>W+charm</td>
<td>395±122</td>
<td>109±34</td>
</tr>
<tr>
<td>W+light</td>
<td>340±56</td>
<td>102±17</td>
</tr>
<tr>
<td>Z+jets</td>
<td>27±4</td>
<td>11±2</td>
</tr>
<tr>
<td>diboson</td>
<td>63±6</td>
<td>22±2</td>
</tr>
<tr>
<td>Multijets</td>
<td>60±24</td>
<td>21±9</td>
</tr>
<tr>
<td>Total background</td>
<td>1492±269</td>
<td>755±91</td>
</tr>
<tr>
<td>Data</td>
<td>1535</td>
<td>752</td>
</tr>
</tbody>
</table>
3.2 CDF Likelihood Function

A multivariate likelihood is built from several kinematic variables that each separate the single top quark signal from the backgrounds. One special variable is a specially developed b-tagging neural network that aids in separating b-quark jets from light quark and c-quark jets. An additional special variable is a kinematic solver using constraints from the W boson mass and the top quark mass to determine if an event is well reconstructed. Another special variable is the t-channel matrix element, which uses the kinematic information provided by the kinematic solver. The likelihood discriminant for the t-channel likelihood is shown in Fig. 4 (left).

The measured cross section is obtained as the peak of a Bayesian posterior probability. The likelihood analysis measures a cross section of $\sigma(tb+tqb) = 1.8^{+0.9}_{-0.8}$ pb, below the SM expectation.

3.3 CDF Neural Network

Several kinematic variables as well as the b-tagging neural network output are combined in a neural network. Four different networks are built with 10-14 variables each, trained separately for 2-jet and 3-jet as well as 1-tag and 2-tag events. The full neural network output distribution is shown in Fig. 4 (center), and the signal region is shown in Fig. 4 (right). The neural network analysis measures a cross section of $\sigma(tb + tqb) = 2.0^{+0.9}_{-0.8}$ pb, below the SM expectation but consistent with the SM within uncertainties.

3.4 CDF Matrix Element

The matrix element analysis uses the same approach as described above, but also includes a top pair matrix element in the 2-jet bin. The matrix element for top quark pair events has more final state particles than the single top process, and these additional particles have to be integrated out. This is done by integrating over the kinematics of the hadronically decaying W-boson in a lepton+jets top pair event.

The Bayesian posterior probability density for the Matrix element analysis is shown in Fig. 5, showing the measured cross section and the measurement uncertainty. The measured cross section is $\sigma(tb + tqb) = 2.2^{+0.9}_{-0.7}$ pb, again below the SM expectation but consistent with the SM within uncertainties. The CKM matrix element $|V_{tb}|$ is also extracted from the posterior probability and a lower limit is found to be $|V_{tb}| > 0.59$ at the 95% confidence level.
4 Summary

Both Tevatron experiments have found better than 3 sigma evidence for single top quark production and have made the first direct measurement of the CKM matrix element $|V_{tb}|$ using advanced multivariate techniques. The CKM matrix element $|V_{tb}|$ can be measured to better than 15%. Further improvements to the analyses are in progress and both experiments are working towards observation of single top quark production at the 5 sigma level.

Acknowledgments

We thank the Fermilab staff and the technical staffs of the participating institutions for their vital contributions.

References

B, C and S physics
MEASUREMENTS OF $\phi_1$ AND $\phi_2$ BY BELLE AND BABAR

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We report recent measurements of the Unitarity triangle angles $\phi_1$ and $\phi_2$ using large data samples collected with Belle and BaBar detectors at $e^+e^-$ asymmetric-energy colliders.

1 Introduction

In the Standard Model (SM), $CP$ violation in $B^0$ meson decays originates from an irreducible complex phase in the $3 \times 3$ Cabibbo-Kobayashi-Maskawa (CKM) mixing matrix$^1$. The angles $\phi_1$ and $\phi_2$ of the CKM unitarity triangle have been measured in several $B$ decay modes$^{2,3,4,5}$. Extra studies in different decay modes are important to check the self-consistency between measurements to probe the existence of New Physics.

The results reported in this paper were obtained by two experiments, Belle and BaBar, working at $e^+e^-$ asymmetric-energy colliders, KEKB$^6$ and PEP-II, correspondingly, with the center-of-mass (CM) energy at $\Upsilon(4S)$ resonance ($\sqrt{s} = 10.58 \text{ GeV}$). The Belle detector$^7$ is a large-solid-angle magnetic spectrometer that consists of a silicon vertex detector (SVD), a 50-layer central drift chamber (CDC), a mosaic of aerogel threshold Cherenkov counters (ACC), time-of-flight scintillation counters (TOF), and an array of CsI(Tl) crystals (ECL) located inside a superconducting solenoid coil that provides a 1.5 T magnetic field. An iron flux-return located outside of the coil is instrumented to detect $K_L$ mesons and to identify muons (KLM). For the results from Belle experiment the data sample of 657 million $B\bar{B}$ pairs is used.

The BaBar detector is described in detail elsewhere$^8$. Charged particle momenta are measured with a tracking system consisting of a five-layer silicon vertex tracker (SVT) and a 40-layer $e^-i$: aushev@itep.ru
drift chamber (DCH) surrounded by a 1.5 T solenoidal magnet. An electromagnetic calorimeter (EMC) comprising 6580 CsI(Tl) crystals is used to measure photon energies and positions. Charged hadrons are identified with a detector of internally reflected Cherenkov light (DIRC) and ionization measurements in the tracking detectors. The results from BaBar experiment are based on 383 million $B\bar{B}$ pairs data sample.

2 Study of $B^+ \to D^+ \bar{D}^0$ and search for $B^0 \to D^0 \bar{D}^0$

Recently, evidence of direct $CP$ violation in the decay $B^0 \to D^+D^-$ was observed by Belle\textsuperscript{9}, while BaBar measured an asymmetry consistent with zero\textsuperscript{10}. A similar effect might occur in the charged mode $B^+ \to D^+ \bar{D}^0$\textsuperscript{11}. This decay has already been observed by Belle\textsuperscript{12} and confirmed by BaBar\textsuperscript{13}.

Now, Belle updated their result with larger data sample\textsuperscript{14}. 366 ± 32 events were found from the fit to the $\Delta E - M_{bc}$ distribution (Fig. 1(a,b)), where $\Delta E = E_B - E_{\text{beam}}$, $M_{bc} = \sqrt{E_{\text{beam}}^2 - p_B^2}$, $E_B(p_B)$ is the energy (momentum) of $B$ candidate in the CM system, $E_{\text{beam}}$ is the CM beam energy. The branching fraction of $B^+ \to D^+ \bar{D}^0$ is measured to be $B(B^+ \to D^+ \bar{D}^0) = (3.85 \pm 0.31 \pm 0.38) \times 10^{-4}$, where the first error is statistical and the second one is systematic. The charge asymmetry for this decay is measured to be consistent with zero: $A_{CP}(B^+ \to D^+ \bar{D}^0) = 0.00 \pm 0.08 \pm 0.02$. Belle also searched for the decay $B^0 \to D^0 \bar{D}^0$. An upper limit is established for the branching fraction: $B(B^0 \to D^0 \bar{D}^0) < 0.43 \times 10^{-4}$ (Fig. 1(c,d)).

![Figure 1](image_url)

Figure 1: $\Delta E$ (a,c) and $M_{bc}$ (b,d) distributions for the $B^+ \to D^+ \bar{D}^0$ (a,b) and $B^0 \to D^0 \bar{D}^0$ (c,d) candidates. Each distribution is the projection of the signal region of the other parameter. Points with errors represent the experimental data, open curves show projections from the 2D fits and cross-hatched curves show the $B\bar{B}$ background component only.
3 Study of $B^0 \to D^{*+}D^{*-}$

Another interesting decay mode to study the $CP$ asymmetry is $B^0 \to D^{*+}D^{*-}$. Both experiments have updated their results for this decay mode and obtained high statistics signals shown in Fig. 2(a,c)\(^{15}\). The time-dependent decay rates of $B^0$ and $\bar{B}^0$ to a $CP$ eigenstate, like $D^{*+}D^{*-}$, is given by formula:

$$P(\Delta t) = \frac{e^{-\Delta t/\tau_{B^0}}}{4\tau_{B^0}}\left\{1 + q\left[S_{FCP} \sin(\Delta m_d \Delta t_{B^0}) + A_{FCP} \cos(\Delta m_d \Delta t_{B^0})\right]\right\},$$

where $q$ is the $b$-flavor charge: $q = +1(-1)$ when the tagging $B$ meson is a $B^0$ ($\bar{B}^0$), $\tau_{B^0}$ is the neutral $B$ lifetime, $\Delta m_d$ is the mass difference between two $B^0$ mass eigenstates, $\Delta t_{B^0} = t_{CP} - t_{tag}$. The tree diagram dominates in this decay mode, which according to the SM gives $S_{FCP} = \xi_{D^{*+}D^{*-}} \sin 2\phi_1$ and $A_{FCP} = 0$. The parameter $\xi_{D^{*+}D^{*-}}$ is the $CP$ eigenvalue of the $D^{*+}D^{*-}$, which is +1 when the decay proceeds via $S$ and $D$ waves, or −1 for a $P$ wave. Therefore, the $CP$ measurement requires helicity study to obtain the $CP$-odd fraction $R_{odd}$ of the decay. It is done in both analyses from Belle and BaBar in the so-called transversity basis. The fit results are presented in Fig. 2(b,d). The parameter $R_{odd}$ is found to be equal to $0.143 \pm 0.034$(stat) $\pm 0.008$(syst) by BaBar and $0.116 \pm 0.042$(stat) $\pm 0.004$(syst) by Belle.

![Figure 2](image-url)

Figure 2: Measured distributions of $M_{bc}$ (a,c) and $\cos\theta_{tr}$ in the region $M_{bc} > 5.27$ GeV/$c^2$ (b,d) for BaBar (a,b) and Belle (c,d) of $B^0 \to D^{*+}D^{*-}$ events. The solid lines are the projections of the fit results. The dotted lines represent the background components.

Finally, the unbinned maximum likelihood fit was performed to obtain the $CP$-violating parameters. The results of the fits are summarized in Table 1 and presented in Fig. 3. Both experiments obtained the results well consistent with each other in both the $CP$-odd fraction and the $CP$-violating parameters. Note that in the BaBar parametrization $A = -C$. The Belle results are preliminary.
Table 1: Results for $B^0 \to D^{+} D^{-}$ decay mode.

<table>
<thead>
<tr>
<th></th>
<th>Yield</th>
<th>$R_{old}$</th>
<th>$A = -C$</th>
<th>$S$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Belle</td>
<td>545 ± 29</td>
<td>0.116 ± 0.042 ± 0.004</td>
<td>+0.16 ± 0.13 ± 0.02</td>
<td>−0.93 ± 0.24 ± 0.15</td>
</tr>
<tr>
<td>BaBar</td>
<td>638 ± 38</td>
<td>0.143 ± 0.034 ± 0.008</td>
<td>+0.02 ± 0.11 ± 0.02</td>
<td>−0.66 ± 0.19 ± 0.04</td>
</tr>
</tbody>
</table>

Figure 3: The $\Delta t$ distributions of $B^0 \to D^{+} D^-$ events in the region $M_{bb} > 5.27$ GeV/c$^2$ for $B^0 (\bar{B}^0)$ tagged candidates (a, c) and the raw asymmetry ($N_{ obscene} - N_{neut})/(N_{ obscene} + N_{neut}$), as a function of $\Delta t$ (b, d) for BaBar (a, b) and Belle (c, d). The lines represent the fit results.

4 CP-violation in $B^0 \to K_S\pi^0\pi^0$ and $B^0 \to K_S\pi^0$

In the SM, the CP violation parameters in $b \to s$ “penguin” and $b \to c$ “tree” transitions are predicted to be the same, $S_f \approx -\xi_f \sin 2\phi_1$ and $A_f \approx 0$, with small theoretical uncertainties. Recent measurements however, indicate that the effective $\sin 2\phi_1$ value, $\sin 2\phi_{\text{eff}}$, measured with penguin processes is different from $\sin 2\phi_1 = 0.687 \pm 0.025$ measured in tree decays by $2.6\sigma$. New particles in loop diagrams may shift the weak phase.

Recently, Belle and BaBar measured the CP asymmetry in $B^0 \to K_S^0\pi^0\pi^0$ and $B^0 \to K_S\pi^0$ decays that proceed through $b \to s q\bar{q} (q = u, d)$ transitions. The results of CP-violating parameters measurements are presented in Table 2. Both experiments are perfectly consistent with each other. In the case of $B^0 \to K_S^0\pi^0\pi^0$ the central value of $S$ has a sign opposite to what we expect from the SM, but the errors are still too large to claim the contradiction. The estimated deviation of the average value from the SM is more than $2\sigma$. The fit to the data for Belle for $B^0 \to K_S^0\pi^0\pi^0$ is presented in Fig. 4(a-c) and the BaBar result for $B^0 \to K_S^0\pi^0$ is shown in Fig. 4(d-f).

Table 2: Results for $B^0 \to K_S^0\pi^0\pi^0$ and $B^0 \to K_S^0\pi^0$ decay modes.

<table>
<thead>
<tr>
<th></th>
<th>$A = -C$</th>
<th>$S = -\sin 2\phi_1$</th>
</tr>
</thead>
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<tr>
<td>$B^0 \to K_S^0\pi^0\pi^0$</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Belle</td>
<td>−0.17 ± 0.24 ± 0.06</td>
<td>+0.43 ± 0.49 ± 0.09</td>
</tr>
<tr>
<td>BaBar</td>
<td>−0.23 ± 0.52 ± 0.13</td>
<td>+0.72 ± 0.71 ± 0.08</td>
</tr>
<tr>
<td>Average</td>
<td>−0.18 ± 0.22</td>
<td>+0.52 ± 0.41</td>
</tr>
<tr>
<td>$B^0 \to K_S^0\pi^0$</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Belle</td>
<td>−0.05 ± 0.14 ± 0.05</td>
<td>+0.33 ± 0.35 ± 0.08</td>
</tr>
<tr>
<td>BaBar</td>
<td>−0.24 ± 0.15 ± 0.03</td>
<td>+0.40 ± 0.23 ± 0.03</td>
</tr>
</tbody>
</table>
Figure 4: Distributions for the $M_{bc}$ (a), the $\Delta t$ (b) and the raw asymmetry (c) for $B^0 \rightarrow K^0_S \pi^0 \pi^0$ decay mode from Belle. $\Delta t$ distributions for the $B^0$ (d) and $\bar{B}^0$ (e) tagged events and the raw asymmetry (f) for the decay $B^0 \rightarrow K^0_S \pi^0$ from BaBar. The lines represent the fit result.

5 $\phi_2$ measurements

The CKM angle $\phi_2$ have been measured in decay modes like $B^0 \rightarrow \pi \pi, \rho \rho, \rho \pi$. Addition of new decay modes allows to improve an accuracy of $\phi_2$ measurement and to check a consistency of measurements in different final states. The decay $B^0 \rightarrow a_1^\pm (1260) \pi^\mp$ proceeds through $b \rightarrow u$ transitions, hence its time-dependent $CP$ violation is also sensitive to $\phi_2$. Belle measured the branching fraction for this decay mode to be $B(B^0 \rightarrow a_1^\pm (1260) \pi^\mp)B(a_1^\pm (1260) \rightarrow \pi^\pm \pi^\mp \pi^\mp) = (14.9 \pm 1.6 \pm 2.3) \times 10^{-6}$, while BaBar has updated their previous measurements now with $CP$ violation study: $\mathcal{A}_{CP} = -0.07 \pm 0.07 \pm 0.02$ and $S = +0.37 \pm 0.21 \pm 0.07$. The angle $\phi_2$ was measured to be $\phi_2^{\text{eff}} = 78.6^\circ \pm 7.3^\circ$. The result is presented in Fig. 5(a-c).

Figure 5: $\Delta t$ distributions of the decay $B^0 \rightarrow a_1^\pm \pi^\mp$ for $B^0$ (a) and $\bar{B}^0$ (b) tagged events, and the raw asymmetry (c). The solid lines show the fit results, while the dotted lines show the background component. Projection of the signal region onto (d) $\Delta E$ and (e) $M_{bc}$ for $B^0 \rightarrow \rho \rho$ candidates. The fit result is shown as the thick solid curve; the hatched region represents the signal component. The dotted, dot-dashed and dashed curves represent, respectively, the cumulative background contributions from continuum processes, $b \rightarrow c$ decays, and charmless $B$ decays.

Belle also performed the search for the decay $B^0 \rightarrow \rho^0 \rho^0$ and other decay modes with four pions in the final state. In the absence of the signals, the upper limits on the branching fraction were established. The signal distributions for the $B^0 \rightarrow \rho^0 \rho^0$ are shown in Fig. 5(a,b).
results are preliminary.

Also a number of the decay modes potentially usable for the $\phi_2$ measurements have been studied by BaBar\textsuperscript{24,25,26}. All the results of these studies are summarized in Table 3.

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|c|}
\hline
Mode & Yield & $\epsilon$ (%) & $S(\sigma)$ & $\mathcal{B}(\times 10^{-6})$ & UL($\times 10^{-6}$)\% CL. \\
\hline
\textbf{Belle results} & & & & & \\
$\rho^0\rho^0$ & $24.5_{-22.1}^{+23.8} \pm 9.7$ & 9.16 & 1.0 & $0.4 \pm 0.4^{+0.2}_{-0.2}$ & < 1.0 \\
$\rho_0\pi^+\pi^-$ & $112.5_{-65.6}^{+67.4} \pm 53.9$ & 2.90 & 1.3 & $5.9_{-3.3}^{+3.5} \pm 2.7$ & < 11.9 \\
$4\pi^+$ & $161.2_{-59.4}^{+61.2} \pm 26.0$ & 1.98 & 2.5 & $12.4_{-4.6}^{+4.7} \pm 2.9$ & < 19.0 \\
$\rho^0 f^0$ & $-11.8_{-12.9}^{+14.5} \pm 4.6$ & 5.10 & 2.75 & - & < 0.6 \\
$f^0 f^0$ & $-7.7_{-3.3}^{+4.1} \pm 2.9$ & 2.75 & 2.75 & - & < 0.4 \\
$f^0\pi^+\pi^-$ & $6.3_{-34.7}^{+37.0} \pm 18.0$ & 1.55 & 1.55 & $0.6_{-3.4}^{+2.6} \pm 1.8$ & < 7.3 \\
\hline
\textbf{BaBar results} & & & & & \\
b_{1}^{0}\pi^+ & $178_{-37}^{+39}$ & 6.78 & 4.0 & $6.7 \pm 1.7 \pm 1.0$ & \\
b_{1}^{0}K^+ & $219_{-36}^{+37}$ & 6.73 & 5.3 & $9.1 \pm 1.7 \pm 1.0$ & \\
b_{1}^{0}\pi^\pm & $387_{-39}^{+41}$ & 9.54 & 8.9 & $10.9 \pm 1.2 \pm 0.9$ & \\
b_{1}^{0}K^- & $267_{-32}^{+33}$ & 9.43 & 6.1 & $7.4 \pm 1.0 \pm 1.0$ & \\
\hline
a_{1}^{0}\pi^+ & $382 \pm 79$ & 7.2 & 3.8 & $20.4 \pm 4.7 \pm 3.4$ & \\
a_{1}^{0}K^0 & $241 \pm 32$ & 9.6 & 6.2 & $34.9 \pm 5.0 \pm 4.4$ & \\
a_{1}^{0}\pi^0 & $459 \pm 78$ & 12.5 & 4.2 & $26.4 \pm 5.4 \pm 4.2$ & \\
a_{1}^{0}K^+ & $272 \pm 44$ & 7.9 & 5.1 & $16.3 \pm 2.9 \pm 2.3$ & \\
\hline
\end{tabular}
\caption{Fit results for decays relevant to $\phi_2$ measurements.}
\end{table}

6 \hspace{1cm} \textbf{CP-violation in $\Upsilon(4S)$ decays}

In the decay $\Upsilon(4S) \rightarrow B^0\bar{B}^0 \rightarrow f_1 f_2$, where $f_1$ and $f_2$ are $CP$ eigenstates, the $CP$ eigenvalue of the final state $f_1 f_2$ is $\xi = -\xi_1\xi_2$. Here the minus sign corresponds to odd parity from the angular momentum between $f_1$ and $f_2$. If $f_1$ and $f_2$ have the same $CP$ eigenvalue, i.e. ($\xi_1, \xi_2$) = (+1, +1) or (−1, −1), $\xi$ is equal to −1. Such decays, for example ($f_1, f_2$) = ($J/\psi K^0_S, J/\psi K^0_S$), violate $CP$ conservation since the $\Upsilon(4S)$ meson has $J^{PC} = 1^{--}$ and thus has $\xi_{\Upsilon(4S)} = +1$. The branching fraction within the SM is suppressed by the factor

$$F \approx \frac{x^2}{1 + x^2} (2\sin 2\phi_1)^2 = 0.68 \pm 0.05,$$

where $x = \Delta m_d/\Gamma = 0.776 \pm 0.008$\textsuperscript{27}.

This decay was studied by Belle. Due to a small branching fractions to the final state and low reconstruction efficiencies the expected yield is very small, 0.04 events. In order to increase the signal yield, a partial reconstruction technique was used\textsuperscript{28}. One $B^0$ was fully reconstructed, while only $K_{S}^0$ was reconstructed from another one. The signal was searched in the recoil mass distribution to the reconstructed particles where, in principle, signals from $\eta_c, J/\psi, \chi_{c1}$, or $\psi(2S)$ can be seen. The method was checked using charged $B$ decay control samples, $\Upsilon(4S) \rightarrow B^+B^- \rightarrow (f_{B^+}, J/\psi_{tag}K^-)$ and $\eta_{c}K_{S}^0K_{S}^0$, where $f_{B^+}$ stands for $J/\psi K^+$ and $\bar{B}^0\pi^+$. Also neutral $B$ decays were examined in the decay $\Upsilon(4S) \rightarrow B^0\bar{B}^0 \rightarrow (f_{B^0}, J/\psi_{tag}K_{S}^0$ and $\eta_{c}K_{S}^0K_{S}^0)$ with $f_{B^0} = B^0 \rightarrow D^{(*)-}h^+$. The fit yields 206 ± 57 for charged $B$ and 35 ± 16 for neutral $B$ signal events, which is in good agreement with the MC expectation (Fig. 6(a,b)).

The results of the final fit are shown in Fig. 6(c). The extracted signal yield, $-1.5_{-2.8}^{+3.6}$ events, is consistent with zero as well as with the SM prediction (1.7 events). An upper limit for the branching fraction was obtained $\mathcal{B}(\Upsilon(4S) \rightarrow B^0\bar{B}^0 \rightarrow J/\psi K_{S}^0, (J/\psi, \eta_c)K_{S}^0) < 4 \times 10^{-7}$ at the
90% confidence level, where the SM prediction is $1.4 \times 10^{-7}$. This corresponds to $F < 2$ at the 90% confidence level.

![Figure 6: Recoil mass distributions for samples reconstructed as $\Upsilon(4S) \rightarrow (B^+, (J/\psi, \eta_c)^{ss} K^-)$ (a), $(B^0 \rightarrow D^{(*)-} h^+, (J/\psi, \eta_c)^{ss} K^0_S)$ (b) and $(J/\psi K^0_S, (J/\psi, \eta_c)^{ss} K^0_S)$ (c). The solid lines show the fits to signal plus background distributions while the dashed lines show the background distributions.](image)

7 Summary

The $CP$ violating parameters have been measured in various decay modes. Most of the measurements are in a good agreement with the SM expectations. Although a room for New Physics becomes smaller and smaller, there is still some sign that it can be found in $b \rightarrow s$ transitions. More statistics is necessary to test these possibilities.

References

11. The inclusion of charge conjugate modes is implied throughout this Letter.
Improved Measurement of Inclusive Radiative $B$-meson decays

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University of Melbourne 3010, Australia

We report a fully inclusive measurement of the flavor changing neutral current decay $B \to X_s \gamma$ in the energy range $1.7\text{GeV} \leq E_{c.m.s}^\gamma \leq 2.8\text{GeV}$, covering 97% of the total spectrum, where c.m.s is the center of mass system. Using 605 fb$^{-1}$ of data, we obtain in the rest frame of the $B$-meson

$$
B(B \to X_s \gamma ; E_\gamma^2 > 1.7\text{GeV}) = (3.31 \pm 0.19 \pm 0.37 \pm 0.01) \times 10^{-4},
$$

where the errors are statistical, systematic and from the boost correction needed to transform from the rest frame of the $\Upsilon(4S)$ (c.m.s) to that of the $B$-meson, respectively. We also measure the first and second moments of the photon energy spectrum as functions of various energy thresholds, which extend down to 1.7 GeV. The results are preliminary.

1 Introduction

Radiative $B$-meson decays may offer a view of phenomena beyond the Standard Model of particle physics (SM). In the SM, these decays proceed via a flavor changing neutral current (FCNC) decay, which consists of a loop process. Yet to be discovered particles, such as charged Higgs or supersymmetric particles, may be produced virtually in the loop and produce a measurable deviation from the branching fraction predicted by the SM.

The predictions of the branching fraction at order $\alpha_s^2$ (NNLO - next to next to leading order) $(3.15 \pm 0.23) \times 10^{-41}$, $(2.98 \pm 0.26) \times 10^{-42}$ and the average of experiment measured values $(3.55 \pm 0.26) \times 10^{-43}$ are in tacit agreement. An updated experimental measurement would further test this agreement, and, moreover, give stronger constraints on extensions to the SM e.g. Minimal Supersymmetric Standard Model and left-right symmetric models. The photon energy spectrum is also of great importance. At the parton level, the photon is monochromatic with energy $E \approx m_b/2$ in the $b$-quark rest frame. The energy is smeared by the motion of the $b$-quark inside the $B$ meson and gluon emission. A measurement of the moments of this spectrum allows for a determination of the $b$-quark mass and of its Fermi motion. This information can then be used in the extraction of the CKM matrix elements $|V_{cb}|$ and $|V_{ub}|$ from inclusive semileptonic $B$ decays. A measurement of the low-energy tail of the photon spectrum is important in this context.

Belle has previously measured the $B \to X_s \gamma$ branching fraction with 5.8 fb$^{-1}$ and 140 fb$^{-1}$ of data using semi-inclusive$^9$ and fully inclusive approaches$^{10}$, respectively. Other measurements include those from CLEO$^{11}$ and BaBar$^{12,13,14}$.

Here we present an update of our fully inclusive measurement$^{10}$, based on a much larger dataset and with significant refinements, which includes an unfolding of detector effects on the measured spectrum that improve the measurements of the branching fraction and spectral moments, respectively. We also extend the photon energy range to $E_{c.m.s}^\gamma > 1.7\text{GeV}$, covering more of the spectrum than ever before, where c.m.s refers to the centre of mass system, which is equivalent to the rest frame of the $\Upsilon(4S)$. 

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2 Detector and Data sample

The $B \to X_s \gamma$ decay is studied using the Belle detector at the KEKB asymmetric $e^+e^-$ storage ring\textsuperscript{15}. The data consists of a sample of 604.6 fb\textsuperscript{−1} taken at the $\Upsilon(4S)$ resonance corresponding to $(656.7 \pm 8.9) \times 10^{6}$ $BB$ pairs. Another 68.3 fb\textsuperscript{−1} sample has been taken at an energy 60 MeV below the resonance and is used to measure the non-$BB$ background. Throughout this manuscript, we refer to these data samples as the ON and OFF samples, respectively.

The Belle detector is a large-solid-angle magnetic spectrometer described in detail elsewhere\textsuperscript{16}. The main component relevant for this analysis is the electromagnetic calorimeter (ECL) made of 16.2 radiation lengths long CsI(Tl) crystals. The photon energy resolution is about 2\% for the energy range relevant in this analysis.

3 Analysis Strategy

The strategy to extract the signal $B \to X_s \gamma$ spectrum is to collect all high-energy photons, vetoing those originating from $\pi^0$ and $\eta$ decays to two photons. The contribution from continuum $e^+e^- \to q\bar{q}$ ($q = u, d, s, c$) and QED type events is subtracted using the OFF sample. The remaining backgrounds from $BB$ events are subtracted using Monte Carlo (MC) distributions scaled by data control samples.

Photon candidates are selected from ECL clusters of 5 $\times$ 5 crystals in the barrel region ($-0.35 \leq \cos \theta \leq 0.70$, where $\theta$ is the polar angle with respect to the beam axis, subtended from the direction opposite the positron beam. They are required to have an energy $E^\gamma_{\text{cluster}}$ larger than 1.4 GeV. We require 95\% of the energy to be deposited in the central 3 $\times$ 3 crystal array and use isolation cuts to veto photons from bremsstrahlung and interaction with matter. The center of the cluster has to be displaced from any other ECL cluster with $E > 20$ MeV by at least 30 cm at the surface of the calorimeter, and from any reconstructed track by 30 cm, or by 50 cm for tracks with a measured momentum above 1 GeV/c. Moreover, the angle between the photon and the highest energy lepton in the event has to be larger than 0.3 radians at the interaction point.

In the Belle detector, a non-negligible background (1\%) is due to the overlap of a hadronic event with energy deposits left in the calorimeter by previous QED interactions (mainly Bhabha scattering). Such composite events are completely removed using timing information for calorimeter clusters associated with the candidate photons. The cluster timing information is stored in the raw data, and is available in the reduced format used for analysis only for data processed after the summer of 2004. This divides our data set into 253.7 fb\textsuperscript{−1} and 350.9 fb\textsuperscript{−1} samples of reprocessed data without and with timing information, respectively. To minimise composite background due to Bhabha scattering and two-photon processes that contaminate both $\Upsilon(4S)$ and continuum data samples, we veto any candidate that contains an ECL cluster with energy exceeding 1 GeV within a cone of 0.2 radians in the direction opposite our photon candidate as measured in the c.m.s frame. In the second data set only photons that are in time with the rest of the event are retained. The efficiency of this selection on signal events is larger than 99.5\%. We veto candidate photons from $\pi^0$ and $\eta$ decays to two photons by combining each $B \to X_s \gamma$ candidate photon with all other photons in the event. We reject the photon candidate if the likelihood of being a $\pi^0$ or $\eta$ is larger than 0.1 and 0.2, respectively, this yield, on average, background suppression factors of 4 and 2, respectively. These likelihoods are determined from MC and are functions of the laboratory energy of the other photon, its polar angle $\theta$ and the mass of the two-photon system.

In order to reduce the contribution from continuum events, we use two Fisher discriminants calculated in the c.m.s frame. The first discriminant exploits the topology of $B \to X_s \gamma$ events and combines three energy flows around the photon axis. These energy flow variables are obtained using all particles, except for the photon candidate, we measure the energy in the three regions defined by $\Theta < 30^\circ$, $30^\circ \leq \Theta \leq 140^\circ$, $\Theta > 140^\circ$, where $\Theta$ is the angle of the particle to the candidate photon. The second exploits the spherical shape of $BB$ events and is built using ten event-shape variables including Fox-Wolfram moments\textsuperscript{17} for the full event and for the partial event with the photon removed, the full- and partial-event thrusts and the angles of the thrust axis with respect to the beam and the photon direction. To optimise these selection criteria, we use a MC simulation\textsuperscript{18} containing large samples of $BB$, $q\bar{q}$ and signal weighted according to the luminosities of the ON and OFF samples. In the optimisation step the signal MC used is generated as inclusive $B \to X_s \gamma$ and exclusive $B \to K^+\gamma$. The inclusive component $X_s$ is defined as a resonance of spin-1 with a Breit-Wigner form and a mass of 2.4 GeV/$c^2$ and width 1.5 GeV/$c^2$. The $X_s$ system is hadronised by JETSET and subsequently reweighted to match the prediction of the DGE.
model\textsuperscript{21} with $m_q(M_S) = 4.20\text{ GeV}/c^2$, with the mass extending no lower than 1.18 GeV/c\textsuperscript{2} to agree with the corresponding world average branching fractions\textsuperscript{3}. To improve the understanding of the photon energy spectrum at low energies, the selection criteria are optimised to maximize the sensitivity to the signal in the energy bin $1.8\text{ GeV} < E_\gamma^{c.m.s} < 1.9\text{ GeV}$.

After these selection criteria we observe $4.15 \times 10^6$ and $0.25 \times 10^6$ photon candidates in the ON and OFF data samples, respectively.

4 Background subtraction

The spectrum measured in OFF data is scaled by luminosity to the expected number of non-$B\bar{B}$ events in ON data and subtracted. The formula used to subtract continuum background is as follows:

\begin{equation}
N^{BB}(E_\gamma^{c.m.s(ON)}) = N_{ON}(E_\gamma^{c.m.s(ON)}) - \alpha \cdot \frac{N_{ON}^{Hadronic}}{N_{OFF}^{Hadronic}} \cdot \frac{\epsilon_{B \rightarrow X,\gamma}}{\epsilon_{B \rightarrow X,\gamma}^{OFF}} \cdot F_N \cdot N^{OFF}(F_E E_\gamma^{c.m.s(OFF)})
\end{equation}

where $\epsilon$ is the efficiency of Belle’s hadronic selection\textsuperscript{19} or of this analysis’ ($B \rightarrow X,\gamma$) selection criteria in continuum events at either ON resonance ($\sqrt{s} = 10.58\text{ GeV}$) or OFF resonance ($\sqrt{s} = 10.52\text{ GeV}$) energies, and $\alpha$ is the ratio of ON to OFF resonance integrated luminosity corrected for the energy difference ($\alpha = 8.7557(\pm 0.3\%)$). The factors $F_E$ and $F_N$ compensate for the slightly lower mean energy and multiplicity of particles in OFF compared to ON events. We find $F_N = 1.0009 \pm 0.0001$, $F_E = 1.0036 \pm 0.0001$. $\epsilon_{B \rightarrow X,\gamma}^{ON} = 0.9986 \pm 0.0001$, and $\epsilon_{B \rightarrow X,\gamma}^{ON} = 0.9871 \pm 0.0014$. The ON and scaled OFF spectra and their difference are shown in Fig. 1.

We then subtract the backgrounds from $B$ decays from the obtained spectrum. Six background categories are considered: (i) photons from $\pi^0 \rightarrow \gamma\gamma$; (ii) photons from $\eta \rightarrow \gamma\gamma$; (iii) other real photons (mainly decays of $\omega$, $\eta'$, and $J/\psi$, and bremsstrahlung, including the short distance radiative correction (modelled with PHOTOS\textsuperscript{22}); (iv) ECL clusters not due to single photons (mainly $K_L^0$’s and $\bar{n}$’s); (v) Electrons misidentified as photons and; (vi) beam background. The spectra of the background of photons from $B$-meson decays with respect to the expected signal is shown in Fig. 2, their relative contributions are also listed. The net background of this type is a factor five greater than the signal.

For each of these categories we take the predicted background from MC and scale it according to measured yields wherever possible. The inclusive $B \rightarrow \pi^0 X$ and $B \rightarrow \eta X$ spectra are measured in data using pairs of photons with well-balanced energies and applying the same ON–OFF subtraction procedure. The yields obtained in data are on average 10% larger and 5% lower for $\pi^0$ and $\eta$ than MC expectations. The observed discrepancy between the measured and simulated $\pi^0$ $\eta$ spectra is attributed to the branching fraction assumptions used for the generator\textsuperscript{23}. Beam background is measured using a sample of randomly triggered events and added to the $BB$ MC.

For each selection criterion and each background category we determine the $E_\gamma^{c.m.s}$-dependent selection efficiency in OFF-subtracted ON data and MC using appropriate control samples. We then scale the MC background sample according to the ratio of these efficiencies. The efficiencies of the $\pi^0$ and $\eta$ vetoes for photons not from $\pi^0$ and $\eta$ are measured in data using one photon from a reconstructed $\pi^0$, where the other photon of the $\pi^0$ is excluded from the search over the remaining photons for the next best $\pi^0$ or $\eta$ candidate (highest $\pi^0$ or $\eta$ likelihood). Consequently the best formed $\pi^0$ or $\eta$ candidate used in the calculation of the likelihoods is most likely a random combination, and therefore suited to measuring the effect of the vetoes. The $\pi^0$ veto efficiency is measured using a sample of photons coming from measured $\pi^0$ decays. We use partially reconstructed $D^+ \rightarrow D^0\pi^+$, $D^0 \rightarrow K^-\pi^+\pi^0$ decays where the $\pi^0$ is replaced by the candidate photon in the reconstruction. The $\eta$ veto efficiency for photons from $\pi^0$’s and event-shape criteria efficiencies are measured using a $\pi^0$ anti-veto sample, which is made of photons with a $\pi^0$ likelihood larger than 0.75 (i.e., no $\pi^0$ veto) and passing all other selection criteria. Other efficiencies are measured using the signal sample. Beam background is negligible after the application of the OFF time veto. In the sample of data where the veto is unavailable we scale the background according to a comparison of yields between MC and data for high energy ($E_\gamma^{c.m.s} > 2.8\text{ GeV}$) photon candidates found in the endcaps of the ECL. This sample after continuum subtraction is a clean sample of ECL clusters from beam backgrounds.

\footnote{In the optimisation step the choice of signal model has a negligible effect on the measure of optimisation, suffice to say the choice of signal model should not be construed as preferential.}
The ratios of data and MC efficiencies versus $E_{\gamma}^{\text{c.m.s}}$ are fitted using first or second order polynomials, which are used to scale the background MC. Most are found to be statistically compatible with unity. An example is the effect of the $\pi^0$ veto on photons from $\pi^0$'s that escape the veto in the partially reconstructed $D^*$ sample, which is shown in Fig. 3.

An exception is the efficiency of the requirement that 95% of the energy be deposited in the central nine cells of the $5 \times 5$ cluster, which is found to be poorly modelled by our MC for non-photon backgrounds. We estimate the efficiency for data using a sample of candidate photons in OFF-subtracted ON data after subtracting the known contribution from real photons. This increases the yield of background (iv) by 50%. The yield from the six background categories, after having been properly scaled by the above described procedures, are subtracted from the OFF-subtracted spectrum. The result is shown in Fig. 1. After these subtractions the yield in the spectrum above the endpoint of $B$ decays is compatible with zero, $1245 \pm 4349$ candidates.

5 Correction for Acceptance

To measure the branching fraction and the moments we correct the raw spectrum using a three step procedure: (i) divide by the efficiency of the selection criteria i.e. the probability of a photon candidate passing cuts given a cluster has been found in the ECL, as a function of the measured energy in the c.m.s frame; (ii) perform an unfolding procedure based on the Singular Value Decomposition (SVD) algorithm, which maps the spectrum from measured energy to true energy thereby undoing the distortion caused by the ECL; (iii) divide by the efficiency of detection i.e. the probability that a photon originating at the interaction point is reconstructed in the ECL, as a function of the true energy. Data are divided into 50 MeV wide bins. Step (ii), which was not performed in our previous analysis, is essential for a consistent extraction of partial branching fractions and moments as a function of lower energy thresholds. The unfolding matrix, derived from signal MC, is calibrated to data using the results of a study of radiative di-muon events, which gave the ECL response in data and MC in an energy and acceptance range consistent with our analysis. We use five signal models: KN 25, BLNP 26, 27, DGE 20, BBU 28 and GG 29.

Values of the parameters of the signal model used in the signal MC are derived from fits to the signal spectrum shown in Fig 1. The two error bars for each point show the statistical and the total error, including the systematic error which is correlated among the points. In order to obtain the total $B \rightarrow X_{\gamma} \gamma$ branching fraction we apply corrections for the contribution from Cabibbo suppressed $B \rightarrow X_d \gamma$ decays. The ratio of the $B \rightarrow X_{\gamma} \gamma$ and $B \rightarrow X_d \gamma$ branching fractions is assumed to be $R_{d/s} = (4.0 \pm 0.4)\%$ 30. We apply corrections to derive the measurements in the $B$-meson rest frame, using a toy MC approach. We generate photon 4-momentum in the rest frame of the $B$-meson using signal models referred to earlier, and generate $B$-meson 4-momentum using their known fixed energy and $1 - \cos \theta^2$ distribution in the c.m.s. Repeating this exercise many times yields photon energy spectra in the rest frame of the $B$-meson and the c.m.s, from which we extract corrections used to yield measurements in the $B$-meson frame. The correction is derived as a mean over all signal models while the root-mean-square is assigned as the uncertainty. After correcting for the acceptance we derive distributions of the partial branching fractions, first moment (mean) and second central moment (variance) of $B \rightarrow X_{\gamma} \gamma$ as measured in the c.m.s and $B$ rest frame for lower energy thresholds as shown in Fig. 1. In the range from 1.7 to 2.8 GeV in the rest frame of the $B$-meson, we obtain a partial branching fraction, and the first two moments of the energy spectrum:

$$B (B \rightarrow X_{\gamma} \gamma) = (3.31 \pm 0.19 \pm 0.37 \pm 0.01) \times 10^{-4}$$

$$\langle E_{\gamma} \rangle = 2.281 \pm 0.032 \pm 0.053 \pm 0.002 \text{GeV}$$

$$\langle E_{\gamma}^2 \rangle - \langle E_{\gamma} \rangle^2 = 0.0396 \pm 0.0156 \pm 0.0214 \pm 0.0012 \text{GeV}^2,$$

where the errors are statistical, systematic and from the boost correction, respectively.

6 Results

The full results, the systematic error budget and correlation coefficients for five lower energy thresholds ($E_{\gamma}^B = 1.7, 1.8, 1.9, 2.0, 2.1 \text{ GeV}$) are listed in Table 1. The total systematic error is derived from a sum in quadrature over all sources. We vary the number of $B \bar{B}$, the ON to OFF ratio of integrated luminosities and the correction factors applied to the OFF data photon candidates and assign the observed variation as the systematic associated with continuum subtraction. The parameters of the correction functions...
Figure 1: (1ST ROW-left) ON data (open circle), scaled OFF data (open square) and continuum background subtracted (filled circle) photon energy spectra of candidates in the c.m.s frame. (1ST ROW-right) The extracted photon energy spectrum of $B \rightarrow X_s \gamma$. The two error bars show the statistical and total errors. (2ND ROW) Partial branching fractions, (3RD ROW) mean, and (4TH ROW) variance of $B \rightarrow X_s \gamma$ in the (LEFT) c.m.s and (RIGHT) rest frame of the $B$-meson for lower energy thresholds. The two error bars show the statistical and total errors.
Table 1: The measurements and correlation coefficients of the branching fraction, mean and variance of the photon energy spectrum for various lower energy thresholds, $E_B^{\gamma}$, as measured in the rest frame of the $B$-meson and the contributions to the systematic uncertainty.

<table>
<thead>
<tr>
<th>$E_B^{\gamma}$ (GeV)</th>
<th>$B(B \to X_{\gamma}^{\prime} \gamma)$ ($10^{-4}$)</th>
<th>$(E_{\gamma})$ (GeV)</th>
<th>$\Delta E_{\gamma}^2 = \langle E_{\gamma}^2 \rangle - \langle E_{\gamma} \rangle^2$ (GeV$^2$)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>1.7</td>
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<tr>
<td>±systematic</td>
<td>0.37</td>
<td>0.24</td>
<td>0.16</td>
</tr>
<tr>
<td>±boost</td>
<td>0.01</td>
<td>0.01</td>
<td>0.02</td>
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| Continuum           | 0.18 | 0.11 | 0.08 | 0.07 | 0.07 | 0.030  | 0.016  | 0.008  | 0.004  | 0.002  |
| Selection           | 0.20 | 0.15 | 0.11 | 0.08 | 0.06 | 0.023  | 0.012  | 0.006  | 0.003  | 0.001  |
| $\pi^0/\eta$        | 0.07 | 0.05 | 0.04 | 0.02 | 0.01 | 0.012  | 0.006  | 0.003  | 0.002  | 0.001  |
| Other $B$           | 0.24 | 0.13 | 0.06 | 0.02 | 0.01 | 0.033  | 0.016  | 0.007  | 0.002  | 0.000  |
| Beam                | 0.02 | 0.02 | 0.01 | 0.01 | 0.01 | 0.001  | 0.001  | 0.001  | 0.000  | 0.000  |
| resolution          | 0.01 | 0.01 | 0.02 | 0.02 | 0.03 | 0.006  | 0.005  | 0.005  | 0.004  | 0.004  |
| Unfolding           | 0.01 | 0.00 | 0.00 | 0.01 | 0.01 | 0.002  | 0.001  | 0.001  | 0.001  | 0.002  |
| Model               | 0.03 | 0.02 | 0.01 | 0.00 | 0.00 | 0.005  | 0.003  | 0.002  | 0.001  | 0.000  |
| $\gamma$ Detection | 0.03 | 0.02 | 0.01 | 0.00 | 0.00 | 0.005  | 0.003  | 0.002  | 0.001  | 0.000  |
| $B \to X_{\gamma}^{\prime}$ | 0.01 | 0.01 | 0.01 | 0.01 | 0.01 | 0.000  | 0.000  | 0.000  | 0.000  | 0.000  |

Correlation coefficients (combined statistical and systematic):

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<th>$\Delta B$</th>
<th>$(E_{\gamma})$</th>
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<tr>
<td>1.7</td>
<td>0.959</td>
<td>0.811</td>
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<tr>
<td>1.8</td>
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<td>0.942</td>
</tr>
<tr>
<td>2.0</td>
<td>1.000</td>
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<tr>
<td>2.1</td>
<td>1.000</td>
<td>0.923</td>
</tr>
<tr>
<td>1.7</td>
<td>1.000</td>
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<tr>
<td>1.8</td>
<td>1.000</td>
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<tr>
<td>(E_{\gamma})</td>
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<td>0.954</td>
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<tr>
<td>2.1</td>
<td>1.000</td>
<td>0.954</td>
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<table>
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<tr>
<th>$\Delta E_{\gamma}^2$</th>
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<th>(2.0)</th>
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<tr>
<td>(1.7)</td>
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<td>(1.8)</td>
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Rencontres de Moriond 2008
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<tr>
<td>Decays of $\pi^0$</td>
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</tr>
<tr>
<td>Decays of $\eta$</td>
<td>0.163</td>
</tr>
<tr>
<td>Other secondary $\gamma$</td>
<td>0.081</td>
</tr>
<tr>
<td>Mis-Identified electrons</td>
<td>0.061</td>
</tr>
<tr>
<td>Mis-Identified hadrons</td>
<td>0.017</td>
</tr>
<tr>
<td>Beam background</td>
<td>0.013</td>
</tr>
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</table>

Figure 2: Relative contributions of the $B\bar{B}$ backgrounds after selection in the $1.7 < E_{\gamma}^{c.m.s.}/(\text{GeV}) < 2.8$ range. The spectra of photons from $B$-meson decays passing selection criteria as predicted using a MC sample.

Figure 3: (LEFT) The $\pi^0$ veto efficiency in the partially reconstructed $D^*$ sample for both Data (circles) and MC (squares) and (RIGHT) their ratio fitted with a first order polynomial.

applied to the $\pi^0$ and $\eta$ yields are varied taking into account their correlations. As we do not measure the yields of photons from sources other than $\pi^0$’s and $\eta$’s in $B\bar{B}$ events, we independently vary the expected yields of these additional sources by ±20%. For the model dependence in correcting for the acceptance we use four signal models in addition to the default model, and assign the maximum deviation from the default as the uncertainty. The error on the photon detection efficiency in the ECL is measured to be 2% using radiative $\mu$-pair events, and also affects the estimation of photons from $B\bar{B}$. For the uncertainties related to the unfolding procedure, we vary the effective rank parameter up and down by one in the SVD algorithm.

7 Summary

In conclusion, we have measured the branching fraction and photon energy spectrum of $B \to X_s \gamma$ in the energy range $1.7 \text{ GeV} \leq E_{\gamma}^{c.m.s.} \leq 2.8 \text{ GeV}$ in a fully inclusive way. For the first time 97% of the spectrum is measured allowing the theoretical uncertainties to be reduced to a very low level. Using 605 fb$^{-1}$ of data taken at the $\Upsilon(4S)$ and 68 fb$^{-1}$ taken below the resonance, we obtain $\mathcal{B}(B \to X_s \gamma : E_B^{\gamma} > 1.7 \text{ GeV}) = (3.31 \pm 0.19 \pm 0.37 \pm 0.01) \times 10^{-4}$, where the errors are statistical, systematic and due to the boost correction, respectively. This result is in agreement with the latest theoretical calculations$^{1,2,20}$. The results can be used to place constraints on new physics$^{22}$ and determine SM parameters such as the $b$-quark mass$^{33}$.

Acknowledgments

We thank the KEKB group for excellent operation of the accelerator, the KEK cryogenics group for efficient solenoid operations, and the KEK computer group and the NII for valuable computing and SINET3 network support. We acknowledge support from MEXT and JSPS (Japan); ARC, DEST and A.J. Slocum (Australia); NSFC (China); DST (India); MOEHRD, KOSEF and KRF (Korea); KBN
References

We present selected new results on leptonic $B$ meson decays from the BABAR experiment: searches for the decays $B^0 \to \ell^+\ell^-$, $B^+ \to \ell^+\nu$ and $B^0 \to \ell^+\tau^-$, and $B \to K\nu\pi$, where $\ell = e$ or $\mu$. We observe no evidence for these decays and set upper limits on their branching fractions.

1 Introduction

Leptonic $B$ meson decays provide an important tool to investigate the Standard Model (SM) and physics beyond the SM. They are highly suppressed in the SM, because they involve a $b \to d$ transition, require an internal quark annihilation, and there are also helicity suppression for $B^0 \to \ell^+\ell^-$ and $B^+ \to \ell^+\nu$ modes, and because the flavor-changing neutral-currents are forbidden at the tree level for $B \to K\nu\pi$ mode. The decay rates can be enhanced or reduced when heavy virtual particles like Higgs or super-symmetric (SUSY) particles replace the $W$ boson or show up at higher orders in loop diagrams. Constraints on these decays can provide information on important SM parameters, such as $B$ meson decay constant. They have identifiable final states with low multiplicity, but they are mostly below our sensitivity. These decay modes will play an important role at the future colliders, such as a Super-B factory, ILC, and LHC (for muon modes).

The analyses described in this paper use data recorded with the BABAR detector at the PEP-II asymmetric energy $e^+e^-$ storage rings. A detailed description of the BABAR detector can be found elsewhere. A full BABAR Monte Carlo (MC) simulation using GEANT4 is used to evaluate signal efficiencies and to identify and study background sources.
Table 1: Result of $B^0 \to \ell^+\ell^-$ analysis. Efficiency ($\epsilon$), number of signal events ($N_{\text{sig}}$) from ML fit, and 90% confidence level upper limit on the branching fraction (UL(BF)) for the three leptonic decays $B^0 \to e^+e^-$, $B^0 \to \mu^+\mu^-$, and $B^0 \to e^\pm\mu^\mp$ are shown.

<table>
<thead>
<tr>
<th>Decay</th>
<th>$\epsilon$ (%)</th>
<th>$N_{\text{sig}}$</th>
<th>UL(BF) x 10^{-8}</th>
</tr>
</thead>
<tbody>
<tr>
<td>$B^0 \to e^+e^-$</td>
<td>16.6 ± 0.3</td>
<td>0.6 ± 2.1</td>
<td>11.3</td>
</tr>
<tr>
<td>$B^0 \to \mu^+\mu^-$</td>
<td>15.7 ± 0.2</td>
<td>-4.9 ± 1.4</td>
<td>5.2</td>
</tr>
<tr>
<td>$B^0 \to e^\pm\mu^\mp$</td>
<td>17.1 ± 0.2</td>
<td>1.1 ± 1.8</td>
<td>9.2</td>
</tr>
</tbody>
</table>

2 $B^0 \to \ell^+\ell^-$

The leptonic decays $B^0 \to \ell^+\ell^-$ are studied using $383.6 \times 10^6 B\bar{B}$ events. The SM prediction on the branching fractions (BFs) are $1.9 \times 10^{-15}(8.0 \times 10^{-11})$ for the $e^+e^-(\mu^+\mu^-)$ mode, and the $B^0 \to e^\pm\mu^\mp$ decay is forbidden. The best upper limits (UL) on the BFs have been set at the order of $10^{-8}$ by the $\text{B}\bar{\text{A}}\text{B}\bar{\text{A}}\text{R}^5$ experiment for $e^+e^-$ and $e^\pm\mu^\mp$ modes using $111\,\text{fb}^{-1}$, and by CDF$^6$ experiment for $\mu^+\mu^-$ mode with $2\,\text{fb}^{-1}$.

The $B^0$ candidate is reconstructed by combining two oppositely charged tracks originating from a common vertex. We use two kinematical quantities: $m_{ES} = \sqrt{(E_{\text{beam}}^*)^2 - (\Sigma p_i^*)^2}$ and $\Delta E = \sum_i \sqrt{m_i^2 + (p_i^*)^2} - E_{\text{beam}}^*$, where $E_{\text{beam}}^*$ is the beam energy in the CM frame, $p_i^*$ and $m_i$ are the momenta in the CM frame and the masses of the daughter particles $i$ of $B$ meson. $E_{\text{beam}}^*$ is used instead of the measured $B$ meson energy in the CM frame because $E_{\text{beam}}^*$ is more precisely known. For correctly reconstructed $B^0$ mesons, the $m_{ES}$ distribution has a maximum at the $B^0$ mass with a standard deviation of about 2.5 MeV/c$^2$ and the $\Delta E$ distribution has a maximum near zero with a standard deviation of about 25 MeV.

Stringent requirements on particle identification$^7$ are made to reduce the contamination from misidentified hadrons and leptons. We retain about 93% (73%) of the electrons (muons), with a misidentification rate for pions of less than about 0.1% (3%). The main background are continuum processes where $e^+e^- \to f\bar{f}$, ($f = u, d, s, c, \tau$). A Fisher discriminant$^8$ ($F$) is constructed, using their different event topology with respect to that of the signal events.

A maximum likelihood (ML) fit is performed based on the variables $m_{ES}$, $\Delta E$ and $F$. The results are summarized in Table 1. The event and background $sPlot^9$ distributions are shown in Figure 1. Using a Bayesian approach, a 90% confidence level (CL) UL on the BF is calculated. The systematic uncertainties are included as a Gaussian into the likelihood calculation.

3 $B^+ \to \ell^+\nu$ and $B^0 \to \ell^+\tau^-$

We present searches for the decays $B^+ \to \ell^+\nu$ and the lepton flavor violating decays $B^0 \to \ell^+\tau^-$, where $\ell = e$ or $\mu$ using $378 \times 10^6 B\bar{B}$ events. The SM predictions of the BFs are of the order of $10^{-11}(10^{-7})$ for $B^+ \to e^+\nu$ ($B^+ \to \mu^+\nu$), and $B^0 \to \ell^+\tau^-$ modes are forbidden. The UL on the BFs have been measured by $\text{B}\bar{\text{A}}\text{B}\bar{\text{A}}\text{R}^10$, Belle$^{11}$, and CLEO$^{12}$. The best published limits are from Belle for $B^+ \to \ell^+\nu$, at the order of $10^{-8}$ with $253\,\text{fb}^{-1}$, and CLEO for $B^0 \to \ell^+\tau^-$, at the order of $10^{-4}$ with $9.6 \times 10^6 B\bar{B}$ events.

We fully reconstruct one of the two B mesons ($B_{\text{tag}}$) in the event: $B_{\text{tag}} \to D^{(*)}X_{\text{had}}$, $X_{\text{had}}$ decays in combinations of $K$’s and $\pi$’s. This method has not been used for searches for these modes. To suppress the continuum backgrounds, we use their different event topologies with
Figure 1: The distributions of events in $m_{ES}$ (a,b,c), $\Delta E$ (d,e,f) and $\mathcal{F}$ (g,h,i) for $B^0 \rightarrow e^+e^-$ (left), $B^0 \rightarrow \mu^+\mu^-$ (middle), $B^0 \rightarrow e^\pm\mu^\mp$ (right) are shown. The points with error bars are data. The overlaid solid curve in each plot is the background $sPlot$ distribution obtained by maximizing the likelihood not using the information from the corresponding component. The dotted line, representing the signal probability density function with an arbitrary scaling, indicates where the signal is expected.
The unbinned maximum likelihood fits and the distributions of the lepton momentum for $B^+ \rightarrow \ell^+ \nu$ and $B^0 \rightarrow \ell^+ \tau^-$ analyses. The points with error bars are data, the solid line represents the ML fit. The dashed line, representing the signal probability density function with an arbitrary scaling, indicates where the signal is expected.

respect to that of the signal events. After all selection criteria are applied, it results in a yield of approximately 2500 (2000) correctly reconstructed $B^+$ ($B^0$) candidates per fb$^{-1}$ of data. This hadronic tagging method yields lower statistics than other methods but it provides an almost background-free environment.

All particles not used in the $B_{tag}$ reconstruction are included in the reconstruction of the signal $B$ meson. From the two-body kinematics, we expect a mono-energetic lepton in the signal $B$ rest frame: lepton momentum ($p^*$) of 2.64 (2.34) GeV/c for the $B^+ \rightarrow \ell^+ \nu$ ($B^0 \rightarrow \ell^+ \tau^-$) modes.

We reconstruct $\tau$ in the following modes: $e^- \tau_e \nu_\tau$, $\mu^- \tau_\mu \nu_\tau$, $\pi^- \nu_\tau$, $\pi^- \pi^0 \nu_\tau$, $\pi^- \pi^0 \pi^0 \nu_\tau$, and $\pi^- \pi^- \pi^+ \nu_\tau$. The second highest momentum track in the event excluding the $B_{tag}$ daughters is assumed to be a $\tau$ daughter, and is required to have a charge opposite to the primary signal lepton.

The signal yields are extracted from unbinned ML fits to the signal lepton momentum distributions, as measured in the signal $B$ rest frame. The fits are restricted to the ranges in $p^*$ shown in Fig. 2. Using a Bayesian approach, a 90% CL UL on the BF is determined. The dominant systematic uncertainties are due to the fitting procedure and the determination of $B_{tag}$ efficiencies. The total uncertainty is between 10 and 16% depending on the modes. The uncertainties are incorporated into the final results by varying the BF assumption by its uncertainty when integrating likelihood for the 90% CL UL. The results are summarized in Table 2.

4 $B \rightarrow K \nu\overline{\nu}$

The $B \rightarrow K \nu\overline{\nu}$ decays are studied using 319 fb$^{-1}$ of data. The SM prediction of this mode is $(3.8\pm1.2)\times10^{-6}$ and the best published UL is at $1.4\times10^{-5}$ from Belle with $535\times10^6 B\overline{B}$ events.
Table 2: Result of $B^+ \to \ell^+\nu$ and $B^0 \to \ell^+\tau^-$ analyses. The efficiency ($\epsilon$), number of signal events ($N_{sig}$) and 90% CL UL on the BF (UL(BF)) for the decay modes are shown.

<table>
<thead>
<tr>
<th>Decay</th>
<th>$\epsilon \times 10^{-5}$</th>
<th>$N_{sig}$</th>
<th>UL(BF) $\times 10^{-6}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$B^+ \to e^+\nu$</td>
<td>135 ± 4</td>
<td>-0.07 ± 0.03</td>
<td>5.2</td>
</tr>
<tr>
<td>$B^+ \to \mu^+\nu$</td>
<td>120 ± 4</td>
<td>-0.11 ± 0.05</td>
<td>5.6</td>
</tr>
<tr>
<td>$B^0 \to e^+\tau^-$</td>
<td>32 ± 2</td>
<td>0.02 ± 0.01</td>
<td>28</td>
</tr>
<tr>
<td>$B^0 \to \mu^+\tau^-$</td>
<td>27 ± 2</td>
<td>0.01 ± 0.01</td>
<td>22</td>
</tr>
</tbody>
</table>

We reconstruct one of the two $B$ mesons in the event, where it decays semileptonically: $B^+ \to D^{(*)0}\ell^+\nu$. Compared to hadronic tagging method used in in $B^+ \to \ell^+\nu$ and $B^0 \to \ell^+\tau^-$ analyses, this semileptonic tagging method yields higher statistics with more background.

A multivariate classifier, the Random Forest (RF) tool from StatPatternRecognition\textsuperscript{15} is used to optimize signal separation from background. Several regions of the parameter space (terminal leaf size, maximum number of input variables randomly selected for decision splits) are explored with the RF classifier. We use the Punzi Figure of Merit\textsuperscript{16}, $S/(N_\sigma/2+\sqrt{b})$, where $s$ is signal, $b$ is background and $N_\sigma$ is the sigma level of discovery (we take $N_\sigma = 3$), and found the optimal Punzi Figure of Merit with a terminal leaf size of 35 events, after growing 100 decision trees, and sampling on at most 20 variables. The variables include number of tracks in the event (excluding tracks from the $B_{tag}$ reconstruction), transverse momentum of tracks, event topology variables, missing energy in the event, total energy in the event, total energy deposit in the detector that are not associated with any charged or neutral particles.

The signal box is defined in the 2-dimensional space of $D^0$ mass and the RF output, which is blinded until we finish with all selections and estimations. The RF output ranges between 0 and 1. The signal box is RF output bigger than 0.82 and near $D^0$ mass peak which varies depends on the $D$ modes. We estimate the background level in the signal box using MC events as well as data outside of the signal box.

While 30.71 ± 10.71 events are expected 38 events are observed as shown in Figure 3. The systematic uncertainties, which are estimated using double tag events, in where both $B$ mesons decay semileptonically, are incorporated in the UL BF calculation. We set 90% UL BF at 4.2 × 10$^{-5}$, using a modified frequentist method\textsuperscript{17}.

5 Summary

New leptonic $B$ meson decays from BABAR are presented: $B^0 \to \ell^+\ell^-$, $B^+ \to \ell^+\nu$, $B^0 \to \ell^+\tau^-$ and $B \to K\nu\ell$ decays. We have not observed signal and set upper limits on all of these decays. With much more statistics from Super-B factory or ILC, exploiting the hadronic tagging method may be powerful. The leptonic $B$ meson decays will provide us important information on nature with more data.

References

1. Throughout this paper, decay modes imply also their charge conjugations.
Figure 3: The output of the Random Forest for $B \to K\nu\nu$ analysis. The right side of the black line is the signal box. The dashed line is data and the solid line is expected background from MC.

B-hadron lifetimes and rare decays at Tevatron

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Here we present most recent results on the measurements of b-hadron lifetimes, and rare decays using data collected by both CDF and DØ experiments at Fermilab Tevatron. With large dataset collected by both experiments, most stringent limits on some of the rare decays are set.

1 Introduction

Accelerator division at Fermilab Tevatron has so far done an excellent job in delivering large amount of data. Both CDF, and DØ experiments have accumulated more than 3 fb$^{-1}$ of data (at the time of the conference). A large number of measurements of b-hadron lifetimes, and rare decays are performed by these experiments. B-physics program at Tevatron is complementary to the one at B factories, where clear understanding of $B^0$, and $B^+$ has been achieved. Although $par{p}$ collision environment is not as clean, Tevatron enjoys having high $bar{b}$ production cross section, high integrated luminosity, and the possibility of producing heavier b-hadrons, thus a rich b-physics program.

2 Measurement of lifetimes

Lifetime measurement of b hadrons serve as a tool to understand the interaction between heavy and light quarks. Theoretical model known as Heavy Quark Effective Theory (HQET) \cite{HQET}, considers in the leading order, all light quarks as spectator and predicts all b hadrons having same lifetime. Up to about 10% difference between lifetimes of b-hadrons is predicted by HQET originating from the higher order corrections that are proportional to $1/m_b^2$. In order to measure lifetime of b hadron experimentally first we determine the distance traveled by the b hadron in the plane transverse to the beam direction, and correct it for Lorentz boost. We then define
proportional decay length as
\[
\lambda = \frac{L_{xy}}{(\beta\gamma)^T_B} = \frac{L_{xy} cM_B}{p_T}
\]  

here \((\beta\gamma)^T_B\) and \(M_B\) are the transverse boost and the mass of the \(b\) hadron. Finally the lifetime is obtained by performing a simultaneous unbinned maximum likelihood fit to the mass and proper decay length. Both CDF and DØ reported measurement of \(\Lambda_b\) lifetime in exclusive decay channel \(\Lambda_b \rightarrow J/\psi(\rightarrow \mu^+\mu^-)\Lambda(\rightarrow p\pi)\). Figure 1 shows proper decay length distribution for \(\Lambda_b\) decay at CDF using 1.0 \(f_{b^{-1}}\) of data. Measurement of lifetime of \(\Lambda_b\) is also presented as a ratio with the lifetime of \(B^0_d\) decay that has very similar event topology. CDF measurement of lifetime of \(\Lambda_b\), \(\tau(\Lambda_b)\) is 1.580 ± 0.077 (stat) ± 0.012 (syst) ps and \(\tau(\Lambda_b)/\tau(B^0) = 1.018±0.062(stat)±0.007(syst)\). This measurement is about 3\(\sigma\) higher than the theoretical prediction and world average. DØ measurement of \(\tau(\Lambda_b)\) in the same channel using 1.2 \(f_{b^{-1}}\) of data is 1.218±0.130(stat) ± 0.042(syst) ps, and \(\tau(\Lambda_b)/\tau(B^0) = 0.811±0.034\). DØ had also done a measurement of \(\Lambda_b\) lifetime in the semileptonic decay channel \(\Lambda_b \rightarrow \mu\nu\Lambda_c(\rightarrow K_s^0p)X\). This measurement benefits from having large statistics, but as full reconstruction is not possible one cannot observe \(\Lambda_b\) peak. The measured lifetime in this channel is \(\tau(\Lambda_b) = 1.290±0.110(stat)±0.087(syst)\) ps. The most recent status of all \(\tau(\Lambda_b)\) measurement is summarized in Figure 2. CDF has looked into the exclusive decay channels for \(B^+ \rightarrow J/\psi K^+, B^0 \rightarrow J/\psi K^0\), and \(B^0 \rightarrow J/\psi K_s^0\). The measured lifetimes are \(\tau(B^+) = 1.630±0.016(stat)±0.011(syst)\) ps, and \(\tau(B^0) = \)
1.551 ± 0.019 (stat) ± 0.011 (syst) ps. From these measurements $\tau(B^+)/\tau(B^0)$ is found to be 1.015 ± 0.023 (stat) ± 0.004 (syst). These measurements are in good agreement with theoretical prediction 1.

DØ has recently reported on the measurement of the lifetime of $B_c^\pm$ meson. $B_c^\pm$ is one of the most interesting meson studied at Tevatron in that it comprises of two different heavy quarks competing each other for decay. $B_c^\pm$ has the shortest lifetime of weakly decaying b-hadron with explicit predictions of its lifetime to be 0.55 ± 0.15 ps using Operator Product Expansion (OPE), and 0.48 ± 0.05 using QCD sum rules 6. This is about 1/3 of the lifetime of other B mesons. DØ has looked into the decay of $B_c^\pm \to J/\psi \mu \nu$ using 1.3 $fb^{-1}$ of data. Due to the escaping $\nu$ DØ measured the pseudo-proper decay length (PPDL), and corrected it by using a factor that takes boost into account. Finally the lifetime is determined by using simultaneous fit to three-muon invariant mass and PPDL. Presence of $B_c$ signal in the sample is demonstrated in Figure 3(left) that shows the fit to the three muon invariant mass distribution after subtracting $J/\psi$ sideband component and $B^+$ component. A requirement is put on the transverse decay length significance, $L_{xy}/\sigma(L_{xy}) > 4$, where $\sigma(L_{xy})$ is the uncertainty on the measurement of $L_{xy}$. The probability of background fluctuating up to the signal is found to be more than 5$\sigma$. It is important to note that the transverse decay length cut that would bias the lifetime measurement is not applied in the full simultaneous mass and PPDL fit. Figure 3(right) shows the PPDL distribution. DØ measured $^7 \tau(B_c^\pm) = 0.444^{+0.039}_{-0.036} (stat)^{+0.034}_{-0.034} (syst)$ ps. This result is in good agreement with earlier CDF measurement 8, and theoretical prediction 6.

Measurements of $B_S$ lifetimes are recently done by CDF, and DØ experiment. Flavor specific lifetime measurement of $B_S$ is done by CDF experiment using the decays of $B_s \to D^-_s(\phi \pi^-)\pi^+$, and $B_s \to D^-_s \rho^+(\pi^+ \pi^0)$. The second decay channel cannot be fully reconstructed due to the presence of $\pi^0$. Both fully and partially reconstructed channels yields about 1100 events each in 1.3 $fb^{-1}$ of data. The lifetime measurement depends on two fits done sequentially. First relative fraction of events from different signal, and background decay modes is determined by performing a fit on the reconstructed mass of $B_s$ candidates. Then using the fractions obtained from the first fit, a maximum likelihood fit for $B_s$ meson lifetime is performed. The values of the lifetimes obtained in the fully reconstructed and partially reconstructed channels are $1.456 \pm 0.067$ ps, and $1.545 \pm 0.051$ ps respectively. Result of the combination of these two modes is $\tau(B_s) = 1.518 \pm 0.041 \pm 0.025$ ps 9.

DØ experiment has measured the average lifetime of $B_s$, $\bar{B}_s$ states in the decay of $B_s \to J/\psi \phi$, using 2.8 $fb^{-1}$ of data. Value of $\tau(B_s)$ is found to be $1.52 \pm 0.05 \pm 0.01$ ps 10. CDF measurement in the same decay channel using 1.7 $fb^{-1}$ of data is $\tau(B_s) = 1.52 \pm 0.04 \pm 0.02$ ps 11.
3 Rare Decays

Flavor Changing Neutral Current (FCNC) processes are excellent place to study new physics beyond Standard model. The FCNC decays of $B_s(B_d^0) \rightarrow \mu^+ \mu^-$ can only go through higher order Feynman diagrams, and are suppressed by the helicity factor $(m_\mu/m_B)^2$. The decay of $B_d^0$ is further suppressed with respect to the decay of $B_s$ by the ratio of CKM elements, $|V_{td}/V_{ts}|^2$. The predicted branching ratios for $B_s \rightarrow \mu^+ \mu^-$, and $B_d^0 \rightarrow \mu^+ \mu^-$ are $(3.42 \pm 0.54) \times 10^{-9}$, and $(1.00 \pm 0.14) \times 10^{-10}$ respectively. Various extensions of the SM predicts branching ratios that are up to 3 orders magnitude higher. For instance Minimal supersymmetric Standard Model (MSSM) predicts enhancement proportional to $\tan^6 \beta$. CDF experiment has analyzed $2 \text{ fb}^{-1}$ of data to look for $B_s \rightarrow \mu^+ \mu^-$, and $B_d^0 \rightarrow \mu^+ \mu^-$. To achieve best separation between signal and background, at the final stage of the analysis CDF has used a neural network variable, comprising of proper decay length, proper decay length significance, 3D opening angle between dimuon system and the displacement vector between primary vertex and dimuon vertex, and the track isolation of the B candidate. Figure 4 shows the distribution of the neural network output vs invariant mass of the dimuon. Indicated in the boxes are signal windows for $B_s$, and $B_d^0$. No excess of signal over background estimation is observed. CDF has put the worlds best limits on the branching ratio (Br) of $B_s(B_d^0) \rightarrow \mu^+ \mu^-$. These limits at 95% (90%) C.L are $\text{Br}(B_s \rightarrow \mu^+ \mu^-) < 5.8 \times 10^{-8} (4.7 \times 10^{-8})$, and $\text{Br}(B_d^0 \rightarrow \mu^+ \mu^-) < 1.8 \times 10^{-9} (1.5 \times 10^{-8})$. DØ has used $2 \text{ fb}^{-1}$ of data also and did not find any excess in their search for $B_s \rightarrow \mu^+ \mu^-$ signal. Limits obtained by DØ experiment are $\text{Br}(B_s \rightarrow \mu^+ \mu^-) < 9.3 \times 10^{-8} (7.5 \times 10^{-8})$ at 95% (90%) C.L.\textsuperscript{15}

FCNC decays are further suppressed through GIM mechanism\textsuperscript{16} in charmed mesons like D, where the standard model expectation of the branching ratio for $D^+ \rightarrow \pi^+ \mu^+ \mu^-$ is less than $10^{-9}$\textsuperscript{17}. DØ experiment has performed a search for the continuum decay of $D^+ \rightarrow \pi^+ \mu^+ \mu^-$. In order to exclude events coming from the decays $D^+, D_s^+ \rightarrow \phi(\rightarrow \mu^+ \mu^-) \pi^+$ the region where dimuon invariant mass is consistent with $\phi$ mass is excluded. In $1.3 \text{ fb}^{-1}$ of data sample 19 candidate events are observed, whereas $25 \pm 4.6$ events from background sources are expected. This leads to the limit of $\text{Br}(D^+ \rightarrow \pi^+ \mu^+ \mu^-) < 3.9 \times 10^{-6} (6.1 \times 10^{-6})$ at 90% (95%) C.L.\textsuperscript{18}. This is currently the world’s most stringent limit on the decay mediated by $c \rightarrow u \mu^+ \mu^-$ transition.
4 Summary

We have presented a lot of measurements on b-hadron lifetimes, and rare decays with improved uncertainties. For some of the FCNC rare decays most stringent limits in the world have been obtained. In some lifetime measurements uncertainties at the level of 1% have been achieved. Many uncertainties are still dominated by statistics. As we expect to double our dataset by the end of the Tevatron running, we look forward to exciting prospects on both precision lifetime measurements, and rare decays. In fact for some of the FCNC rare decays we expect to get close to the standard model prediction.

5 Acknowledgments

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**B_s MIXING, ∆Γ_s AND CP VIOLATION AT THE TEVATRON**

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We discuss the results from the Tevatron experiments on mixing and CP violation in the \( B^0_s - \bar{B}^0_s \) system, with particular emphasis to the updated measurements of the decay-width difference \( \Delta \Gamma_s \) and the first measurement of the CP-violating phase \( \beta_s \) using flavor tagging information. We also briefly review the charge asymmetry measurements in semileptonic \( B^0_s \) decays and in \( B^\pm \rightarrow J/\psi K^\pm \) decays.

## 1 Introduction

The Tevatron is a \( p\bar{p} \) collider operating at the Fermi National Accelerator Laboratory. The protons and anti-protons collide at a center-of-mass energy of \( \sqrt{s} = 1.96 \) TeV in two interaction points, where the CDF II and DØ detectors are located. The two experiments have collected an integrated luminosity of 3 fb\(^{-1}\) and the measurements presented here span from 1.0 fb\(^{-1}\) to 2.8 fb\(^{-1}\). The physics of the \( b \) quark is a very active research area to challenge the Standard Model predictions. Precise measurements in \( B^0 \) and \( B^+ \) meson decays, performed at the \( B \) factories, improved the understanding of flavor dynamics and proved the Standard Model description very successful. On the other hand, a comparable experimental knowledge of \( B^0_s \) decays has been lacking. The \( B^0_s \) oscillation observation at CDF\(^1\) strongly constrained the magnitude of New Physics contributions in the \( B^0_s \) mixing, while its phase, responsible for CP violating effects, is not precisely determined yet. The \( B^0_s \) sector offers a large variety of interesting processes in which large CP violation effects are still allowed by the current experimental constraints, but are negligible small in the Standard Model. Thus, the Tevatron collider, providing a simultaneous access to large samples of strange and non-strange \( b \)-mesons necessary for precision measurements, offers a great opportunity to study the \( B^0_s \) flavor sector, before the start-up of CERN Large Hadronic Collider (LHC).

## 2 Phenomenology of the \( B^0_s \) System

Flavor oscillation, or mixing, is a very well established phenomenon in particle physics. In the Standard Model the mass and the flavor eigenstates of neutral \( B \) mesons differ. This give rise to particle-antiparticle oscillations, which proceed through forth-order flavor changing weak interactions, whose phenomenology depends on the Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix. The rate at which the neutral \( B - \bar{B} \) transitions occur is governed by the mass difference, \( \Delta m \) of the two mass eigenstates, \( B^L \) and \( B^H \), where the superscripts \( L \) and \( H \) stay for
“light” and “heavy”. The phenomenon of mixing in $B_s^0$ and $\bar{B}_s^0$ mesons is, then, characterized by the mass difference of the two mass eigenstates, $\Delta m_s$, as well as by the decay width-difference $\Delta \Gamma_s = \Gamma_s^L - \Gamma_s^H = 1/\tau_{B_s} - 1/\tau_{\bar{B}_s}$. The latter depends on the CP violating phase defined as $\phi_s = \arg(-M_{12}/\Gamma_{12})$, through the relationship $\Delta \Gamma_s = 2|\Gamma_{12}| \cos(\phi_s)$. $M_{12}$ and $\Gamma_{12}$ are the off-diagonal elements of the $B_s^0 - \bar{B}_s^0$ decay matrix from the Schrödinger equation describing the time evolution of $B_s^0$ mesons$^{2,3}$. While the Standard Model expectations are small$^4$, $\phi_s = 4 \times 10^{-3}$, New Physics could significantly modify the observed phase value contributing with additional processes, $\phi_s = \phi_s^{SM} + \phi_s^{NP}$. The same phase would alter the observed phase between the mixing and the $b \to c\bar{c}s$ transitions, $2\beta_s = 2\beta_s^{SM} - \phi_s^{NP}$, in which the Standard Model contribution is defined as $-2\beta_s^{SM} = -2 \arg(-V_{ts}/V_{td}) \approx \mathcal{O}(0.04)$, where $V_{ij}$ are the elements of the CKM matrix. Since both $\phi_s^{SM}$ and $\beta_s^{SM}$ are tiny with respect to the current experimental resolution, we can approximate $\phi_s = -2\beta_s$. A measurement of sizable value of $2\beta_s$ ($\phi_s$) would be a clear indication of New Physics$^{2,3}$.

3 $B_s^0$ Mixing

While $\Delta m_d$ was precisely determined at the $B$ factories$^{5,6}$, the $B_s^0$ mixing frequency has been first measured at CDF experiment$^1$. The $B_s^0 - \bar{B}_s^0$ oscillation observation was achieved through a combination of several data-sets of 1 fb$^{-1}$, in integrated luminosity, which results in:

$$\Delta m_s = 17.77 \pm 0.10 \text{ (stat.)} \pm 0.07 \text{ (syst.) ps}^{-1}, \tag{1}$$

with a significance greater than 5 standard deviations. Two independent types of flavor tags are used to identify the $B^0_s$ flavor at production: the Opposite Side Tagger (OST) and the Same Side Kaon Tagger (SSKT). The performance of flavor taggers are quantified by the efficiency $\epsilon$ and the dilution $D$, defined as the probability to correctly tag a candidate. The tagging effectiveness, $\epsilon D^2$ of the OST is 1.8%. The SSKT has $\epsilon D^2 = 3.5\%$ (hadronic) and 4.8% (semileptonic) and thus contributes most to the sensitivity of the CDF analysis. The accurate measurement of the $B_s^0 - \bar{B}_s^0$ mixing frequency offers a powerful constraint to the ratio $|V_{ts}|^2/|V_{td}|^2$ of CKM matrix elements:

$$\frac{|V_{ts}|^2}{|V_{td}|^2} = 0.2060 \pm 0.0007 \text{ (stat.)}^{+0.0081}_{-0.0060} \text{ (theory).} \tag{2}$$

DØ recently reported a measurement of the $B_s^0$ oscillation frequency$^7$ using a large sample of semileptonic $B_0^s$ decays and their first hadronic mode, $B_0^s \to D_s^- [\to \phi (\to K^+K^-) \pi^-] \pi^+$. DØ combines the tagging algorithms using a likelihood-ratio method, obtaining a total effective tagging power $\epsilon D^2 = (4.49 \pm 0.88)\%$. With a data-set of approximately 2.4 fb$^{-1}$, they obtains:

$$\Delta m_s = 18.56 \pm 0.87 \text{ (stat.) ps}^{-1}. \tag{3}$$

The result statistically exceeds the $3\sigma$ significance and it is compatible with the CDF measurement. The $\Delta m_s$ is well consistent with the Standard Model unitarity hypothesis for the CKM matrix.

4 Phase of the Mixing Amplitude and Decay-Width Difference in the $B_s^0$ System

We present the time-dependent angular analyses of $B_s^0 \to J/\psi (\to \mu^+\mu^-) \phi (\to K^+K^-)$ decay mode performed at the Tevatron experiments. The decay $B_s^0 \to J/\psi\phi$ proceeds through the $b \to c\bar{c}s$ transition and gives rise to both CP-even and CP-odd final states. Through the angular distributions of the $J/\psi$ and $\phi$ mesons, it is possible to statistically separate the two final states.
with different CP eigenvalues, thus allowing to determine the phase $\beta_s$ and to separate lifetimes for the mass eigenstates, so to measure the decay-width difference $\Delta \Gamma_s$. After the DØ analysis\(^8\) of untagged $B^0_s \to J/\psi \phi$ decay sample of 1.1 fb\(^{-1}\), and reported at Moriond 2007, the CDF Collaboration presents a similar analysis with a sample of 1.7 fb\(^{-1}\) in integrated luminosity\(^9\). CDF measures $\Delta \Gamma_s = 0.076^{+0.059}_{-0.063}$ (stat.) \pm 0.006 (syst.) ps\(^{-1}\), $c\tau_s = 456 \pm 13$ (stat.) \pm 7 (syst.) \mu m, assuming CP conservation ($\beta_s = 0$) results. To date, this is one of the most precise $B^0_s$ lifetime measurements and it is in excellent agreement with both the DØ results and the theoretical expectations predicting $\tau_s = \tau_d \pm \mathcal{O}(1\%)$. Allowing CP violation, a bias and non-Gaussian fit estimates are observed in pseudo-experiments for statistics similar to the present data-sets. The observed bias originates from the loss of degree of freedom of the likelihood for certain values of the parameters of interest and does not permit a point estimation of $\Delta \Gamma_s$ and $\beta_s$. Thus, CDF provides confidence level regions in the $2\beta_s - \Delta \Gamma_s$ plane using the likelihood ratio ordering of Feldman and Cousins\(^10\). For the Standard Model expectation ($\Delta \Gamma_s \approx 0.096$ ps\(^{-1}\) and $2\beta_s = 0.04$ rad\(^4\)), the probability to get equal or greater likelihood ratio than the one observed in data is 22\%, which corresponds to 1.2 Gaussian standard deviations. Figure 1 shows the CDF and the DØ results in the $2\beta_s - \Delta \Gamma_s$ plane. Furthermore, the CDF Collaboration performed an angular analysis on the $B^0 \to J/\psi(\rightarrow \mu^+ \mu^-)K^0(\rightarrow 2\pi^-)$ decay mode for the measurement of the transversity amplitudes and strong phases. Such an analysis plays a key role in the validation of the entire framework used for the $B^0_s \to J/\psi \phi$ angular analysis. The results obtained for the transverse linear polarization amplitudes at $t = 0$, $A_\parallel$ and $A_\perp$, corresponding to CP even and CP odd final states respectively, as well as the strong phases $\delta_\parallel = \arg(A_\parallel^* A_0)$ and $\delta_\perp = \arg(A_\perp^* A_0)$, are $|A_\parallel|^2 = 0.569 \pm 0.009$ (stat.) $\pm 0.009$ (syst.), $|A_\perp|^2 = 0.211 \pm 0.012$ (stat.) $\pm 0.006$ (syst.), $\delta_\parallel = -2.96 \pm 0.08$ (stat.) $\pm 0.03$ (syst.) and $\delta_\perp = 2.97 \pm 0.06$ (stat.) $\pm 0.01$ (syst.), which are in agreement and competitive with the current B factories results\(^11\).

We present the first Tevatron studies of the $B^0_s \to J/\psi \phi$ decay mode when the initial state of the $B^0_s$ meson is identified exploiting the flavor tagging information. In fact, such information allows to separate the time evolution of mesons originally produced as $B^0_s$ or $\bar{B}^0_s$. The angular analyses which do not use the flavor tagging are sensitive to $|\cos(2\beta_s)|$ and $|\sin(2\beta_s)|$, leading to a 4-fold ambiguity in the likelihood for the determination of $2\beta_s$ (see Figure 1). On the other hand, utilizing flavor tagging algorithms, the analyses gain sensitivity to the sign of $\sin(2\beta_s)$ reducing by half the allowed region for $\beta_s$. CDF performed a flavor tagged analysis on a 1.35 fb\(^{-1}\) data-set of $B^0_s \to J/\psi \phi$ reconstructed events, which yields $\approx 2,000$ signal candidates\(^12\).
The measured efficiencies for OST and SSKT are \( \epsilon_{OST} = (96 \pm 1)\% \) and \( \epsilon_{OST} = (50 \pm 1)\% \). The dilutions are respectively \( D_{OST} = (11 \pm 2)\% \) for the OST and \( D_{SSKT} = (27 \pm 4)\% \) for the SSKT. The addition of tagging information improves the regularity of the likelihood with respect to the untagged case, but still non-Gaussian uncertainties and biases are observed in simulated experiments with the available statistics. Thus, CDF reports a confidence region constructed according to the Feldman Cousins criterion with rigorous inclusion of systematics uncertainties. In fact, any \( \Delta \Gamma_s - \beta_s \) pair is excluded at a given CL only if it can be excluded for any choice of all other fit parameters, sampled uniformly within \( \pm 5 \sigma \) of the values determined in their estimate on data. Assuming the Standard Model predicted values of \( 2\beta_s = 0.04 \) rad and \( \Delta \Gamma_s = 0.096 \) ps\(^{-1} \), the probability of a deviation as large as the observed data is 15\%, which corresponds to 1.5 Gaussian standard deviations. Moreover, if \( \Delta \Gamma_s \) is treated as a nuisance parameter, thus fitting only for \( 2\beta_s \), CDF finds \( 2\beta_s \in [0.31, 2.82] \) rad at the 68\% confidence level. By exploiting the current experimental and theoretical information, CDF extracts tighter bounds on the CP violation phase \( \beta_s \). Imposing the constraint on \( |\Gamma_{12}| = 0.048 \pm 0.018 \) ps\(^{-1} \) in \( \Delta \Gamma_s = 2|\Gamma_{12}| \cos(2\beta_s) \) \(^3\), \( 2\beta_s \in [0.24, 1.36] \cup [1.78, 2.90] \) rad at the 68\% CL. Additionally constraining the strong phases \( \delta_0 \) and \( \delta_\perp \) to the B factories results on \( B^0 \to J/\psi K^{*0} \) \(^{11}\) and the \( B^0 \) mean width to the world average \( B^0 \) width \(^{13}\), it is found \( 2\beta_s \in [0.40, 1.20] \) rad at the 68\% CL. The DØ Collaboration reports an analysis \(^{14}\) on 2,000 signal \( B^0 \to J/\psi \) candidates, reconstructed in 2.8 fb\(^{-1} \). DØ combines the tagging algorithms, as done in their \( B^0 \) mixing analysis. The total tagging power is \( \epsilon D^2 = (4.68 \pm 0.54)\% \) and a tag is defined for 99.7\% of the events. To overcome the likelihood pathologies described above, DØ decides to vary the strong phases around the world-averaged values for the \( B^0 \to J/\psi K^{*0} \) decay \(^{15}\), applying a Gaussian constraint. This removes the 2-fold ambiguity, inherent the measurement for arbitrary strong phases. The strong phases in \( B^0 \to J/\psi K^{*0} \) and \( B^0 \to J/\psi \) cannot be exactly related in the \( SU(3) \) symmetry limit, so the width of the Gaussian is chosen to be \( \pi/5 \), allowing for some degree of \( SU(3) \) symmetry violation. The fit with all floating parameters yields to the measurements

\[
\begin{align*}
\phi_s &= -0.57^{+0.24}_{-0.30} \text{ (stat.)}^{+0.07}_{-0.02} \text{ (syst.)} \text{ rad}, \\
\Delta \Gamma_s &= 0.19 \pm 0.07 \text{ (stat.)}^{+0.02}_{-0.01} \text{ (syst.)} \text{ ps}^{-1}, \\
\tau_s &= 1.52 \pm 0.05 \text{ (stat.)} \pm 0.01 \text{ (syst.)} \text{ ps}.
\end{align*}
\]

The allowed ranges at the 90\% CL for the parameters of interest are found to be \( \phi_s \in [-1.20, 0.06] \) rad and \( \Delta \Gamma_s \in [0.06, 0.30] \) ps\(^{-1} \). The expected confidence level contours in the \( \phi_s - \beta_s \) plane at 68\% and 90\% CL are depicted in Figure 2. The level of agreement with the Standard Model corresponds to 6.6\%, which is obtained by generating pseudo-experiments with the initial value for \( \phi_s \) set to \(-0.04 \) rad and counting the events whose obtained fitted value of the phase is lower than the measured \(-0.57 \) rad. The results supersede the previous DØ untagged analysis on a smaller \( B^0 \) \( \to J/\psi \) sample.

5 Charge Asymmetry in \( B^0 \) Semileptonic Decays

Another way of studying the CP violation induced by the \( B_s \) mixing, is to measure the charge asymmetry in semileptonically decaying mesons. The charge asymmetry is connected to the CP violating phase \( \phi_s \), through the relationship \( A_{SL}^s = \Delta \Gamma_s / \Delta m_s \times \tan(\phi_s) \). With the underlying assumption of \( \phi_s = -2\beta_s \) (see Section 2), an independent measurements on charge asymmetry could be used to constrain the CP violating phase \( \beta_s \). \(^{16}\) DØ Collaboration performed two independent analyses to extract \( A_{SL}^s \). The first result is based on the di-muon charge asymmetry measurement \(^{17}\), defined as

\[
A_{SL}^{\mu\mu} = \frac{N(bb \to \mu^+\mu^+) - N(bb \to \mu^-\mu^-)}{N(bb \to \mu^+\mu^+) + N(bb \to \mu^-\mu^-)}.
\]
The following asymmetry gets its contributions from both $B^0$ and $B^0_s$: by using the world average value for $B^0$ and $B^0_s$ production fractions and the $B^0$ charge asymmetry measurements from the $B$ factories, DØ extracts the $B^0_s$ charge asymmetry on a data-set of 1.0 fb$^{-1}$:

$$A_{\mu \mu, B^0_s}^{\mu \mu, B^0} = -0.0064 \pm 0.0101 \text{ (stat. + syst.).} \quad (6)$$

CDF Collaboration also released a similar measurement of the di-muon charge asymmetry on a sample of 1.6 fb$^{-1}$ data. In this analysis, the unbinned likelihood is performed using the impact parameter information of the two muons, in order to separate the $b - \bar{b}$ component of the sample from the others which arise from prompt and charm sources:

$$A_{\mu \mu, B^0_s}^{\mu \mu, B^0} = 0.020 \pm 0.021 \text{ (stat.)} \pm 0.016 \text{ (syst.)} \pm 0.009 \text{ (inputs)}. \quad (7)$$

Additionally to the statistical and systematic uncertainties, the last uncertainty term comes from the world average value for $B^0$ and $B^0_s$ production fractions and the $B^0$ charge asymmetry measurements already discussed in the description of DØ results. Compared to CDF, DØ analysis has strongly reduced systematics uncertainties thanks to a regular flipping of the magnet polarity. Such technique, removing most of the artificial asymmetry in the detector response, is constantly used by DØ to measure all the charge asymmetries described along this paper.

The DØ Collaboration probes the $\phi_s$ phase also by measuring the charge asymmetry in an untagged sample of $B^0_s \to \mu D_s$ decays, with $D_s \to \phi (\to K^+ K^-) \pi$. With a data-set of 1.3 fb$^{-1}$ the charge asymmetry is found to be:

$$A_{\mu \mu, D_s}^{\mu \mu, B^0_s} = 0.0245 \pm 0.0193 \text{ (stat.)} \pm 0.0035 \text{ (syst.).} \quad (8)$$

### 6 Charge Asymmetry in $B^+ \to J/\psi K^+$ Decay

We present a search for direct $CP$ violation in $B^+ \to J/\psi K^+$ decays. The event sample is selected from 2.8 fb$^{-1}$ of $p\bar{p}$ collisions recorded by DØ experiment. The charge asymmetry is defined as

$$A_{CP}(B^+ \to J/\psi K^+) = \frac{N(B^- \to J/\psi K^-) - N(B^+ \to J/\psi K^+)}{N(B^- \to J/\psi K^-) + N(B^+ \to J/\psi K^+)}.$$  

By using a sample of approximately 40,000 $B^+ \to J/\psi K^+$ decays, the asymmetry is measured to be $A_{CP} = 0.0075 \pm 0.0061 \text{ (stat.)} \pm 0.0027 \text{ (syst.).}$ The result is consistent with the
world average $^{13}$ and the Standard Model expectation $A_{CP}(B^+ \rightarrow J/\psi K^+) \simeq 0.003$ $^{21}$, but has a factor of two improvement in precision, thus representing the most stringent bound for new models which predict large values of $A_{CP}(B^+ \rightarrow J/\psi K^+)$. Furthermore, DØ provides the direct CP violating asymmetry measurement in $B^+ \rightarrow J/\psi \pi^+$, $A_{CP}(B^+ \rightarrow J/\psi \pi^+) = -0.09 \pm 0.08$ (stat.) $\pm 0.03$ (syst.). The result agrees with the previous measurements of this asymmetry $^{13}$ and has a competitive precision.

7 Conclusions

After the successful $B_s^0$ oscillation observation, the CDF and DØ Collaboration directed their effort in the exploration of the mixing-induced CP violation effect in the $B_s^0$ system. We described the first tagged measurement in $B_s^0 \rightarrow J/\psi \phi$ performed at the CDF II detector, which improved the sensitivity to the CP violating phase $\beta_s$, excluding negative and large values for the phase itself. The DØ Collaboration promptly delivered a similar analysis confirming the results. The agreement of the analyses of $B_s^0 \rightarrow J/\psi \phi$ decays, shows an interesting fluctuations in the same direction from CDF and DØ experiments and they will certainly need further investigations to support an evidence, which would be possible exploiting the full Run II data sample, if these first indications are confirmed in the future. We also reviewed the charge asymmetry measurements of $B_s^0$ semileptonic decays, which provide another independent test for the CP violation in $B_s^0$ mixing and can be combined with the analyses on $B_s^0 \rightarrow J/\psi \phi$ to get a better understanding of the CP violating phenomena. Finally, we presented the world most precise direct CP violating asymmetry in the $B^+ \rightarrow J/\psi K^+$ decay mode. The Tevatron experiments are becoming increasingly competitive with B factories results on $B^0/B^+$ decays and complementary to them in corresponding $B_s^0$ modes. Since many of the analyses reported do not even use half of the statistics available, significant improvements are expected in the future, as the Tevatron keeps producing data.

References

BELLE NEW RESULTS ON $B \to D^{**}\ell\nu$ DECAYS

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We present a study of semileptonic $B$ decays to $P$-wave $D^{**}$ mesons at Belle. Semileptonic decay to a $D_{2}^{*}$ meson is observed for the first time and its product branching ratio is measured to be $B(B^{+} \to D_{2}^{0}\ell^{+}\nu) \times B(D_{2}^{0} \to D^{-}\pi^{+}) = 0.22 \pm 0.03(\text{stat.}) \pm 0.04(\text{syst.})\%$.

1 Introduction

Heavy Quark Effective Theory (HQET) has proven to be very successful at describing semileptonic decays of $B$-mesons, especially inclusive transitions. However, some difficulties arise when it is applied to exclusive decays. For example, certain sum rules (in particular, the Uraltsev sum rule\textsuperscript{1}) imply the strong dominance of decays to the narrow excited $D$ mesons over those to the wide ones, while some experimental data show the opposite trend\textsuperscript{2,3}. However, no complete experimental study of such semileptonic decays to excited $D$ mesons exists, and thus no direct comparison with theoretical predictions can be performed. Here we present Belle study of $B \to D^{(*)}\pi\ell\nu$ decays and measurement of the excited $D$ contributions to the $D^{(*)}\pi$ final state\textsuperscript{4}.

According to HQET there are two doublets of orbitally excited ($P$-wave) charmed mesons ($D^{**}$), differentiated by their light quark angular momentum $j_{q} = 1/2$ or $j_{q} = 3/2$. Members of the $j_{q} = 3/2$ doublet are predicted to decay only via a $D$-wave and be relatively narrow, while members of the $j_{q} = 1/2$ doublet are predicted to decay only via an $S$-wave and be relatively broad\textsuperscript{5}. The $D^{**}$ states with spin-parity and light quark angular momentum combinations $0^{+}(j_{q} = 1/2), 1^{+}(j_{q} = 1/2), 1^{+}(j_{q} = 3/2)$ and $2^{+}(j_{q} = 3/2)$ are usually labelled $D_{0}^{*}, D_{1}^{*}, D_{1}$ and $D_{2}^{*}$, respectively. The $D^{**}$ states have previously been observed and studied in hadronic $B$ decays\textsuperscript{6}. Semileptonic $B$ decays to narrow $D_{1}$ and $D_{2}^{*}$ mesons have been studied by a number of experiments\textsuperscript{7}. The semileptonic branching fractions of $B \to D^{(*)}\pi\ell\nu$ decays were measured by Belle\textsuperscript{8} and BaBar\textsuperscript{9}.
This measurement is based on a data sample that contains 657 million $B\bar{B}$ pairs, which corresponds to 605 fb$^{-1}$, collected at the $\Upsilon(4S)$ resonance with the Belle detector\cite{2} operating at the KEKB asymmetric-energy $e^+e^-$ collider\cite{3}. An additional 68 fb$^{-1}$ data sample taken at a center-of-mass energy 60 MeV below the $\Upsilon(4S)$ resonance is used to study continuum $e^+e^- \rightarrow q\bar{q}$ ($q = u, d, s, c$) background.

2 Data analysis

To suppress the large combinatorial background expected in the reconstruction of final states including a neutrino, we use a full reconstruction tagging method. The first $B$ meson (denoted as $B_{\text{sl}}$) is reconstructed in the semileptonic mode of interest, i.e. as a combination of all final particles $D^{(*)}\pi\ell$ except for the neutrino. The remainder of the event is combined into either a $D^{(*)}\pi\ell\nu$ (n = 6) or $D^{(*)}\rho\ell\nu$ combination to form the tagging $B$ meson (referred to below as $B_{\text{tag}}$). Semileptonic decays are identified by a peak around zero in the missing mass squared spectrum, $M^2_{\text{miss}} = (P_{\text{beams}} - P_{\text{tag}} - P_{\text{sl}})^2$, where $P_{\text{beams}}$ is the total four-momentum of the beams and $P_{\text{tag}}$ and $P_{\text{sl}}$ are the reconstructed four-momenta of the $B_{\text{sl}}$ and $B_{\text{tag}}$, respectively. This method provides significantly improved resolution in the missing momentum in comparison with non-tagging methods, thus allowing background suppression, separation of different decay modes and precise calculation of the decay kinematics. The $M^2_{\nu}$ spectra for the four semileptonic decays $B \rightarrow D^{(*)}\pi\ell\nu$ are shown in Figs. 1, 1a)–1d) as points with error bars.

We divide the backgrounds into the following categories: (1) Continuum, (2) Backgrounds with the $B_{\text{sl}}$ misreconstructed from particles belonging to the other $B$ meson or fake tracks, (3) $B_{\text{sl}}$ background with the $B_{\text{tag}}$ reconstructed correctly, which can be further separated by their source: (3a) Combinatorial background under the $D^{(*)}$ signal from $B_{\text{sl}}$, (3b) Hadrons misidentified as leptons, (3c) Feed-down from $B \rightarrow D^*\pi\ell\nu$ reconstructed as $B \rightarrow D\pi\ell\nu$ with lost neutral(s). All backgrounds except for (3c) are reliably determined and finally subtracted directly from the data. Background (3c) is observed only in the $B \rightarrow D\pi\ell\nu$ channels and is estimated from a Monte Carlo (MC) simulation with normalization fixed to the data using $B \rightarrow D^*(\pi)\ell\nu$ signal yields. This contribution is plotted in Figs. 1, 1a), 1c) as open histograms.

The background-subtracted $M^2_{\nu}$ distributions are shown in Figs. 1, 2a)–2d). These distributions are fitted with signal functions, the shapes of which are fixed from MC studies. Fitted signal yields, reconstruction efficiencies and branching ratios are summarized in Table 1. The branching ratios are calculated relative to the normalization modes $B \rightarrow D\ell\nu$ to cancel out the $B_{\text{tag}}$ reconstruction efficiency according to the formula: $\mathcal{B}(\text{mode}) = \mathcal{B}(\text{norm}) \times N_{\text{norm}}/N_{\text{mode}} \times \epsilon_{\text{norm}}/\epsilon_{\text{mode}}$, where $N_{\text{norm(mode)}}$ and $\epsilon_{\text{norm(mode)}}$ are the signal yield and reconstruction efficiency of the normalization mode (mode of interest) and the normalization mode $B$ is taken from the PDG\cite{4}. Relative efficiencies are obtained from MC simulation. Intermediate branching fractions are included, while the tagging efficiency is not. The reconstruction and background subtraction procedures for the $B \rightarrow D\ell\nu$ mode are identical to those applied for the studied channels. The obtained branching fractions are in good agreement with our previous measurement\cite{5} and with BaBar results\cite{6}.

Signals for semileptonic $B$ decays to orbitally excited $D^{*\ast}$ are extracted from the $D^{(*)}\pi$ invariant mass distributions. We define a signal window for $B \rightarrow D^{(*)}\pi\ell\nu$ decays by the requirement $|M^2_{\nu}| < 0.1 \text{GeV}^2/c^4$. The backgrounds are estimated in the same way as in the $M^2_{\nu}$ distribution study. The $D^{(*)}\pi$ invariant mass spectra from the signal window after subtraction of backgrounds (1–3) are shown in Fig. 2. The mass distributions before background subtraction, restricted to the region near the $j_\pi = 3/2$ states, are shown in the insets.

To extract the $D^{*\ast}$ signals we perform simultaneous unbinned likelihood fits to the signal and background $D^{(*)}\pi$ mass spectra. The signal function includes all orbitally excited $D^{*\ast}$ contributing to the given final state ($D_0$ and $D_2$ to $D\pi$ and $D_1$, $D_1'$, $D_2^*$ to $D^*\pi$), each of
Rencontres de Moriond 2008

Figure 1: $M_\nu^2$ spectra before (1) and after (2) background subtraction for: a) $B^+ \to D^-\pi^+\ell^+\nu$, b) $B^+ \to D^{*-}\pi^+\ell^+\nu$, c) $B^0 \to D^0\pi^-\ell^+\nu$, d) $B^0 \to \bar{D}^{*0}\pi^-\ell^+\nu$. The curves are the fits, which are described in the text.

which is described by a relativistic Breit-Wigner function for a known orbital momenta, and a non-resonant part described by the Goity-Roberts model\textsuperscript{14}. $D^{**}$ masses and widths are fixed to measured values\textsuperscript{6}. To further investigate the $D\pi$ mass spectrum we also test a $D^+_\nu + D^*_2$ hypothesis. Despite the $D^{0}\pi^+$ mass region corresponding to $D^{**}$ being excluded from the study, and while $D^{*0}$ is below the $D^{+}\pi^+$ threshold, a virtual $D^*_\nu$ can be produced off-shell. We describe the $D^*_\nu$ contribution by a tail of the Breit-Wigner function with floating normalization. Fit results are shown as a dashed line for this combination.

Fitted resonance yields and corresponding product branching ratios are listed in Table 2. The contribution of the non-resonant component in all cases is consistent with zero. The $B \to D^{**}\ell\nu$ decay significance is defined as $\sqrt{-2\ln L_{\text{max}}/L_0}$, where $L_0$ is the likelihood value returned by the fit to the $D^{(*)}\pi$ distribution with the $D^{**}$ contribution fixed to zero. Our result for $B \to \bar{D}_1\ell^+\nu$ is in good agreement with previous measurements\textsuperscript{7}. For a $D^*_0 + D^*_2$ hypothesis the branching

Table 1: Results for $B \to D^{(*)}\pi\ell\nu$ where the first error is statistical and the second is systematic.

<table>
<thead>
<tr>
<th>Mode</th>
<th>Yield</th>
<th>Eff.,%</th>
<th>$\mathcal{B}$ (mode),%</th>
</tr>
</thead>
<tbody>
<tr>
<td>$B^+ \to D^0\ell^+\nu$</td>
<td>$2320 \pm 60$</td>
<td>6.4</td>
<td>$2.15 \pm 0.22$\textsuperscript{13}</td>
</tr>
<tr>
<td>$B^+ \to D^-\pi^+\ell^+\nu$</td>
<td>$192 \pm 19$</td>
<td>2.8</td>
<td>$0.40 \pm 0.04 \pm 0.06$</td>
</tr>
<tr>
<td>$B^+ \to D^{*-}\pi^+\ell^+\nu$</td>
<td>$123 \pm 14$</td>
<td>1.14</td>
<td>$0.64 \pm 0.08 \pm 0.09$</td>
</tr>
<tr>
<td>$B^0 \to D^-\ell^+\nu$</td>
<td>$760 \pm 30$</td>
<td>3.7</td>
<td>$2.12 \pm 0.20$\textsuperscript{13}</td>
</tr>
<tr>
<td>$B^0 \to \bar{D}^{0}\pi^-\ell^+\nu$</td>
<td>$150 \pm 20$</td>
<td>3.7</td>
<td>$0.42 \pm 0.07 \pm 0.06$</td>
</tr>
<tr>
<td>$B^0 \to \bar{D}^{*0}\pi^-\ell^+\nu$</td>
<td>$22 \pm 8$</td>
<td>0.40</td>
<td>$0.56 \pm 0.21 \pm 0.08$</td>
</tr>
</tbody>
</table>
ratio of the decay to the wide $D_0^*$ is large, in contrast to theoretical predictions. However, the present statistics do not definitely exclude an interpretation of broadly distributed $D\pi^+$ events as the $D_v^*$ tail.

For $D^{*,**}$'s decaying into $D\pi$ we perform a study of the helicity angle distributions, which is the angle between $\pi$ momentum and the direction opposite to $B_{\text{sl}}$-momentum in the $D^{*,**}$ rest frame. To extract the $D_v^*$, $D_0^*$ and the $D_2^*$ helicity distributions we perform a combined fit of the $M(D\pi)$ spectra for $D\pi$ combinations from both $B^+$ and $B^0$ in bins of helicity angle. The fit procedure is identical to that used for the $B(B \to D^{*,**}\ell\nu)$ calculation. The results corrected for the efficiency are plotted in Fig. 3. $D_2^*$ distributions for $D_v^*$ and $D_0^*$ hypothesis coincide within errors, so that only for the $D_0^* + D_2^*$ case is shown in Fig. 3 c. The $D_0^*$ helicity distribution is consistent with the $J = 0$ hypothesis ($\chi^2/ndf = 6.0/4$, where ndf is the number of degrees of freedom). The $D_2^*$ helicity distribution is fitted with the function $a_0^2|Y_0^1|^2 + a_1^2|Y_1^1|^2 + a_2^2|Y_2^1|^2$, where the $Y_j^i$ are spherical harmonics and $a_0^2 + 4a_1^2 + 4a_2^2 = 1$. The fit yields $a_0^2 = 0.74 \pm 0.10$, $a_1^2 = 0.04 \pm 0.02$ and $a_2^2 = 0.02 \pm 0.02$; the fit quality is $\chi^2/ndf = 2.0/3$. The fit is consistent with the assumed quantum numbers and demonstrates that the $D_2^*$ from semileptonic decay is dominantly in the $s_z = 0$ spin projection. Helicity distributions, predicted by theory, are shown as dashed lines. For evaluating the $D_v^* + D_2^*$ hypothesis, the obtained $D_v^*$ helicity distribution (Fig. 3 b) is fitted with the function $b_0^2|Y_0^1|^2 + b_1^2|Y_1^1|^2$. This fit yields
$b_0^2 = 0.15 \pm 0.09$, $b_1^2 = 0.85 \pm 0.09 \ (\chi^2/\text{ndf} = 18.8/4)$ in poor agreement with expectations from theory, shown as a dashed line.

We also study the dependence of the $B \to D^{**}$ transition on $q^2$ or, equivalently, on the conventional HQET variable $w$, which is the dot-product of $B$ and $D^{**}$ four-velocities: $w = v_B \cdot v_{D^{**}}$. The $w$-dependence is obtained from fits of $D \pi$ invariant mass in bins of $w$. The results are presented in Fig. 4. As with the helicity study the $D^*_3$ distribution is shown only for the $D_0^*$ hypothesis in Fig. 4 c. The $w$ distribution is fitted according to the model given in Ref. 15. In HQET, the matrix elements between the $B$ and $D$ states to leading order in $\Lambda_{\text{QCD}}/m_Q$ are expressed in terms of three universal Isgur-Wise functions $\xi(w)$, $\tau_{1/2}(w)$ and $\tau_{3/2}(w)$ for $(D, D^*)$, $(D_0^*, D_0^*')$ and $(D_1, D_2)$ doublets, respectively. We assume a “pole” form for $\xi(w)$: $\xi = (2/(1 + w))^{2/3} w^2$ and a linear form for $\tau_i(w)$ functions: $\tau_i(w) = \tau_i(1)[1 + \hat{\tau}_i^2(w - 1)]$, and the following relation: $\hat{\tau}_{1/2}^2 = \hat{\tau}_{3/2}^2 + 0.5$. A simultaneous fit to the $w$-distributions for $D_0^*$ and $D^*_3$ gives $\hat{\tau}_{3/2}^2 = -1.8 \pm 0.3$. Using the measured branching ratios of $B \to D_0^* \ell \nu$, we also calculate $\tau_{3/2}(1) = 0.75$ and $\tau_{1/2}(1) = 1.28$. All parameters are in agreement with expectations except for $\tau_{1/2}(1)$, which is larger than predicted due to the large value of $B(B \to D_0^* \ell \nu)$.

Table 2: Results of the $D^{(*)} \pi^+$ pair invariant mass study. $B(\text{mode}) \equiv B(B \to D^{**} \ell \nu) \times B(D^{**} \to D^{(*)} \pi^+)$. The first error is statistical and the second is systematic.

<table>
<thead>
<tr>
<th>Mode</th>
<th>Yield</th>
<th>$B(\text{mode}),%$</th>
<th>Signif.</th>
</tr>
</thead>
<tbody>
<tr>
<td>$B^+ \to D_0^{0*} \ell^+ \nu$</td>
<td>$102 \pm 19$</td>
<td>$0.24 \pm 0.04 \pm 0.06$</td>
<td>5.4</td>
</tr>
<tr>
<td>$B^+ \to D_2^{0*} \ell^+ \nu$</td>
<td>$94 \pm 13$</td>
<td>$0.22 \pm 0.03 \pm 0.04$</td>
<td>8.0</td>
</tr>
<tr>
<td>$B^0 \to D_0^{-} \ell^+ \nu$</td>
<td>$61 \pm 22$</td>
<td>$0.20 \pm 0.07 \pm 0.05$</td>
<td>2.6</td>
</tr>
<tr>
<td>$B^0 \to D_2^{-} \ell^+ \nu$</td>
<td>$68 \pm 13$</td>
<td>$0.22 \pm 0.04 \pm 0.04$</td>
<td>5.5</td>
</tr>
<tr>
<td>$B^+ \to D_1^{0*} \ell^+ \nu$</td>
<td>$-5 \pm 11$</td>
<td>$&lt; 0.07 @ 90% \text{ C.L.}$</td>
<td></td>
</tr>
<tr>
<td>$B^+ \to D_2^{0*} \ell^+ \nu$</td>
<td>$81 \pm 13$</td>
<td>$0.42 \pm 0.07 \pm 0.07$</td>
<td>6.7</td>
</tr>
<tr>
<td>$B^+ \to D_3^{0*} \ell^+ \nu$</td>
<td>$35 \pm 11$</td>
<td>$0.18 \pm 0.06 \pm 0.03$</td>
<td>3.2</td>
</tr>
<tr>
<td>$B^0 \to D_1^{-} \ell^+ \nu$</td>
<td>$4 \pm 8$</td>
<td>$&lt; 0.5 @ 90% \text{ C.L.}$</td>
<td></td>
</tr>
<tr>
<td>$B^0 \to D_2^{-} \ell^+ \nu$</td>
<td>$20 \pm 7$</td>
<td>$0.54 \pm 0.19 \pm 0.09$</td>
<td>2.9</td>
</tr>
<tr>
<td>$B^0 \to D_3^{-} \ell^+ \nu$</td>
<td>$1 \pm 6$</td>
<td>$&lt; 0.3 @ 90% \text{ C.L.}$</td>
<td></td>
</tr>
</tbody>
</table>

Figure 3: Helicity distributions for a) $D_0^*$, b) $D_1^*$, c) $D_2^*$. The curves represent the fits, described in the text.
3 Conclusion

In conclusion, we measured the branching fractions for $B \rightarrow D^{(*)}\pi\ell\nu$ decays. We also performed an analysis of the final state $D^{(*)}\pi$ hadronic system and obtained branching ratios for the $B \rightarrow D^{**}\ell\nu$ components. Semileptonic decay to $D_S^2$ meson is observed and measured for the first time. Helicity and $w$ distributions are studied for this decay. We observe a broad enhancement in the $D\pi$ mass distribution consistent with wide $D_S^0$ production. The branching ratio of the decay to $B \rightarrow D_S^0\ell\nu$ is found to be large, in contrast with theoretical predictions. However there is no indication of a broad $D_S^0$ in the $B \rightarrow D^*\pi\ell\nu$ channel, which should be of the same order. The combined likelihood of fits to the $D\pi$ mass, helicity and $w$ distributions for $D_S^0 + D_S^2$ hypothesis is higher than that for the $D_v^* + D_S^2$ combination by $2.8\sigma$. However, the present data sample cannot exclude the interpretation of this enhancement as a $D_v^*$ tail.

References

13. Used as a reference.
SEMILEPTONIC $B$ AND $D$ DECAYS — A REVIEW OF RECENT PROGRESS

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We present a review of semileptonic decays of $B$ and $D$ mesons, highlighting recent results from the $B$ factories. We discuss measurements of both inclusive and exclusive decays, measurements of the CKM quark-mixing matrix elements $|V_{cb}|$ and $|V_{ub}|$, studies of nonperturbative QCD effects, and a search for new physics effects using decays to $\tau$ leptons.

1 Introduction

Semileptonic decays provide an excellent laboratory in which to study electroweak physics, QCD, and to search for physics beyond the Standard Model. We present recent results on semileptonic $B$ and $D$ meson decays from the three $B$ factories, $B\bar{A}B\bar{R}$, Belle, and CLEO.

2 $|V_{cb}|$ and Heavy-quark Parameters from Inclusive $B$ Decays

The inclusive decay mode $B \rightarrow X_c \ell^- \nu_\ell$, where $X_c$ indicates any charmed hadronic system, can be used both to measure the CKM matrix element $|V_{cb}|$ and to study nonperturbative QCD effects of quarks bound inside hadrons. The differential decay rate for this process is described in Heavy Quark Effective Theory (HQET) as an expansion in terms of $\alpha_s$, whose effects are perturbatively calculable, and in the $b$ quark mass, $m_b$, whose effects are nonperturbative and must be measured in data. At second order in $1/m_b$, two nonperturbative parameters arise, corresponding to the kinetic energy and chromomagnetic moment of the $b$ quark in the $B$ meson and denoted $\mu_1^2$ and $\mu_2^2$, respectively; at third order in $1/m_b$, two further parameters arise, $\rho_3^{LS}$ and $\rho_3^D$. By measuring moments of the lepton energy spectrum and the $X_c$ mass spectrum in $B \rightarrow X_c \ell^- \nu_\ell$ decays and the photon energy spectrum in $B \rightarrow X_s \gamma$ decays — and by studying the variation of these moments as a function of a low-energy cut on the lepton (or photon, in the case of $B \rightarrow X_s \gamma$) energy — we can measure these nonperturbative heavy-quark expansion
parameters. By measuring the total rate of $B \to X_s \ell^- \nu_\ell$ decays, we can simultaneously extract the value of $|V_{cb}|$.

The Belle and BABAR Collaborations recently presented measurements of the moments of the $E_\ell$ and $m_X$ spectra. These measurements use a tagging technique where one of the two $B$ mesons in an $Y(4S) \to B \bar{B}$ event is fully reconstructed in a hadronic decay channel; by tagging one $B$ meson, the second $B$ can be reconstructed with reduced background and additional kinematic constraints, both of which are helpful when reconstructing decays with unobserved neutrinos. Corrections are applied to the observed kinematic variables $E_\ell$ and $m_X$ to account for finite detector resolution and the effects of unobserved particles. The Belle measurements use an unfolding technique, based on the Singular Valued Decomposition technique, while the BABAR analysis uses a set of calibration curves to make event-by-event corrections.

A global fit for $|V_{cb}|$ and the heavy-quark expansion parameters is shown in Figure 1. The average includes the recent measurements from Belle and BABAR, as well as older measurements from CLEO, CDF, and DELPHI, and includes up to the third $E_\ell$ moment, the third $m_X$ moment, and the second $E_\ell$ moment, all for a variety of lepton or photon energy cuts. The measured moments are highly correlated with one another taking into account the individual covariance matrices as well as a number of external constraints from theory and from other measurements. The measured value of $|V_{cb}|$ is $(42.04 \pm 0.34 \pm 0.59) \times 10^{-3}$, with a total error less than 2%, and the $b$ quark mass is measured as $(4.597 \pm 0.034)$ GeV$/c^2$, with an error less than 1%.

Inclusion of the $B \to X_s \gamma$ photon energy moments in this global fit is somewhat problematic, both from a theoretical and an experimental point of view. Theoretically, including these moments is difficult, in part because all calculations are model dependent to some degree, and in part because the operator product expansion must take into account non-local operators which are difficult to estimate. Additionally, the experimental results display some tension, with the $B \to X_s \gamma$ results pulling down the value of $m_b$ by about 1%. If the $B \to X_s \gamma$ moments are excluded from the fit, we instead obtain $|V_{cb}| = (41.85 \pm 0.38 \pm 0.59) \times 10^{-3}$ and $m_b = (4.660 \pm 0.053)$ GeV$/c^2$. While the effect on $|V_{cb}|$ is rather small, the effect of this change on the value of $|V_{ub}|$ is much larger, $\approx 10\%$; for this reason, the extraction of $|V_{ub}|$ presented below uses only the $B \to X_s \ell^- \nu_\ell$ moments.

3. $|V_{ub}|$ from Inclusive $B$ Decays

Precision measurement of $|V_{ub}|$ is one of the main goals of the $B$ factory physics program since, together with the angle $\beta$, $|V_{ub}|$ helps determine the apex of the Unitarity Triangle. The most precise measurements of $|V_{ub}|$ come from the inclusive $B \to X_u \ell^- \nu_\ell$ decay rate, which is proportional to $|V_{ub}|^2$.

The $B \to X_u \ell^- \nu_\ell$ decay rate is difficult to measure because background from $B \to X_c \ell^- \nu_\ell$ decays is 50 times larger than the signal. Measurements of $|V_{ub}|$ use cuts on kinematic variables — including the lepton energy, $m_X$, $q^2$, and $P_+ \equiv E_X - |p_X|$ — to suppress this $|V_{cb}|$ background, taking advantage of the fact that the $c$ quark is much heavier than the $u$. The partial decay rate in this restricted phase space is then extrapolated back to the full decay rate using theoretical models based on heavy-quark parameters which are determined from $B \to X_c \ell^- \nu_\ell$ decays as described above.

BABAR presented a measurement of $|V_{ub}|$ using three kinematic variables: $m_X$, $q^2$, and $P_+$. One $B$ meson is fully reconstructed and a high-momentum lepton is identified in the recoil. Combinatorial backgrounds are subtracted by fitting distributions of the tag $B$ mass in bins of the three kinematic variables, and a fit to the resulting kinematic distributions is used to distinguish $B \to X_u \ell^- \nu_\ell$ signal from the residual $B \to X_c \ell^- \nu_\ell$ events and other backgrounds. Several values of $|V_{ub}|$ are reported for different kinematic cuts and in different theoretical frameworks. A global average of inclusive $|V_{ub}|$ measurements, including this latest
Figure 1: Projection of a global fit for $|V_{cb}|$ and heavy-quark expansion parameters using moments measurements, showing the error ellipse in the $m_b - \mu^2_\pi$ plane. The ellipse is shown for three configurations of the fit: including all moments in the fit, including just the $B \to X_c \ell^- \bar{\nu}_\ell$ moments, and including just the $B \to X_s \gamma$ moments.

one and a similar analysis from Belle\textsuperscript{11}, is shown in Figure 2 for the BLNP framework; $|V_{ub}|$ is measured to be $(3.98 \pm 0.15 \pm 0.30) \times 10^{-3}$, with a total error of 8%, while similar results are obtained in the other theoretical frameworks\textsuperscript{9}.

4 Charm Semileptonic Decays and Form Factors

Studies of exclusive semileptonic decays, in which particular final state hadronic systems are selected, provide us with another approach to measuring CKM matrix elements and another way to help shed light on perturbative QCD processes. The dynamics of exclusive semileptonic decays are described by a set of form factors which are functions of the squared momentum transfer, $q^2$. A variety of theoretical techniques have been used to calculate these form factors\textsuperscript{12}. Decays of charm mesons provide a clean environment in which to measure the dynamics of semileptonic decay and to study these form factors; testing form factor models in the charm sector also leads to improved understanding of the form factors in the bottom sector, improving the extraction of $|V_{cb}|$ and $|V_{ub}|$.

CLEO-c presented recent results on the semileptonic $D$ decays $D \to \pi \ell^- \bar{\nu}_\ell$ and $D \to K \ell^- \bar{\nu}_\ell$ for both charged and neutral $D$ mesons\textsuperscript{13}. This analysis uses the missing four-momentum in the event to estimate the neutrino momentum, taking advantage of the good hermeticity of the detector. Signal events are required to have a squared missing mass, $m^2_{\text{miss}}$, consistent with zero, indicating that a single neutrino was undetected. Signals are further discriminated from background events using two kinematic variables, the mass and energy of the reconstructed $D$ candidate.

A fit is performed in bins of $q^2$ in order to measure the branching fractions and to extract
is also used to discriminate signal from $B$ 

$m_{\ell\ell}$ system and the beam energy. For the $\rightarrow 0$ study of the exclusive modes decay modes, such as $B B\pi$ factories. 

$h\pi\ell$ as well as to test form factor models in heavy-to-light meson $B\rightarrow$ normalizations, they measure $\rho\pi\ell$ from exclusive decays are orthogonal to those $\rightarrow \rho K\ell$, consistent with previous 

$\eta h\ell$ from Exclusive $B \rightarrow h\ell$ B factories. 

Figure 2: Global averages of inclusive $|V_{ub}|$ measurements (left) and the exclusive $B^0 \rightarrow \pi^- \ell^+\nu_\ell$ branching fraction (right), highlighting the consistency between many different measurement techniques as well as the precision obtained in recent years from the $B$ factories.

information about the form factors. The branching fractions measured are $B(D^0 \rightarrow \pi^+ \ell^-\nu_\ell) = (0.299 \pm 0.011 \pm 0.009)\%$, $B(D^- \rightarrow \pi^0 \ell^-\nu_\ell) = (0.373 \pm 0.022 \pm 0.013)\%$, $B(D^0 \rightarrow K^+ \ell^-\nu_\ell) = (3.56 \pm 0.03 \pm 0.09)\%$, and $B(D^- \rightarrow K^0\ell^-\nu_\ell) = (8.53 \pm 0.13 \pm 0.23)\%$, consistent with similar recent results from Belle and BABAR. The form factors for these decays are measured using both a model-independent series expansion and pole models. The expansion results are generally consistent with previous measurements. While the pole models also consistent with previous measurements, they only give a reasonable description of the data for unphysical parameter values. By using lattice QCD calculations of the form factor normalizations, they measure $|V_{cd}| = 0.217 \pm 0.009 \pm 0.004 \pm 0.023$ and $|V_{cs}| = 1.015 \pm 0.010 \pm 0.011 \pm 0.106$, in good agreement with previous measurements.

5 $|V_{ub}|$ from Exclusive $B$ Decays

Exclusive $b \rightarrow u$ decay modes, such as $B \rightarrow \pi\ell^-\nu_\ell$, $B \rightarrow \rho\ell^-\nu_\ell$, $B \rightarrow \omega\ell^-\nu_\ell$, and $B \rightarrow \eta(1^+)\ell^-\nu_\ell$, allow us to measure $|V_{ub}|$ as well as to test form factor models in heavy-to-light meson decays. The experimental and theoretical errors on $|V_{ub}|$ from exclusive decays are orthogonal to those in inclusive decays, making these modes complementary to the inclusive studies discussed above.

The CLEO Collaboration recently published a study of the exclusive modes $B \rightarrow h\ell^-\nu_\ell$, where $h = \{\pi^+/\pi^0/\rho^+/\rho^0/\omega/\eta/\eta'\}$. As in the previous analysis, the missing momentum in the event must be consistent with a single neutrino, which is then used to reconstruct the $B \rightarrow h\ell^-\nu_\ell$ candidate. Signal and background events are identified using two kinematic variables: $m_{h\ell^-\nu_\ell}$, the mass of the $h\ell^-\nu_\ell$ system after correcting for the neutrino energy resolution, and $\Delta E$, the difference between the observed energy of the $h\ell^-\nu_\ell$ system and the beam energy. For the $\rho$ and $\omega$ modes, the invariant mass $m_h$ of the $\rho$ or $\omega$ is also used to discriminate signal from background. A binned fit is performed to the joint distribution of $m_{h\ell^-\nu_\ell}$, $\Delta E$, $q^2$, $m_h$, and, for the $\rho$ mode, $\cos \theta_W$, the cosine of the angle between the lepton and the $W$ in the $B$ rest frame; this last variable is sensitive to the helicity of the $\rho$. 

Figure 2: Global averages of inclusive $|V_{ub}|$ measurements (left) and the exclusive $B^0 \rightarrow \pi^- \ell^+\nu_\ell$ branching fraction (right), highlighting the consistency between many different measurement techniques as well as the precision obtained in recent years from the $B$ factories.
Using isospin to combine the $\pi^+$ with $\pi^0$ results and the $\rho^+$, $\rho^0$, and $\omega$ results, they obtain $\mathcal{B}(\bar{B}^0 \to \pi^+ \ell^- \nu_\ell) = (1.31 \pm 0.15 \pm 0.11) \times 10^{-4}$ and $\mathcal{B}(\bar{B}^0 \to \rho^+ \ell^- \nu_\ell) = (2.93 \pm 0.37 \pm 0.37) \times 10^{-4}$, results which are among the most precise measurements to date. The branching fraction for $\bar{B}^0 \to \pi^+ \ell^- \nu_\ell$ can be compared to the world average\(^6\), which is shown in Figure 2. From the $\pi$ channel, they also measure $|V_{ub}| = (3.6 \pm 0.4 \pm 0.2^{+0.6}_{-0.4}) \times 10^{-3}$, comparable in precision to recent results from the \textit{BABAR} and Belle Collaborations\(^16\) and consistent with the world average. They find 3$\sigma$ evidence for the $\eta'$ mode with $\mathcal{B}(B^- \to \eta' \ell^- \nu_\ell) = (2.66 \pm 0.80 \pm 0.56) \times 10^{-4}$ and set a 90% upper limit $\mathcal{B}(B^- \to \eta' \ell^- \nu_\ell) < 1.01 \times 10^{-4}$; these results are consistent with a previous \textit{BABAR} upper limit at the 5% level, and may suggest a significant singlet contribution to the $\eta'$.}

6 \hspace{1cm} B \to D\ell^-\nu_\ell, \hspace{1cm} D^*\ell^-\nu_\ell, \hspace{1cm} \text{and} \hspace{1cm} D^{**}\ell^-\nu_\ell

Understanding the exclusive $b \to c$ semileptonic decays is another important part of the $B$ factory physics program, particularly since these modes have among the largest $B$ meson branching fractions. The dominant decay modes $B \to D\ell^-\nu_\ell$ and $B \to D^*\ell^-\nu_\ell$ make up about 70% of the total inclusive rate\(^6\), with the remaining 30% not yet well measured. These decay modes provide us with complementary measurements of $|V_{ub}|$ and allow us to study decay form factors and HQET. Additionally, these processes are backgrounds in many other analyses, so improved understanding of these decays will lead to improvements in extraction of $|V_{ub}|$ and $|V_{cb}|$.

The \textit{BABAR} Collaboration has presented a simultaneous measurement of the branching fractions $B \to D\ell^-\nu_\ell$, $B \to D^*\ell^-\nu_\ell$, $B \to D\pi^\pm\ell^-\nu_\ell$, and $B \to D^*\pi^\pm\ell^-\nu_\ell$, for both charged and neutral $B$ mesons\(^17\). Each of these modes is reconstructed in the recoil of a fully reconstructed $B$ meson, and signals are extracted using a fit to the $m^2_{\text{miss}}$ distribution, where correctly reconstructed events with just one missing neutrino peak at zero $m^2_{\text{miss}}$. Each of these eight branching fractions is the most precise measurement to date. The sum of these measurements, together with the inclusive branching fraction, suggests that $11 \pm 4\%$ of $B \to X_c \ell^-\nu_\ell$ decays are still unaccounted for, and may likely be due to $B \to D^{(*)}_1 n \pi \ell^-\nu_\ell$ decays with $n > 1$ pions in the final state.

Studies of the decays $B \to D^{**}\ell^-\nu_\ell$ (where $D^{**}$ means either a charm resonance heavier than the $D^*$ or a nonresonant $D^{(*)}_1 n \pi$ system) are interesting because, as mentioned above, the known exclusive decay modes do not saturate the inclusive decay rate, and $D^{**}$ is expected to make up most of the remainder. These decays are also interesting because of what is known as the 1/2–3/2 puzzle: HQET strongly favors production of resonances where the light quark has angular momentum $j_q = 3/2$ (the $D_1$ and $D_2^*$ states) over those with angular momentum $j_q = 1/2$ (the $D_0^*$ and $D'_1$), but experimental results\(^18\) suggest that the rates of the two angular momentum states are comparable.

Belle and \textit{BABAR} recently presented studies of $B \to D^{**}\ell^-\nu_\ell$ decays where the individual $D^{**}$ states are distinguished\(^19\). Both analyses identify a clean sample of $B \to D^{(*)}_1 \pi^\pm\ell^-\nu_\ell$ decays by reconstructing them in the recoil of a fully reconstructed $B$ meson and using $m^2_{\text{miss}}$ to identify signal events. A fit to the $D\pi$ and $D^*\pi$ mass spectra is used to disentangle the individual $D^{**}$ contributions, and the branching fractions are summarized in Table 1. The results of the two analyses are largely consistent with one another and with previous results. The branching fractions for the $j_q = 1/2$ states are of the same magnitude as the $j_q = 3/2$ states, confirming earlier results yet perpetuating the 1/2–3/2 puzzle in HQET. Neither measurement sees evidence for a nonresonant $B \to D^{(*)}_1 \pi^-\ell^-\nu_\ell$ state. The most significant difference between the two sets of results is in the $B \to D'_1 \ell^-\nu_\ell$ state. Belle sees no evidence for these decays and sets an upper limit, while \textit{BABAR}, with comparable sensitivity, sees a significant signal ($>6\sigma$). It is difficult to accommodate a large rate for the $B \to D'_1 \ell^-\nu_\ell$ state without a similarly large rate in $B \to D_0^* \ell^-\nu_\ell$, so further study of these modes will help to resolve this discrepancy.
Table 1: Measured product branching fractions $B(B \to D^{(*)} \ell^- \nu_{\ell}) \times B(D^{(*)} \to D^{(*)}\pi)$. Both analyses observe nonresonant $B \to D^{(*)}\pi \ell^- \nu_{\ell}$ yields consistent with zero.

<table>
<thead>
<tr>
<th>Mode</th>
<th>$B(B \to D^{(<em>)} \ell^- \nu_{\ell}) \times B(D^{(</em>)} \to D^{(*)}\pi)$ (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>Belle</td>
</tr>
<tr>
<td>$D^\pi$ invariant mass fit</td>
<td></td>
</tr>
<tr>
<td>$B^- \to D_1^0 \ell^- \nu_{\ell}$</td>
<td>0.24 ± 0.04 ± 0.06</td>
</tr>
<tr>
<td>$B^- \to D_2^0 \ell^- \nu_{\ell}$</td>
<td>0.22 ± 0.03 ± 0.04</td>
</tr>
<tr>
<td>$B^0 \to D_0^{*+} \ell^- \nu_{\ell}$</td>
<td>0.20 ± 0.07 ± 0.05</td>
</tr>
<tr>
<td>$B^0 \to D_2^{*-} \ell^- \nu_{\ell}$</td>
<td>0.22 ± 0.04 ± 0.04</td>
</tr>
<tr>
<td>$D^{\pi}\pi$ invariant mass fit</td>
<td></td>
</tr>
<tr>
<td>$B^- \to D_1^0 \ell^- \nu_{\ell}$</td>
<td>&lt; 0.07 (90% CL)</td>
</tr>
<tr>
<td>$B^- \to D_1^0 \ell^- \nu_{\ell}$</td>
<td>0.42 ± 0.07 ± 0.07</td>
</tr>
<tr>
<td>$B^- \to D_2^0 \ell^- \nu_{\ell}$</td>
<td>0.18 ± 0.06 ± 0.03</td>
</tr>
<tr>
<td>$B^0 \to D_1^{*+} \ell^- \nu_{\ell}$</td>
<td>&lt; 0.5 (90% CL)</td>
</tr>
<tr>
<td>$B^0 \to D_2^{*-} \ell^- \nu_{\ell}$</td>
<td>0.54 ± 0.19 ± 0.09</td>
</tr>
<tr>
<td>$B^0 \to D_2^{*-} \ell^- \nu_{\ell}$</td>
<td>&lt; 0.3 (90% CL)</td>
</tr>
</tbody>
</table>

7 $B \to D^{(*)}\tau^- \pi^\tau$

Semileptonic decays with $\tau$ leptons provide a new source of information on SM processes as well as a window into physics beyond the SM since the large $\tau$ mass gives sensitivity to decays mediated by a charged Higgs boson $^20$. Because the corresponding decays to light leptons have been studied and the form factors have been measured, theoretical predictions for the $\tau$ modes are quite clean, making these modes attractive probes of new physics. These decays are extremely challenging experimentally, however, due to the presence of multiple neutrinos in the final state.

Belle and BABAR recently presented the first results on exclusive semitauonic $B$ decays $^21$. Both experiments fully reconstruct one of the two $B$ mesons in the event and use the kinematic constraints to measure the missing four-momentum from the second $B$. Care must be taken to be sure that the decay products of both $B$ mesons are correctly reconstructed and account for all of the visible particles in the event, since mistakes tend to fake the missing momentum signature of signal events.

The Belle analysis reconstructs $B^0 \to D^{*+}\tau^- \pi^\tau$ with $\tau^- \to \ell^- \nu_{\ell} \nu_{\tau}$ and $\tau^- \to \pi^- \nu_{\tau}$ and requires events to have a large value of $X_{\text{miss}}$, a kinematic variable closely related to the missing mass. This cut preferentially selects events in which multiple neutrinos have escaped detection. The signal yield is then extracted by fitting the tag $B$ mass distribution, yielding the result $B(B^0 \to D^{*+}\tau^- \pi^\tau) = (2.02^{+0.34}_{-0.37} \pm 0.37)\%$.

The BABAR analysis reconstructs four modes, $B^- \to D^0\tau^- \pi^\tau$, $B^- \to D^{*0}\tau^- \pi^\tau$, $B^0 \to D^+\tau^- \pi^\tau$, and $B^0 \to D^{*+}\tau^- \pi^\tau$, with $\tau^- \to \ell^- \nu_{\ell} \nu_{\tau}$. The signal is extracted with a fit to the $m^2_{\text{miss}}$ and lepton momentum distributions (for signal events, this lepton is secondary), performed simultaneously in the $D^0$, $D^{*0}$, $D^+$, and $D^{*+}$ final states, as well as a set of control samples which simultaneously constrain background from $B \to D^{*+}\ell^- \nu_{\ell}$ decays. Combining results from charged and neutral $B$ modes, they obtain $B(B^0 \to D^{*+}\tau^- \pi^\tau) = (0.86 \pm 0.24 \pm 0.11 \pm 0.06)\%$ and $B(B^0 \to D^{*+}\tau^- \pi^\tau) = (1.52 \pm 0.31 \pm 0.10 \pm 0.05)\%$, where the $D^*$ result is consistent with that of Belle.

Both the Belle and BABAR results are about one standard deviation higher than the SM prediction. These measurements are statistically limited, however, and with increased statistics, studies of these modes are expected to add significant constraints to new physics models. In addition to the branching fractions, several other observables are sensitive to possible non-SM contributions, including $q^2$ distributions and $D^*$ and $\tau$ polarization $^20$, which would add to the
8 Conclusion

We have presented an overview of recent results in semileptonic decays from the $B$ factories. $|V_{ub}|$ has been measured with several different techniques and is now known to better than 10%, while $|V_{cb}|$ is now known to better than 2%. Both of these measurements are fundamental to the $B$ factory goal of overconstraining the Unitarity Triangle. Work is ongoing to understand the composition of the exclusive states which make up $B \rightarrow X_c \ell^- \nu_\ell$, particularly in disentangling the various $D^{**}$ contributions. New decay modes with $\tau$ leptons have been observed for the first time, opening up a new window into physics beyond the Standard Model.

References

1. Charge conjugate processes are implied throughout.
6. E. Barberio et al. (HFAG), arXiv:0704.3575 and online updates.
19. D. Liventsev et al. (Belle Collab.), Phys. Rev. D 77 091503 (2008); see also the contribution from D. Liventsev in these proceedings.

Charm and tau decays at $B$ factories

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We discuss recent results on charm and tau physics obtained by the Belle and BaBar collaborations. In the charm section we present measurements of $D^0 - \bar{D}^0$ mixing parameters, measurements searches for $CP$ violation in $D^0$ decays and a measurement of $D_s$ meson decay constant. In the tau section the recent results on lepton flavor violation in tau decays to three leptons or a lepton and a vector meson are discussed.

1 Introduction

The cross-sections for $c\bar{c}$ and $\tau$ pair production are very similar to the $b\bar{b}$ production cross-section at the $B$ factories. The Belle$^1$ and BaBar$^2$ detectors at the KEKB$^3$ and PEP-II colliders have accumulated together over 1 ab$^{-1}$ of data and therefore provide large samples and an excellent environment to study charm and $\tau$ decays.

2 $D^0 - \bar{D}^0$ mixing and search for $CP$ violation in $D^0$ decays

Particle-antiparticle mixing has been observed in several systems of neutral mesons: neutral kaons, $B_d$ and $B_s$ mesons. Last year at this conference the first evidence for $D^0 - \bar{D}^0$ mixing$^4,5$ was presented by both Belle and BaBar collaborations. As in the kaon and $B$-meson systems, the $D^0 - \bar{D}^0$ are produced in flavor eigenstates. The mixing occurs through weak interactions between the quarks and gives rise to two different mass eigenstates

$$|D_{1,2} >= p|D^0 > \pm q|\bar{D}^0 >,$$

where $|p|^2 + |q|^2 = 1$. The time evolution of flavor eigenstate is then given by

$$|D^0(t) >= \left[ |D^0 > \cosh \left( \frac{ix + y}{2} t \right) + \frac{q}{p} |\bar{D}^0 > \sinh \left( \frac{ix + y}{2} t \right) \right] \times e^{-\frac{1}{2}(1 + \frac{\Gamma}{m})t},$$

where $\Gamma$ is the width of the $D^0$ meson.
Table 1: The mixing parameter $y_{\text{CP}}$ and CP violating parameter $\Delta Y$ measured by BaBar using the ratios of lifetimes for the decays of $D^0$ mesons to $K^- K^+$, $\pi^- \pi^+$ and $K^- \pi^+$.

<table>
<thead>
<tr>
<th>Sample</th>
<th>$y_{\text{CP}}$</th>
<th>$\Delta Y$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$K^- K^+$</td>
<td>(+1.60 ± 0.46 ± 0.17)%</td>
<td>(-0.40 ± 0.44 ± 0.12)%</td>
</tr>
<tr>
<td>$\pi^- \pi^+$</td>
<td>(+0.46 ± 0.65 ± 0.25)%</td>
<td>(+0.05 ± 0.64 ± 0.32)%</td>
</tr>
<tr>
<td>Combined</td>
<td>(+1.24 ± 0.39 ± 0.13)%</td>
<td>(-0.26 ± 0.36 ± 0.08)%</td>
</tr>
</tbody>
</table>

where the two parameters that describe the $D^0 - \bar{D}^0$ mixing $x$ and $y$,

\[
x = \frac{m_1 - m_2}{\Gamma}, \quad (3)
\]

\[
y = \frac{\Gamma_1 - \Gamma_2}{2\Gamma}, \quad (4)
\]

\[
\Gamma = \frac{\Gamma_1 + \Gamma_2}{2} \quad (5)
\]

are the mass and width difference of the two mass eigenstates. In the Standard Model (SM), $D^0 - \bar{D}^0$ mixing is strongly GIM and CKM suppressed, and is dominated by long distance effects\(^6\). As the mixing rate is expected to be small within the SM, it is sensitive to the contribution of new, as of now unobserved processes and particles. The largest SM predictions for the parameters $x$ and $y$, which include the impact of long distance dynamics, are of order 1\(^%\)\(^6\).

$CP$ violating effects in decays of neutral $D$ meson system would appear as a difference in the partial decay widths of $D^0$ and $\bar{D}^0$ mesons decaying to a $CP$ eigenstate $f$

\[
A_{\text{CP}} = \frac{\Gamma(D^0 \rightarrow f) - \Gamma(\bar{D}^0 \rightarrow \bar{f})}{\Gamma(D^0 \rightarrow f) + \Gamma(\bar{D}^0 \rightarrow \bar{f})}. \quad (6)
\]

The contribution to the time-integrated asymmetry in neutral $D$ meson decays can be separated into three parts: direct $CP$ violation in decays to specific states, indirect $CP$ violation in $D^0 - \bar{D}^0$ mixing, and indirect $CP$ violation in interference between mixing and decay. Indirect $CP$ violation is to a good approximation predicted to be universal for amplitudes with final $CP$ eigenstates, but direct $CP$ violation can be non-universal depending on the specifics of the new physics. Within the SM the expected level of $CP$ violation is below the current experimental sensitivity\(^7\), therefore any positive signal would indicate physics beyond the SM.

BaBar measured $D^0 - \bar{D}^0$ mixing parameters using the ratios of lifetimes for the decays of neutral $D$ mesons to $CP$ even eigenstates $K^- K^+$ and $\pi^- \pi^+$ to the mixed-$CP$ state $K^- \pi^+$\(^8\). The ratio of lifetimes

\[
y_{\text{CP}} = \frac{\tau_{K \pi}}{\tau_{hh}} - 1, \quad h = K, \pi, \quad (7)
\]

corresponds in the limit of conserved $CP$ symmetry to the mixing parameter $y$ defined above. By measuring the lifetime difference of $D^0$ and $\bar{D}^0$ mesons decaying to $CP$ eigenstates the $CP$ violating parameter

\[
\Delta Y = \frac{\tau_{K \pi}}{<\tau_{hh}>} A_{\Gamma}, \quad A_{\Gamma} = \frac{\tau_{hh}(D^0) - \tau_{hh}(\bar{D}^0)}{\tau_{hh}(D^0) + \tau_{hh}(\bar{D}^0)} \quad (8)
\]

is measured. In the limit of $CP$ conservation $\Delta Y = 0$.

The $D^0$ meson is required to be produced in a $D^{*+} \rightarrow D^0 \pi^+$ decay\(^a\). This requirement suppresses the background and tags the flavor of neutral $D$ meson at the production with the charge of the pion. The $D^0$ lifetime is determined from an unbinned likelihood fit to the

\(^a\)Charge conjugation is implied throughout this paper.
reconstructed decay time and its estimated error, determined by a vertex-constrained combined fit to the $D^0$ decay and production vertices. The obtained value of $\gamma_{CP}$ given in Table 1, combined for both decay modes, represent evidence of $D^0 - \bar{D}^0$ mixing at the $3\sigma$ level. It confirms the lifetime ratio measurement made by Belle.\(^4\). The comparison of measured lifetimes for $D^0$ and $\bar{D}^0$ decaying to $CP$ eigenstates $K^-K^+$, $\pi^\pm \pi^\mp$ shows now evidence for $CP$ violation (Table 1).

Belle performed an improved search for $D^0 - \bar{D}^0$ mixing using semileptonic $D^0 \rightarrow K^{(*)-} \ell^+ \nu_\ell$ decays\(^5\), where the lepton is either an electron or a muon. Neutral $D$ mesons from $D^{*+} \rightarrow D^0 \pi^+$ decays are used and tagged at production by the charge of the pion. The mixing parameter,

$$R_M \equiv \frac{x^2 + y^2}{2} = \frac{N_{WS}}{N_{RS}},$$

(9)

is determined by measuring the numbers of reconstructed wrong (WS) and right sign (RS) events. The non-mixed decay results in a charge combination $\pi^+K^-\ell^+$ referred to as the RS charge combination while the mixing process results in a charge combination $\pi^+K^-\ell^-$ and is referred to as the WS charge combination. The reconstructed masses of $D^0$ and $D^{*+}$ candidates are smeared since the neutrino is not directly reconstructed. The RS and WS yields are determined from the fits to the RS and WS distributions of mass difference $\Delta M = M(K\nu\pi) - M(K\ell\nu)$, in which the uncertainty due to the neutrino four momentum cancels to a large extent. No significant WS signal is found in either the electron or muon samples and the most stringent experimental limit, obtained from semileptonic decays, on time time integrated mixing rate is given, $R_M < 6.1 \times 10^{-4}$ at $90\%$ C.L. The $R_M$ values obtained for each subsample, $e$ and $\mu$, are shown on Fig. 1.

The Belle and BaBar collaborations performed measurements searching for $CP$ violation in decays of neutral $D$ mesons to $K^-K^+$, $\pi^\mp \pi^\pm$\(^10\), $\pi^\pm \pi^\mp\pi^0$\(^11,12\) and $K^+K^-\pi^0$\(^12\). The main experimental challenge in these analyses is precise tagging of a neutral $D$ meson decaying to a $CP$ eigenstate. The flavor of the $D^0$ meson at production is tagged, as in the mixing analyses described above, by reconstructing $D^{*+} \rightarrow D^0 \pi^+$ decays. Beside the intrinsic asymmetry $A_{CP}$, defined by Eq. 6, there are two other contributions that create a difference in the numbers of reconstructed $D^0$ and $\bar{D}^0$ events. The first one is the forward-backward (FB) asymmetry in the production of $D^{*+}$ in $e^+e^- \rightarrow c\bar{c}$ arising from $\gamma-Z$ interference and higher order QED effects and is an odd function of the cosine of the $D^{*+}$ production polar angle in the center-of-mass system (CMS)\(^13\). The second one is the asymmetry in the reconstruction efficiencies of oppositely charged pions from $D^{*+}$ decays. The effect of the latter is evaluated and corrected for by measuring the relative detection efficiency for tagging pions using the $D^0 \rightarrow K^-\pi^+$ decays with and without flavor tag. CP violation would appear as an asymmetry in the $D^0 - \bar{D}^0$ yields independent of any kinematic variable. However, the reconstruction efficiency of the tagging pion is polar angle dependent, therefore the $CP$ asymmetry, $A_{CP} = \frac{N_{\pi^+} - N_{\pi^0}}{N_{\pi^+} + N_{\pi^0}}$, is measured.
in intervals of the cosine of the polar angle in the CMS. Any forward-backward asymmetry is canceled by averaging over symmetric intervals in the cosine of the polar angle in the CMS.

In Table 2 the measured $CP$ asymmetry by the BaBar and Belle Collaborations in $D^0 \to K^- K^+$ and $D^0 \to \pi^- \pi^+$ decays is given. No $CP$ violation is observed in either of the decay modes. The measurements are statistically limited. The main source of the systematic uncertainty is the statistics of the $D^0 \to K^- \pi^+$ samples, used to correct the charged pion reconstruction efficiency asymmetry and will thus also reduce with larger data samples.

The three-body decays $D^0 \to \pi^- \pi^+ \pi^0$, $K^- K^+ \pi^0$ proceed both via $CP$ eigenstates and flavor states, making it possible to probe $CP$ violation in both types of amplitudes and in the interference between them. Measuring interference effects in a Dalitz plot probes asymmetries in both the magnitudes and phases of the amplitudes, not simply interference between them. Measuring interference effects (Eq. 6) in $D^0 \to \pi^- \pi^+ \pi^0$ and $D^0 \to K^- K^+ \pi^0$ decays was evaluated using independent $D^{*+} \to D^0(K_S \pi^0) \pi^+$ data and Monte Carlo simulated samples at Belle, while in BaBar’s measurement it was evaluated using tagged and untagged data samples of $D^0 \to K^- \pi^+$ decays as described above. This difference explains the larger systematic uncertainty on measured $CP$ asymmetry from Belle. The phase-space integrated $CP$ asymmetry is insensitive to differences in the Dalitz plot shapes, so BaBar adopted three other approaches to search for $CP$ violation in $D^0 \to \pi^- \pi^+ \pi^0$, $K^- K^+ \pi^0$ decays. First they quantified differences between the $D^0$ and $\bar{D}^0$ Dalitz plots in two dimensions by plotting normalized residuals (shown in Figure 3)

$$\Delta = (n_{D^0} - Rn_{\bar{D}^0})/\sqrt{\sigma_{n_{D^0}}^2 + R^2\sigma_{n_{\bar{D}^0}}^2} \quad (10)$$

in the Dalitz plot area elements, and where $n$ denotes the number of events, $\sigma$ its uncertainty and $R$ is the efficiency corrected ratio. From the calculated $\chi^2/\nu = (\sum_{DP} \Delta^2)/\nu$ value, where $\nu$ is the number of Dalitz plot elements, the one-sided Gaussian confidence levels for consistency with no $CP$ violation are obtained: 32.8% for $\pi^- \pi^+ \pi^0$ and 16.6% for $K^- K^+ \pi^0$. In BaBar’s second approach differences in the angular moments of the $D^0$ and $\bar{D}^0$ intensity distributions
are looked for. The angular moments of the cosine of the helicity angle of the $D^0$ meson decay products reflect the spin and mass structure of intermediate resonant and nonresonant amplitudes. Similarly to the previous approach the one sided Gaussian confidence levels for consistency with no CP violation are obtained: 28.2% for the $\pi^+\pi^-$, 28.4% for the $\pi^+\pi^0$, 63.1% for the $K^+K^-$, and 23.8% for the $K^+\pi^0$ subsystems. In the third, model dependent approach, BaBar searched for CP violation in the amplitudes describing intermediate states in the $D^0$ and $\bar{D}^0$ decays. The Dalitz plot amplitude $A$ can be parametrized as a sum of amplitudes $A_r(s_+,s_-)$ for all relevant intermediate states $r$, each with a complex coefficient, i.e., $A = \sum_r a_r e^{i\phi_r} A_r(s_+,s_-)$, where $a_r$ and $\phi_r$ are real and $s_+$ and $s_-$ are the squared invariant masses of the pair of final state particles with +1 and −1 net charge. In the absence of CP violation the values of $a_r$ and $\phi_r$ are expected to be identical for $D^0$ and $\bar{D}^0$ decay. Comparison of amplitudes and relative phases, $a_r$ and $\phi_r$, obtained for $D^0$ and $\bar{D}^0$ decays showed, that the CP asymmetry in any amplitude, relative to that of the whole decay, is no larger than a few percent.

3 Measurement of $B(D_s^+ \to \ell^+\nu_\ell)$

One of the more important goals of particle physics is the precise measurement and understanding of the CKM matrix. To interpret results on $B$ meson decays, theoretical calculations of form factors and decay constants are often needed (usually based on lattice gauge theory). Decays of charmed hadrons in turn enable tests of the predictions for analogous quantities in the charm sector. It is necessary to have accurate measurements in the charm sector to check theoretical methods and predictions. In the SM the leptonic decays of mesons are mediated by a single virtual $W^\pm$ boson. The decay rate for e.g. $D_s^+ \to \ell^+\nu_\ell$ is given by

$$\Gamma(D_s^+ \to \ell^+\nu_\ell) = \frac{G_F^2}{8\pi} f_{D_s^+}^2 m_\ell^2 m_{D_s} (1 - \frac{m_\ell^2}{m_{D_s}^2})^2 |V_{cs}|^2,$$

(11)

where $G_F$ is the Fermi coupling constant, $V_{cs}$ is the corresponding CKM matrix element, $m_\ell$ and $m_{D_s}$ are the masses of the lepton and $D_s$ meson, respectively. The effects of the strong interaction are accounted for by the decay constant $f_{D_s}$. Since the decay rate is very small for
Figure 4: Recoil mass spectrum for $D_s$-tags for right-sign (left top) and wrong sign (left bottom) charge combinations of the $D$ meson and kaon. (Right) Spectrum of missing mass squared for $D_s^+ \rightarrow \mu^+ \nu_\mu$ candidates. The signal peaks at zero, the background shape in red is obtained by reconstructing $D_s^+ \rightarrow e^+ \nu_e$ decays, where no signal is expected due to helicity suppression.

Electrons due to helicity suppression and detection of $\tau$’s involves additional neutrinos, the muon mode is experimentally the most accessible one.

The analysis performed at Belle uses events of the type $e^+e^- \rightarrow D_s^0 D^{\pm,0} K^{\pm,0} X$, where $X$ can be any number of pions and up to one photon. The particles in the final state are divided into a tag and signal side. The tag side consists of a $D$ meson and a kaon in any charge combination and tags the flavor of the $D_s$ meson. The signal side is a $D_s^+$ decaying to $D_s \gamma$. Reconstructing the tag side, and allowing for any possible set of particles in $X$, the signal side is identified by reconstruction of the recoil mass as shown in Figure 4. Within this sample of tagged inclusive $D_s$ decays, decays of $D_s$ meson to muon and neutrino are selected by requiring another charged track that is identified as a muon and has the same charge as the $D_s$ candidate. The number of reconstructed $D_s^+ \rightarrow \mu^+ \nu_\mu$ decays is then determined from the fit to the recoil mass squared against all reconstructed particles, including the muon, as shown in Figure 4. Normalizing the number of reconstructed $D_s^+ \rightarrow \mu^+ \nu_\mu$ decays to the number of reconstructed tagged inclusive $D_s$ decays an absolute branching ratio is measured

$$B(D_s^+ \rightarrow \mu^+ \nu_\mu) = [6.44 \pm 0.76{\text{(stat)}} \pm 0.57{\text{(syst)}}] \times 10^{-3},$$

which is consistent with the world average and Babar’s and Cleo-c’s measurements. The obtained value of $f_{D_s}$ using Eq. 11 is

$$f_{D_s} = (275 \pm 16{\text{(stat)}} \pm 12{\text{(syst)}}) \text{ MeV.}$$

A simple average of the $D_s$ meson decay constant obtained from the cited measurements has an uncertainty of 11 MeV. Recently a lattice QCD calculation of significantly improved precision was performed, with preliminary result $f_{D_s} = (241 \pm 3) \text{ MeV}$. This value is somewhat lower than the experimental average and if it proves to be stable the comparison with the experimental results may point to some inconsistency between the two. More precise measurements are needed for a firm conclusion.

4 Search for lepton flavor violating $\tau$ decays

One of the currently most interesting questions in $\tau$ physics is whether there is a sizable lepton flavor violation (LFV) or not. LFV decays are expected even in the SM extended with the massive neutrinos, but the expected rate is very small and far beyond the reach of $B$ factories. Many extensions of the SM however, predict LFV $\tau \rightarrow \ell\ell\ell$ decays at the level of $10^{-10}$ to $10^{-7}$, which can be already probed at B factories with the current accumulated data. $B$ factories provide very clean environment for measurements searching for LFV $\tau$ decays. Candidate signal events are required to have 1-3 topology, where the $\tau$ on the signal side yields three charged
particles, while the second $\tau$ on the tag side yields one charged track. The event is easily divided into two hemispheres in the CMS. The signal side does not include any neutrinos in the final state, therefore signal events should peak at the nominal mass of the tau and at zero in the two dimensional distribution of the invariant mass versus energy difference.

Belle and BaBar reported improved upper limits on $\tau \to \ell \ell \ell$ branching ratios $^{22,23}$, where leptons in the final state are either electrons or muons, leading to six distinct decay modes: $e^- e^+ e^-$, $\mu^+ e^- e^-$, $\mu^- e^+ e^-$, $e^+ \mu^- \mu^-$, $e^- \mu^+ \mu^-$ and $\mu^- \mu^+ \mu^-$. In all cases the observed number of events in the signal region is consistent with the expected background. The improved upper limits on branching ratios, given in Table 4, are of order of $10^{-8}$ and they already restrict the parameter space of some beyond the SM models.

Belle reported improved upper limits on LFV $\tau$ decays to a lepton and vector meson, where the lepton is either an electron or a muon and vector meson is either $\phi$, $K^{*0}$, $\bar{K}^{*0}$ or $\rho^0$ $^{24}$. For the first same a search for $\tau \to \ell \omega$ ($\ell = e, \mu$) decays was performed by Belle and BaBar $^{24,25}$. No significant signal was observed in any of the studied decay modes. The improved upper limits on $\mathcal{B}(\tau \to \ell V^0)$ range from $5.9$–$10 \times 10^{-8}$ and are given in Table 5.

5 Conclusions

Only one year after the first observation of $D^0 - \bar{D}^0$ mixing, the mixing parameter $y_{CP}$ is known with relatively high precision. The current world averages of the mixing parameters $x$ and $y$ $^{26}$ lie at the upper edge of still uncertain theoretical expectations, at the level of 1%, therefore making it impossible to conclude whether $D^0 - \bar{D}^0$ mixing is a purely SM effect or receives contributions from new physics. $CP$ violation is expected to be small in the $D$ meson system, below the sensitivity of current experimental data. If large $CP$ violating phases are present in yet unknown processes the asymmetries could be increased to $\sim 1\%$. All measured $CP$ asymmetries
in $D^0$ decays observe no CP violation.

Further measurements of the $D_s$ meson decay constant are needed to resolve the discrepancy between the latest lattice QCD calculations and the experimental value.

The measurements searching for LFV tau decays are approaching the $10^{-8}$ level and already restrict the parameter space of many beyond the SM models.

References

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CONSTRaining NEW PHYSICS FROM \( D^0 - \bar{D}^0 \) MIXING

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I review constraints on possible New Physics interactions from \( D^0 - \bar{D}^0 \) mixing measurements. I consider the most general low energy effective Hamiltonian and include leading order QCD running of effective operators. I discuss constraints from an extensive list of popular New Physics models, each of which could be discovered at the LHC, that can generate these operators. In most of the scenarios, strong constraints that surpass those from other search techniques could be placed on the allowed parameter space using the existent evidence for observation of \( D \) meson mixing.

1 Introduction

Meson-antimeson mixing has traditionally been of importance because it is sensitive to heavy degrees of freedom that propagate in the underlying mixing amplitudes. Estimates of the charm quark and top quark mass scales were inferred from the observation of mixing in the \( K^0 \) and \( B_d \) systems, respectively, before these particles were discovered directly.

This success has motivated attempts to indirectly detect New Physics (NP) signals by comparing the observed meson mixing with predictions of the Standard Model (SM). \( K^0 - \bar{K}^0 \) mixing has historically placed stringent constraints on the parameter space of theories beyond the SM and provides an essential hurdle that must be passed in the construction of models with NP. The large mixing signal in the \( B_d \) and \( B_s \) systems, observed at the B-factories and the Tevatron collider, can be precisely described in terms of the SM alone, which makes the parameter spaces of various NP models increasingly constrained. These facts influenced theoretical and experimental studies of \( D^0 \) flavor oscillations, where the SM mixing rate is sufficiently small that the NP component might be able to compete. There has been a flurry of recent experimental activity.
regarding the detection of $D^0$-$\bar{D}^0$ mixing, which marks the first time Flavor Changing Neutral Currents (FCNC) have been observed in the charged $+2/3$ quark sector. With the potential window to discern large NP effects in the charm sector and the anticipated improved accuracy for future mixing measurements, the motivation for a comprehensive up-to-date theoretical analysis of New Physics contributions to $D$ meson mixing is compelling.

The phenomenon of meson-anti-meson mixing occurs in the presence of operators that change quark flavor by two units \(^1\). Those operators can be generated in both the Standard Model and many possible extensions of it. They produce off-diagonal terms in the meson-anti-meson mass matrix, so that the basis of flavor eigenstates no longer coincide with the basis of mass eigenstates. Those two bases, however, are related by a linear transformation,

$$|D_1\rangle = p|D^0\rangle \pm q|\bar{D}^0\rangle,$$  \hspace{1cm} (1)

where the complex parameters $p$ and $q$ are obtained from diagonalizing the $D^0 - \bar{D}^0$ mass matrix. Neglecting CP-violation leads to $p = q = 1/\sqrt{2}$. The mass and width splittings between those mass eigenstates are given by

$$x_D = \frac{m_1 - m_2}{\Gamma_D}, \quad y_D = \frac{\Gamma_1 - \Gamma_2}{2\Gamma_D}. \hspace{1cm} (2)$$

It is expected that $x_D$ and $y_D$ should be rather small in the Standard Model, which is usually attributed to the absence of superheavy quarks destroying Glashow-Iliopoulos-Maiani (GIM) cancellation. In Eq. (2), $\Gamma_D$ is the average width of the two neutral $D$ meson mass eigenstates. The quantities which are actually measured in most experimental determinations of the mass and width differences, $y_D^{(CP)}$, $x_D'$, and $y_D'$, are defined as

$$y_D^{(CP)} = y_D \cos \phi - x_D \sin \phi \left( \frac{A_m}{2} - A_{\text{prod}} \right),$$

$$x_D' = x_D \cos \delta_{K\pi} + y_D \sin \delta_{K\pi},$$

$$y_D' = y_D \cos \delta_{K\pi} - x_D \sin \delta_{K\pi}, \hspace{1cm} (3)$$

where $A_{\text{prod}} = (N_D - N_{\bar{D}}^0) / (N_D^0 + N_{\bar{D}}^0)$ is the so-called production asymmetry of $D^0$ and $\bar{D}^0$ (giving the relative weight of $D^0$ and $\bar{D}^0$ in the sample) and $\delta_{K\pi}$ is the strong phase difference between the Cabibbo favored and double Cabibbo suppressed amplitudes \(^2\), which is usually measured in $D \to K\pi$ transitions. In what follows we shall neglect CP-violating parameters $\phi$ and $A_m$. In this limit $y_D^{(CP)} = y_D$. Please see recent reviews \(^1,3,4\) for more complete analysis.

2 Experimental Constraints on Charm Mixing

The recent interest in $D^0$-$\bar{D}^0$ mixing started with the almost simultaneous observations by the BaBar \(^6\) and Belle \(^7\) collaborations of nonzero mixing signals at about the per cent level,

$$y_D' = (0.97 \pm 0.44 \pm 0.31) \cdot 10^{-2} \hspace{1cm} \text{(BaBar)},$$

$$y_D^{(CP)} = (1.31 \pm 0.32 \pm 0.25) \cdot 10^{-2} \hspace{1cm} \text{(Belle)}. \hspace{1cm} (4) \hspace{1cm} (5)$$

This was soon followed by the announcement by the Belle collaboration of mixing measurements from the Dalitz plot analyses of $D^0 \to K\pi^+\pi^-$. \(^8\)

$$x_D = (0.80 \pm 0.29 \pm 0.17) \cdot 10^{-2}, \quad y_D = (0.33 \pm 0.24 \pm 0.15) \cdot 10^{-2}. \hspace{1cm} (6)$$

A fit to the current database by the Heavy Flavor Averaging Group (HFAG) gives \(^5\)

$$x_D = 9.8_{-2.7}^{+2.6} \cdot 10^{-3}, \quad y_D = (7.5 \pm 1.8) \cdot 10^{-3}. \hspace{1cm} (7)$$
which is obtained assuming no CP-violation affecting mixing. It is important to note that the combined analysis of \(x_D\) and \(y_D\) excludes the ”no-mixing” point \(x_D = y_D = 0\) by \(6.7\sigma\). This fact adds confidence that charm mixing has indeed been observed. Then, a correct interpretation of the results is important. In addition, as with any rare low-energy transition, the question arises on how to use it to probe for physics beyond the Standard Model.

3 Standard Model ”background” in \(D^0 - \bar{D}^0\) mixing

Theoretical predictions for \(x_D\) and \(y_D\) obtained in the framework of the Standard Model historically span several orders of magnitude. I will not discuss predictions of the SM for the charm mixing rates here, instead referring the interested reader to recent reviews\(^{1,3,4}\). It might be advantageous to note that there are two approaches to describe \(D^0 - \bar{D}^0\) mixing, neither of which give very reliable results because \(m_c\) is in some sense intermediate between heavy and light.

The inclusive approach\(^{10,11}\) is based on the operator product expansion (OPE). In the formal limit \(m_c \gg \Lambda\) limit, where \(\Lambda\) is a scale characteristic of the strong interactions, \(x_D\) and \(y_D\) can be expanded in terms of matrix elements of local operators. The use of the OPE relies on local quark- hadron duality, and on \(\Lambda/m_c\) being small enough to allow a truncation of the series after the first few terms. This, however, is not realized in charm mixing, as the leading term in \(1/m_c\) is suppressed by four and six powers of the strange quark mass for \(x_D\) and \(y_D\) respectively. The parametrically-suppressed higher order terms in \(1/m_c\) can have less powers of \(m_s\), thus being more important numerically\(^{11}\). This results in reshuffling of the OPE series, making it a triple expansion in \(1/m_c, m_s,\) and \(\alpha_s\). The (numerically) leading term contains over twenty matrix elements of dimension-12, eight-quark operators, which are difficult to compute reliably. A naive power counting then yields \(x_D, y_D < 10^{-3}\). The exclusive approach\(^{12}\) sums over intermediate hadronic states. Since there are cancellations between states within a given \(SU(3)\) multiplet, one needs to know the contribution of each state with high precision. However, the \(D\) is not light enough that its decays are dominated by a few final states. In the absence of sufficiently precise data, one is forced to use some assumptions. Large effects in \(y_D\) appear for decays close to \(D\) threshold, where an analytic expansion in \(SU(3)_F\) violation is no longer possible. Thus, even though theoretical calculations of \(x_D\) and \(y_D\) are quite uncertain, the values \(x_D \sim y_D \sim 1\%\) are quite natural in the Standard Model\(^{13}\).

It then appears that experimental results of Eq. (7) are consistent with the SM predictions. Yet, those predictions are quite uncertain to be subtracted from the experimental data to precisely constrain possible NP contributions. In this situation the following approach can be taken. One can neglect the SM contribution altogether and assume that NP saturates the result reported by experimental collaborations. This way, however, only an upper bound on the NP parameters can be placed. A subtlety of this method of constraining the NP component of the mixing amplitude is related to the fact that the SM and NP contributions can have either the same or opposite signs. While the sign of the SM contribution cannot be calculated reliably due to hadronic uncertainties, \(x_D\) computed entirely within a given NP model can be determined rather precisely. This stems from the fact that NP contributions are generated by heavy degrees of freedom making short-distance OPE reliable. This means that only the part of parameter space of NP models that generate \(x_D\) of the same sign as observed experimentally can be reliably and unambiguously constrained.

4 New Physics contributions to \(D^0 - \bar{D}^0\) mixing

Any NP degree of freedom will generally be associated with a generic heavy mass scale \(M\), at which the NP interaction will be most naturally described. At the scale \(m_c\) of the charm
mass, this description will have been modified by the effects of QCD, which should be taken into account. In order to see how NP might affect the mixing amplitude, it is instructive to consider off-diagonal terms in the neutral D mass matrix,

$$\left( M - \frac{i}{2} \Gamma \right)_{12} = \frac{1}{2M_D} \langle D^0 | H^{\Delta C=-2} | D^0 \rangle + \frac{1}{2M_D} \sum_n \frac{\langle D^0 | H^{\Delta C=-1} | n \rangle \langle n | H^{\Delta C=-1} | D^0 \rangle}{M_D - E_n + i\epsilon}$$

(8)

where the first term contains $H^{\Delta C=-2}$, which is an effective $|\Delta C| = 2$ Hamiltonian, represented by a set of operators that are local at the $\mu \approx m_D$ scale. Note that a $b$-quark also gives a (negligible) contribution to this term. This term only affects $x_D$, but not $y_D$.

The second term in Eq. (8) is given by a double insertion of the effective $|\Delta C| = 1$ Hamiltonian $H^{\Delta C=-1}_w$. This term is believed to give dominant contribution to $D^0 - \bar{D}^0$ mixing in the Standard Model, affecting both $x$ and $y$. It is generally believed that NP cannot give any sizable contribution to this term, since $H^{\Delta C=-1}_w$ Hamiltonian also mediates non-leptonic $D$-decays, which should then also be affected by this NP contribution. I will show that there is a well-defined theoretical limit where NP contribution dominates lifetime difference $y_D$ and consider implications of this limit in "real world".

4.1 New Physics in $|\Delta C| = 1$ interactions.

Consider a non-leptonic $D^0$ decay amplitude, $A[D^0 \to n]$, which includes a small NP contribution, $A[D^0 \to n] = A_n^{(SM)} + A_n^{(NP)}$. Here, $A_n^{(NP)}$ is assumed to be smaller than the current experimental uncertainties on those decay rates. This ensures that NP effects cannot be seen in the current experimental analyses of non-leptonic D-decays. One can then write $y_D$ as

$$y_D \simeq \sum_n \frac{\rho_n}{\Gamma_D} A_n^{(SM)} A_n^{(SM)} + 2 \sum_n \frac{\rho_n}{\Gamma_D} A_n^{(NP)} A_n^{(SM)}$$

denoted as $\text{(SM)}$.

(9)

The first term of Eq. (schematic) represents the SM contribution to $y_D$. The SM contribution to $y_D$ is known to vanish in the limit of exact flavor $SU(3)$. Moreover, the first order correction is also absent, so the SM contribution arises only as a second order effect. This means that in the flavor $SU(3)$ limit the lifetime difference $y_D$ is dominated by the second term in Eq. (9), i.e. New Physics contributions, even if their contributions are tiny in the individual decay amplitudes. A calculation reveals that NP contribution to $y_D$ can be as large as several percent in R-parity-violating SUSY models or as small as $\sim 10^{-10}$ in the models with interactions mediated by charged Higgs particles.

This wide range of theoretical predictions can be explained by two observations. First, many NP affecting $|\Delta C| = 1$ transitions also affect $|\Delta B| = 1$ or $|\Delta S| = 1$ decays or kaon and B-meson mixings, which are tightly constrained. Second, a detailed look at a given NP model that can potentially affect $y_D$ reveals that the NP contribution itself can vanish in the flavor $SU(3)$ limit. For instance, the structure of the NP interaction might simply mimic the one of the SM. Effects like that can occur in some models with extra space dimensions. Also, the chiral structure of a low-energy effective lagrangian in a particular NP model could be such that the leading, mass-independent contribution vanishes exactly, as in a left-right model (LRM). Finally, the NP coupling might explicitly depend on the quark mass, as in a model with multiple Higgs doublets. However, most of these models feature second order $SU(3)$-breaking already at leading order in the $1/m_s$ expansion. This should be contrasted with the SM, where the leading order is suppressed by six powers of $m_s$ and term of order $m_s^2$ only appear as a $1/m_s^6$-order correction.

4.2 New Physics in $|\Delta C| = 2$ interactions.

Though the particles present in models with New Physics may not be produced in charm quark decays, their effects can nonetheless be seen in the form of effective operators generated by the
exchanges of these new particles. Even without specifying the form of these new interactions, we know that their effect is to introduce several $|\Delta C| = 2$ effective operators built out of the SM degrees of freedom.

By integrating out new degrees of freedom associated with new interactions at a scale $M$, we are left with an effective Hamiltonian written in the form of a series of operators of increasing dimension. Operator power counting then tells us the most important contributions are given by the operators of the lowest possible dimension, $d = 6$ in this case. This means that they must contain only quark degrees of freedom and no derivatives. Realizing this, we can write the complete basis of these effective operators, which can be done most conveniently in terms of chiral quark fields,

$$\langle f|\mathcal{H}_{NP}|i\rangle = G \sum_{i=1}^{8} C_i(\mu) \langle f|Q_i|i\rangle(\mu) \ ,$$

where the prefactor $G$ has the dimension of inverse-squared mass, the $C_i$ are dimensionless Wilson coefficients, and the $Q_i$ are the effective operators:

$$
\begin{align*}
Q_1 &= (\bar{u}_L\gamma_\mu c_L)(\bar{u}_L\gamma^\mu c_L) \\
Q_2 &= (\bar{u}_L\gamma_\mu c_L)(\bar{u}_R\gamma^\mu c_R) \\
Q_3 &= (\bar{u}_L\gamma_L)(\bar{u}_R\gamma_R) \\
Q_4 &= (\bar{u}_R\gamma_L)(\bar{u}_R\gamma_R) \\
Q_5 &= (\bar{u}_R\gamma_\mu c_L)(\bar{u}_R\gamma^\mu c_L) \\
Q_6 &= (\bar{u}_R\gamma_\mu c_R)(\bar{u}_R\gamma^\mu c_R) \\
Q_7 &= (\bar{u}_L\gamma_L)(\bar{u}_L\gamma_R) \\
Q_8 &= (\bar{u}_L\gamma_L)(\bar{u}_L\gamma_R) \\
\end{align*}
$$

In total, there are eight possible operator structures that exhaust the list of possible independent contributions to $|\Delta C| = 2$ transitions. Since these operators are generated at the scale $M$ where the New Physics is integrated out, a non-trivial operator mixing can occur when one takes into account renormalization group running of these operators between the scales $M$ and $\mu$, with $\mu$ being the scale where the hadronic matrix elements are computed. We shall work at the renormalization scale $\mu = m_c \simeq 1.3$ GeV. This evolution is determined by solving the RG equations obeyed by the Wilson coefficients,

$$\frac{d}{d\log \mu} \hat{C}(\mu) = \hat{\gamma}^T \hat{C}(\mu) \ ,$$

where $\hat{\gamma}$ represents the matrix of anomalous dimensions of the operators in Eq. (11). Due to the relatively simple structure of $\hat{\gamma}$, one can easily write the evolution of each Wilson coefficient in Eq. (10) from the New Physics scale $M$ down to the hadronic scale $\mu$, taking into account quark thresholds. Corresponding to each of the eight operators $\{Q_i\}$ ($i = 1, \ldots, 8$) is an RG factor $r_i(\mu, M)$. The first of these, $r_1(\mu, M)$, is given explicitly by

$$r_1(\mu, M) = \left( \frac{\alpha_s(M)}{\alpha_s(m_t)} \right)^{2/7} \left( \frac{\alpha_s(m_t)}{\alpha_s(m_b)} \right)^{6/23} \left( \frac{\alpha_s(m_b)}{\alpha_s(\mu)} \right)^{6/25} \ .$$

and the rest can be expressed in terms of $r_1(\mu, M)$ as

$$
\begin{align*}
r_2(\mu, M) &= [r_1(\mu, M)]^{1/2} \\
r_3(\mu, M) &= [r_1(\mu, M)]^{-4} \\
r_4(\mu, M) &= [r_1(\mu, M)]^{(1+\sqrt{241})/6} \\
r_5(\mu, M) &= [r_1(\mu, M)]^{(1-\sqrt{241})/6} \\
r_6(\mu, M) &= r_1(\mu, M) \\
r_7(\mu, M) &= r_4(\mu, M) \\
r_8(\mu, M) &= r_5(\mu, M) \ .
\end{align*}
$$

The RG factors are generally only weakly dependent on the NP scale $M$ since it is taken to be larger than the top quark mass, $m_t$, and the evolution of $\alpha_s$ is slow at these high mass scales. In Table 1, we display numerical values for the $r_i(\mu, M)$ with $M = 1, 2$ TeV and $\mu = m_c \simeq 1.3$ GeV. Here, we compute $\alpha_s$ using the one-loop evolution and matching expressions for perturbative consistency with the RG evolution of the effective Hamiltonian. A contribution to $D^0 - \bar{D}^0$
mixing from a particular NP model can be obtained by calculating matching conditions for the Wilson coefficients $C_i$ at the scale $M$, running their values down to $\mu$ and computing the relevant matrix elements of four-quark operators. A generic model of New Physics would then give the following contribution $x_D$: 

$$x_D^{NP} = G \frac{\mu_2 B_D m_D}{\Gamma_D} \left[ \frac{2}{3} [C_1(m_c) + C_6(m_c)] - \frac{5}{12} [C_4(m_c) + C_7(m_c)] + \frac{7}{12} C_3(m_c) \right.$$ 

$$\left. - \frac{5 C_2(m_c)}{6} + [C_5(m_c) + C_8(m_c)] \right] .$$

(15)

Here we simplified the result by assuming that all non-perturbative ('bag') parameters are equal to $B_D \approx 0.82$. The Wilson coefficients at the scale $\mu$ are related to the Wilson coefficients at the scale $M$ by renormalization group evolution,

$$C_1(m_c) = r_1(m_c, M) C_1(M) ,$$

$$C_2(m_c) = r_2(m_c, M) C_2(M) ,$$

$$C_3(m_c) = \frac{2}{3} [r_3(m_c, M) - r_3(m_c, M)] C_3(M) + r_3(m_c, M) C_3(M) ,$$

$$C_4(m_c) = \frac{8}{\sqrt{241}} [r_5(m_c, M) - r_4(m_c, M)] \left[ C_4(M) + \frac{15}{4} C_5(M) \right]$$

$$+ \frac{1}{2} [r_4(m_c, M) + r_5(m_c, M)] C_4(M) ,$$

$$C_5(m_c) = \frac{1}{8 \sqrt{241}} [r_4(m_c, M) - r_5(m_c, M)] [C_4(M) + 64 C_5(M)]$$

$$+ \frac{1}{2} [r_4(m_c, M) + r_5(m_c, M)] C_5(M) ,$$

$$C_6(m_c) = r_6(m_c, M) C_6(M) ,$$

$$C_7(m_c) = \frac{8}{\sqrt{241}} [r_8(m_c, M) - r_7(m_c, M)] \left[ C_7(M) + \frac{15}{4} C_8(M) \right]$$

$$+ \frac{1}{2} [r_7(m_c, M) + r_8(m_c, M)] C_7(M) ,$$

$$C_8(m_c) = \frac{1}{8 \sqrt{241}} [r_7(m_c, M) - r_8(m_c, M)] [C_7(M) + 64 C_8(M)]$$

$$+ \frac{1}{2} [r_7(m_c, M) + r_8(m_c, M)] C_8(M) .$$

(16)

A contribution of each particular NP model can then be studied using Eq. (15). Even before performing such an analysis, one can get some idea what energy scales can be probed by $D^0 - \overline{D^0}$ mixing. Setting $G = 1/M^2$ and $C_i(M) = 1$, we obtain $M \sim 10^3$ TeV. More realistic models can be probed in the region of several TeV, which is very relevant for LHC phenomenology applications.

A program described above has been recently executed\textsuperscript{15} for 21 well-motivated NP models, which will be actively studied at LHC. The results are presented in Table 2. As can be seen, out of 21 models considered, only four received no useful constraints from $D^0 - \overline{D^0}$ mixing. More informative exclusion plots can be found in that paper\textsuperscript{15} as well. It is interesting to note that

<table>
<thead>
<tr>
<th>$M$(TeV)</th>
<th>$r_1(m_c, M)$</th>
<th>$r_2(m_c, M)$</th>
<th>$r_3(m_c, M)$</th>
<th>$r_4(m_c, M)$</th>
<th>$r_5(m_c, M)$</th>
</tr>
</thead>
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<tr>
<td>1</td>
<td>0.72</td>
<td>0.85</td>
<td>3.7</td>
<td>0.41</td>
<td>2.2</td>
</tr>
<tr>
<td>2</td>
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<td>0.84</td>
<td>4.0</td>
<td>0.39</td>
<td>2.3</td>
</tr>
</tbody>
</table>

Table 1: Dependence of the RG factors on the heavy mass scale $M$.  

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<table>
<thead>
<tr>
<th>Model</th>
<th>Approximate Constraint</th>
</tr>
</thead>
<tbody>
<tr>
<td>Fourth Generation</td>
<td>$</td>
</tr>
<tr>
<td>$Q = -1/3$ Singlet Quark</td>
<td>$s_{2} \cdot m_{S} &lt; 0.27 \text{ (GeV)}$</td>
</tr>
<tr>
<td>$Q = +2/3$ Singlet Quark</td>
<td>$</td>
</tr>
<tr>
<td>Little Higgs</td>
<td>Tree: See entry for $Q = -1/3$ Singlet Quark</td>
</tr>
<tr>
<td>Generic $Z'$</td>
<td>Box: Parameter space can reach observed $x_{D}$</td>
</tr>
<tr>
<td>Family Symmetries</td>
<td>$M_{Z'}/C &gt; 2.2 \cdot 10^{3} \text{ TeV}$</td>
</tr>
<tr>
<td>Left-Right Symmetric</td>
<td>$m_{1}/f &gt; 1.2 \cdot 10^{3} \text{ TeV} \ (\text{with } m_{1}/m_{2} = 0.5)$</td>
</tr>
<tr>
<td>Alternate Left-Right Symmetric</td>
<td>No constraint</td>
</tr>
<tr>
<td>Vector Leptoquark Bosons</td>
<td>$M_{R} &gt; 1.2 \text{ TeV} \ (m_{D_{1}} = 0.5 \text{ TeV})$</td>
</tr>
<tr>
<td>Flavor Conserving Two-Higgs-Doublet</td>
<td>$(\Delta m/m_{D_{1}})/M_{R} &gt; 0.4 \text{ TeV}^{-1}$</td>
</tr>
<tr>
<td>Flavor Changing Neutral Higgs</td>
<td>$M_{VLQ} &gt; 55(\lambda_{PP}/0.1) \text{ TeV}$</td>
</tr>
<tr>
<td>FC Neutral Higgs (Cheng-Sher)</td>
<td>No constraint</td>
</tr>
<tr>
<td>Scalar Leptoquark Bosons</td>
<td>$m_{H}/C &gt; 2.4 \cdot 10^{3} \text{ TeV}$</td>
</tr>
<tr>
<td>Higgsless</td>
<td>$m_{H}/</td>
</tr>
<tr>
<td>Universal Extra Dimensions</td>
<td>See entry for RPV SUSY</td>
</tr>
<tr>
<td>Split Fermion</td>
<td>$M &gt; 100 \text{ TeV}$</td>
</tr>
<tr>
<td>Warped Geometries</td>
<td>No constraint</td>
</tr>
<tr>
<td>MSSM</td>
<td>$M/</td>
</tr>
<tr>
<td>SUSY Alignment</td>
<td>$M_{1} &gt; 3.5 \text{ TeV}$</td>
</tr>
<tr>
<td>Supersymmetry with RPV</td>
<td>$</td>
</tr>
<tr>
<td>Split Supersymmetry</td>
<td>$</td>
</tr>
<tr>
<td></td>
<td>$\tilde{m} &gt; 2 \text{ TeV}$</td>
</tr>
<tr>
<td></td>
<td>$\lambda_{12k}^{*}\lambda_{11k}/m_{d_{R,k}} &lt; 1.8 \cdot 10^{-3}/100 \text{ GeV}$</td>
</tr>
</tbody>
</table>

Table 2: Approximate constraints on NP models from $D^{0}$ mixing.

some models require large signals in the charm system if mixing and FCNCs in the strange and beauty systems are to be small (as in, for example, the SUSY alignment model\textsuperscript{16,17}).

5 Conclusions

I reviewed implications of recent measurement of $D^{0} - \bar{D}^{0}$ mixing rates for constraining models of New Physics. A majority of considered models received competitive constraints from $D^{0} - \bar{D}^{0}$ mixing measurements despite hadronic uncertainties that plague SM contributions. It should be noted that vast majority of predictions of NP models do not suffer from this uncertainty, and can be computed reliably, if lattice QCD community provides calculations of matrix elements of four-fermion operators Eq. (11).

Another possible manifestation of new physics interactions in the charm system is associated with the observation of (large) CP-violation\textsuperscript{1,4,18}. This is due to the fact that all quarks that build up the hadronic states in weak decays of charm mesons belong to the first two generations. Since $2 \times 2$ Cabbibo quark mixing matrix is real, no CP-violation is possible in the dominant tree-level diagrams which describe the decay amplitudes. CP-violating amplitudes can be introduced in the Standard Model by including penguin or box operators induced by virtual $b$-quarks. However, their contributions are strongly suppressed by the small combination of CKM matrix elements $V_{cb}V_{ub}^{*}$. It is thus widely believed that the observation of (large) CP violation in charm decays or mixing would be an unambiguous sign for New Physics.
Acknowledgments

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References

The CKM angle $\gamma/\phi_3$ - B-factories results review

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$\gamma/\phi_3$ is the less precisely known of the Unitarity Triangle angles. The general problematics of measurements of this parameter are discussed and recent experimental results from Babar and Belle are presented.

1 Measurements of the CKM angle $\gamma/\phi_3$

1.1 Introduction

In the Standard Model, $CP$ violation is described by the presence of an irreducible phase in the CKM matrix, the unitary matrix that relates the weak interaction with the mass eigenstates. The CKM can be written as:

$$V_{CKM} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix}$$

where $V_{q_1q_2}$ is the coupling related to the transition $q_1 \rightarrow q_2$. Many parametrizations exist in literature, we use here a generalization of the Wolfenstein parametrization\(^1\) as presented in\(^2\), where the four independent parameters are $\lambda$, $A$, $\bar{\rho}$ and $\bar{\eta}$ (where the latter is the $CP$ violating phase). The matrix is written:

$$V_{CKM} = \begin{pmatrix} 1 - \frac{\lambda^2}{2} & \lambda & A\lambda^3(\bar{\rho} - i\bar{\eta}) \\ -\lambda & 1 - \frac{\lambda^2}{2} & A\lambda^2 \\ A\lambda^3(1 - \bar{\rho} - i\bar{\eta}) & -A\lambda^2 & 1 \end{pmatrix} + O(\lambda^4)$$  \(1\)
The unitarity of the $V_{CKM}$ matrix implies several relations between its elements that can be represented as triangles in the $(\bar{\rho}, \bar{\eta})$ plane. We choose the relation $V^*_a V_b + V^*_c V_d + V^*_d V_e = 0$, whose elements can be determined by $B$ physics measurements. This triangle, represented in fig. 1, is particularly attractive from the experimental point of view, since it has all the sides of order $\lambda^3$. The angles of the triangle are called either $\alpha$, $\beta$ and $\gamma$ or $\phi_2$, $\phi_1$ and $\phi_3$, we adopt here the first notation.

In the Wolfenstein parametrization the only complex elements, up to terms of order $O(\lambda^5)$, are $V_{ub}$ and $V_{td}$ and the phases $\gamma$ and $\beta$ can be directly related to them. In particular, for $\gamma$ it can be written $V_{ub} = |V_{ub}| e^{-i\gamma}$. Several measurements, using different methods, constrains the weak phase $\gamma$ from the analyses of $B^+ \to (D^{(*)0}(D^{(*)0})K^{(*)0}$ and $B^0 \to (D^{(*)0}(D^{(*)0})K^{(*)0}$ decays, exploiting the interference between $b \to u$ and $b \to c$ transitions whose decay amplitudes will be proportional to the $V_{ub}$ and $V_{cb}$ elements respectively.

1.2 Phenomenology of $B \to DK$ decays

The amplitudes for the $B \to DK$ decays of interest can be expressed:

$$A(B^+ \to D^0 K^+) = V_{ub} V^*_c (T + C), \quad A(B^0 \to D^0 K^0) = V_{us} V^*_c C,$$

$$A(B^+ \to D^0 K^+) = V_{cs} V^*_u (\bar{C} + A), \quad A(B^0 \to D^0 K^0) = V_{cs} V^*_u \bar{C}.$$  \hspace{1cm} (2)

(3)

The $T$ parameter will account for a tree diagram, $C$ and $\bar{C}$ for color-suppressed diagrams and $A$ for an annihilation diagram. For the neutral $B \to DK$ decays, both the diagrams for the $b \to c$ and $b \to u$ transitions are color-suppressed and their amplitudes are described by the $C$ and $\bar{C}$ parameters respectively (see$^3$ for a complete treatment).

1.3 Measuring a phase

The idea of measuring a relative phase $\phi$ through the interference between two amplitudes $A_1$ and $A_2 e^{i\phi}$ connecting the same initial and final states is based on the fact that the decay rate between these two states is proportional to: $|A_1 + A_2 e^{i\phi}|^2 = A_1^2 + A_2^2 + 2A_1 A_2 \cos \phi$ and hence the interference term gives sensitivity to the relative phase $\phi$.

In fig. 2 we show an interference scheme for $B^+$ mesons decays giving sensitivity to $\gamma$. The $B^+$ can decay either to $D^0 K^+$ through a $b \to c$ transition or to $D^0 K^+$ through a $b \to u$ transition. If both the $D^0$ and the $\bar{D}^0$ decay to the same final state $f$, the study of the decay $B^+ \to f K^+$ gives sensitivity to the relative phase between the two decay amplitudes. The amplitude for $b \to c$ and $b \to u$ transitions can be written as $A(b \to u) \equiv |V_{ub}| e^{i\gamma} A_u e^{i\delta_u}$ and $A(b \to c) \equiv |V_{cb}| A_c e^{i\delta_c}$, where $A_u(c)$ and $\delta_u(c)$ are the absolute value and the phase of the strong interaction contribution to the amplitude. If the neutral $D$ decay is also considered, a term $A_D e^{i\delta_D}$ (or $A_{\bar{D}} e^{i\delta_{\bar{D}}}$) has to be included. In case of $B^+$, the interference term in the decay rate will be proportional to $\cos(\delta + \gamma)$, where $\delta = \delta_D - \delta_D + \delta_u - \delta_c$. A similar diagram can be

![Figure 1: Unitarity Triangle, represented in the $(\bar{\rho}, \bar{\eta})$ plane.](image-url)
drawn for the $CP$ conjugate decay ($B^- \to [f]K^-$), in this case the interference term will be proportional to $\cos(\delta - \gamma)$, since the strong interactions conserve $CP$.

The example shown in fig. 2 refers to the $B^+ \to \bar{D}^0(D^0)K^+$, but equivalent arguments can be done for all the $B^+ \to \bar{D}^{(*)0}(D^{(*)0})K^+$ and $B^- \to D^{(*)0}(\bar{D}^{(*)0})K^-$ as well as for the $B^0 \to D^{(*)0}(D^{(*)0})K^{(*)0}$ and $\bar{B}^0 \to D^{(*)0}(\bar{D}^{(*)0})K^{(*)0}$ decays.

A fundamental quantity in all the measurements of $\gamma$ is the parameter $r_B = |A(b \to u)|/|A(b \to c)|$. Being the absolute value of the ratio of the $b \to u$ to the $b \to c$ transition amplitudes, $r_B$ leads the sensitivity to $\gamma$ in each channel. Following the expressions for the decay amplitudes in 2, the $r_B$ ratio for charged $B \to DK$ channels can be written as:

$$r_B(D^0K^+) = \frac{|A(B^+ \to D^0K^+)|}{|A(B^+ \to \bar{D}^0K^+)|} = \frac{|V_{cs}V_{ub}|}{|V_{us}V_{cb}|} \left| \frac{\bar{C} + A}{T + C} \right|,$$

and, for neutral decays, as:

$$r_B(D^0K^0) = \frac{|A(B^0 \to D^0K^0)|}{|A(B^0 \to \bar{D}^0K^0)|} = \frac{|V_{cs}V_{ub}|}{|V_{us}V_{cb}|} \left| \frac{\bar{C}}{C} \right|.$$

In the expressions 4 and 5, the term $|V_{cs}V_{ub}|/|V_{us}V_{cb}|$ only depends on absolute values of CKM parameters and is known to be $\sqrt{\rho^2 + \eta^2} = 0.372 \pm 0.012^4$, while the terms depending on the hadronic parameters are not easily predictable. For simple numerical evaluation, the following assumption can be used: $|C|/|T| \approx 0.3$ and $|A|/|T| \approx 0.5^5$, and one would expect $r_B^{CH} \approx 0.1$ for the charged $B \to DK$ channels and $r_B^{NEWUT} \approx 0.4$ for the neutral $B \to DK$ ones.

The measurements of $\gamma$ are difficult because $b \to u$ transitions are strongly suppressed with respect to $b \to c$ ones, as described by $r_B$ ratios$^9$ and, as shown from the sketch in fig. 2, the unknowns in any $\gamma$ analysis are $\gamma$ itself, the $r_B$ ratio and a strong phase $\delta$. These are usually called polar coordinates. Some analyses make use of the cartesian coordinates, defined in terms of the polar coordinates as $x_\pm = r_B \cos(\delta \pm \gamma)$ and $y_\pm = r_B \sin(\delta \pm \gamma)$.

In the following, we denote $r_B^*$ and $\delta_B^*$ the amplitude ratio and strong phase relative to $B^+ \to D^{(*)0}(D^{(*)0})K^+$ decays. In case of a presence of a $K^*$ in the $B$ decay final state, as in the $B^- \to D^0(D^0)K^{*-}$ channel, the natural width of the $K^*$ resonance has to be taken into account and effective variables are used, following the formalism shown in 17. In case of the polar coordinates, these variables are $\gamma$ (which stays unchanged), $k$, $r_S$ and $\delta_S$ while, in case of the cartesian coordinates, they are called $x_{s\pm}$, $y_{s\pm}$.

1.4 Different experimental methods

There are several methods that aim to measure $\gamma$ in $B \to DK$ decays (all based on the strategy sketched in fig. 2) that differ because of the neutral $D$ final states $f$ they reconstruct and

$^a$It has to be stressed that the parameters $r_B$ are ratios between amplitudes, the ratio between number of events from $b \to u$ and $b \to c$ transitions will be proportional to $r_B^2$.  

4.5 Different experimental methods
consequently because of different experimental analysis techniques they use.

The Gronau London Wyler method

In the GLW method \(^6,7\), \(\gamma\) is measured from the study of \(B\) decays to \(D^0\)\(K\) final states, where \(D^0\) is a CP eigenstate (i.e. it is reconstructed in a CP eigenstate final state) with eigenvalues \(\pm 1\), defined starting from \(D^0\) and \(\bar{D}^0\), as \(|D^0_\pm| = \frac{1}{2}(|D^0| \pm |\bar{D}^0|)\).

The following observables are measured:

\[
R_{CP\pm} = \frac{\Gamma(B^+ \to D^0_{CP\pm} K^+) + \Gamma(B^- \to D^0_{CP\pm} K^-)}{\Gamma(B^+ \to D^0 K^+) + \Gamma(B^- \to D^0 K^-)} = 1 + r_B^2 \pm 2r_B \cos \gamma \cos \delta_B
\]

\[
A_{CP\pm} = \frac{\Gamma(B^+ \to D^0_{CP\pm} K^+) - \Gamma(B^- \to D^0_{CP\pm} K^-)}{\Gamma(B^+ \to D^0_{CP\pm} K^+) + \Gamma(B^- \to D^0_{CP\pm} K^-)} = \frac{\pm 2r_B \sin \gamma \sin \delta_B}{R_{CP\pm}}
\]

where \(\delta_B\) is the relative strong phase between the two \(B\) decay amplitudes.

In the GLW method, four observables, \(A_{CP\pm}\) and \(R_{CP\pm}\), are measured to constraint three unknowns, \(\gamma, \delta\) and \(r_B\). This method suffers of an irreducible four-fold ambiguity on the determination of the phases and, with the actual available statistics, is very useful in measuring \(r_B\), but has typically a low sensitivity to \(\gamma\).

The Adwood Dunietz Soni method

In the ADS method \(^8,9\), \(\gamma\) is measured from the study of \(B \to DK\) decays, where \(D\) mesons decay into non CP eigenstate final states. In this method the \(B\) meson is reconstructed in final states which can be reached in two ways: either through a favored \(b \to c B\) decay followed by a suppressed \(D\) decay (\(D^0 \to f\), or \(\bar{D}^0 \to \bar{f}\)), or through a suppressed \(b \to u B\) decay followed by a favored \(D\) decay (\(D^0 \to f\) or \(\bar{D}^0 \to \bar{f}\)). In this way the two amplitudes are comparable and one can expect larger interference terms.

In the ADS method, one measures the observables:

\[
R_{ADS} = \frac{\Gamma(B^+ \to \bar{f}K^+) + \Gamma(B^- \to fK^-)}{\Gamma(B^+ \to fK^+) + \Gamma(B^- \to \bar{f}K^-)} = r_D^2 + r_B^2 + 2r_B r_D \cos\gamma \cos(\delta_D + \delta_D)
\]

\[
A_{ADS} = \frac{\Gamma(B^+ \to \bar{f}K^+) - \Gamma(B^- \to fK^-)}{\Gamma(B^+ \to fK^+) + \Gamma(B^- \to \bar{f}K^-)} = r_B r_D [\cos(\delta + \gamma) + \cos(\delta - \gamma)] / R_{ADS}.
\]

Here \(\delta_D\) is the relative strong phase between the favored and suppressed \(D\) decay amplitudes, and \(r_D\) is the ratio between the absolute values of their amplitudes \(r_D = |A(D^0 \to f)| / |A(D^0 \to \bar{f})|\).

This method is very useful in measuring \(r_B\), but normally it has very low sensitivity to \(\gamma\).

The Giri Grossman Soffer Zupan method

In this method \(^10\), usually called Dalitz method, \(\gamma\) is measured from the \(B \to DK\) decays with the \(D\) decaying to multi-body CP eigenstate final states. Multi-body decays are usually described by the isobar model, in which the decay amplitude is written as a sum of amplitudes with quasi two-body intermediate states and determined on independent neutral \(D\) samples.

This information is used in input to the Dalitz analyses (that directly extracts form data \(\gamma, r_B\) and \(\delta\) or the polar coordinates) where the complete and rich structure of the multi-body \(D\) decay is exploited and detectable interference terms are expected because of the presence of different strong phases. This method is indeed very powerful and it is the one that gives the best error on the weak phase \(\gamma\).
2 Common experimental techniques

We present here the results obtained by the two B-factories experiments: Babar at the PEP-II asymmetric-energy $e^+e^-$ collider, located at the Stanford Linear Accelerator Center (USA) and Belle at the KEK asymmetric-energy $e^+e^-$ collider, located in Tsukuba (Japan). All the analyses presented reconstruct exclusively $B$ decays and make use of some common techniques.

The $B$ mesons are characterized by two almost independent kinematic variables: the beam-energy substituted mass $m_{ES}(M_{bc}) \equiv \sqrt{(E_0^* / 2 + \mathbf{p}_0 \cdot \mathbf{p}_B)^2 / E_0^* - p_B^2}$ and the energy difference $\Delta E \equiv E_B^* - E_0^*/2$, where $E$ and $p$ are the energy and the momentum respectively, the subscript $s$ refers to the candidate $B$ and 0 refers to the $e^+e^-$ system respectively and the asterisk denotes the $e^+e^-$ CM frame.

Since both PEP-II and KEK $e^+e^-$ collide at $\sqrt{s} = M(\Upsilon(4S))$, the $\Upsilon(4S)$ resonance is produced almost at rest in the $e^+e^-$ center of mass frame. Given the values of the masses of the $\Upsilon(4S)$ and of the $B$ mesons, the latter have a very low residual momentum in the $e^+e^-$ center of mass frame. On the other hand, in case of $e^+e^- \rightarrow q\bar{q}$ events, with $q = u, d, s, c$ (called continuum events), the two quarks are produced with large momentum and for this reason, these events have a jet-like spatial shape, different from the spherically distributed one for $B\bar{B}$ events.

Several variables account for these differences and are used in the analyses to fight continuum background, which is typically the main source of background to these analyses.

3 Experimental results on the charged $B$ decays

We present here the recent results on $\gamma$ from Babar and Belle, using the different methods.

3.1 Analyses using the GLW method

We report on the update of the GLW analysis\textsuperscript{12} of $B^- \rightarrow D^0 K^-$, with $D^0 \rightarrow K^+ K^-, \pi^+\pi^-$, $K_S \pi^0$ and $K_S \omega$, using 383 $10^6$ $B\bar{B}$ pairs collected with the Babar detector. In this analysis, after a cut on $m_{ES}$ and on a combination of event shape variables, the observables are extracted using a maximum likelihood fit to the variables $\Delta E$ and the Cerenkov angle of the charged $K$ produced in the charged $B$ decay.

The results obtained for the direct $CP$ asymmetries and the ratios are the following:

\[
R_{CP^+} = 1.06 \pm 0.10 \pm 0.05, \quad A_{CP^+} = 0.27 \pm 0.09 \pm 0.04,
\]
\[
R_{CP^-} = 1.03 \pm 0.10 \pm 0.05, \quad A_{CP^-} = -0.09 \pm 0.09 \pm 0.02,
\]

where the first error is statistical and the second one is systematic. For the first time for a GLW analysis, the results are extracted from data also in terms of the cartesian coordinates:

\[
x_+ = -0.09 \pm 0.05 \pm 0.02,
\]
\[
x_- = +0.10 \pm 0.05 \pm 0.03,
\]
\[
r^2 = +0.05 \pm 0.07 \pm 0.03,
\]

where the first error is statistical and the second one is systematic.

The uncertainties on $A_{CP\pm}$ ($R_{CP\pm}$) are smaller by a factor of 0.7 (0.9) and 0.6 (0.6) than the previous Babar\textsuperscript{13} and Belle\textsuperscript{14} measurements, respectively.

\[\text{the } K^-\pi^+ \text{ mode is also reconstructed for normalization} \]
3.2 Analyses using the ADS method

We report on the update of the ADS analysis of $B^- \to D^0 K^-$, with $D^0 \to K^- \pi^+$ using 657 $10^6$ $B\bar{B}$ pairs collected with the Belle detector. In this analysis, after a cut on $m_{ES}$ and on a combination of event shape variables, the observables are extracted using a maximum likelihood fit to the variable $\Delta E$, giving the following results:

$$R_{ADS} = (8.0^{+6.3+2.0}_{-5.7-2.8}) \times 10^{-3}, \quad A_{ADS} = -0.13^{+0.97}_{-0.88} \pm 0.26,$$

where the first error is statistical and the second one is systematic.

The results obtained for $R_{ADS}$ show that no evidence of $b \to u$ transition is found, even with the very high statistics used. This result implies an upper limit on the ratio $r_B, r_B < 0.19 \ 90\%$ C.L. This result on $r_B$ is consistent with the previous Belle and Babar analyses and confirms the expectation for a small value of $r_B$ ($r_B \sim 0.1$) in charged $B \to DK$ decays.

3.3 Analyses using the GGSZ method

Both the Babar and Belle collaboration have presented at this conference new results using Dalitz technique, that strongly improve the precision on the determination of $\gamma$.

We first report on a new Dalitz analysis of $B^- \to D^0 K^-$ and $B^- \to D^{*0} K^-$, that for the first time uses neutral $D$ reconstructed into the final state $D^0 \to K_s K^+ K^-$ and on the update of the Dalitz analysis of $B^- \to D^0 K^-, B^- \to D^{*0} K^-$ and $B^- \to D^0 K^{*-}$, with $D^0 \to K_S \pi^+ \pi^-$ using 383 $10^6$ $B\bar{B}$ pairs collected with the Babar detector. In this analysis, $m_{ES}$, $\Delta E$ and a combination of event shape variables are used in the maximum likelihood fit to extract the number of signal and background events and then a $CP$ fit is performed to extract the cartesian coordinates for the three channels, $B^- \to D^0 K^-, B^- \to D^{*0} K^-$ and $B^- \to D^0 K^{*-}$. In the $CP$ fit, the $D$ Dalitz distribution, for $D^0 \to K_S \pi^+ \pi^-$ and $D^0 \to K_s K^+ K^-$, as they are determined on independent data samples, are used as an input. The results for the cartesian coordinates are shown in tab. 1, for the three analyzed channels (in the tables, the symbol $D^0$ indicates either a $D^0$ or a $D^{*0}$). The first error is statistical, the second is experimental systematic uncertainty and the third is the systematic uncertainty associated with the Dalitz models.

<table>
<thead>
<tr>
<th>Parameters</th>
<th>$B^- \to D^0 K^-$</th>
<th>$B^- \to D^{*0} K^-$</th>
<th>$B^- \to D^0 K^{*-}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$x_-, x^+,$ $x_+$</td>
<td>$0.090 \pm 0.043 \pm 0.015 \pm 0.011$</td>
<td>$-0.111 \pm 0.069 \pm 0.014 \pm 0.004$</td>
<td>$0.115 \pm 0.138 \pm 0.039 \pm 0.014$</td>
</tr>
<tr>
<td>$y_-, y^+, y_+$</td>
<td>$0.053 \pm 0.056 \pm 0.007 \pm 0.015$</td>
<td>$-0.051 \pm 0.080 \pm 0.009 \pm 0.010$</td>
<td>$0.226 \pm 0.142 \pm 0.058 \pm 0.011$</td>
</tr>
<tr>
<td>$x_-, x^+, x_+$</td>
<td>$-0.067 \pm 0.043 \pm 0.014 \pm 0.011$</td>
<td>$0.137 \pm 0.068 \pm 0.014 \pm 0.005$</td>
<td>$-0.113 \pm 0.107 \pm 0.028 \pm 0.018$</td>
</tr>
<tr>
<td>$y_-, y^+, y_+$</td>
<td>$-0.015 \pm 0.055 \pm 0.006 \pm 0.008$</td>
<td>$0.080 \pm 0.102 \pm 0.010 \pm 0.012$</td>
<td>$0.125 \pm 0.139 \pm 0.051 \pm 0.010$</td>
</tr>
</tbody>
</table>

Table 1: $CP$-violating parameters $x_{\pm}^{(s)}, y_{\pm}^{(s)}, x_{\pm},$ and $y_{\pm}$, as obtained from the $CP$ fit.

Using a frequentist analysis, the experimental results for $x_{\pm}^{(s)}, y_{\pm}^{(s)}, x_{\pm}$, and $y_{\pm}$ are interpreted in terms of the weak phase $\gamma$, the amplitude ratios $r_B$, $r_B^*$, and $r_S$, and the strong phases $\delta_B$, $\delta_B^*$, and $\delta_S$, giving $\gamma = (76 \pm 22 \pm 5 \pm 5) (\text{mod } 180^\circ)$, $r_B = 0.086 \pm 0.035 \pm 0.010 \pm 0.011$, $r_B^* = 0.135 \pm 0.051 \pm 0.011 \pm 0.005$, $kr_S = 0.163^{+0.088}_{-0.105} \pm 0.037 \pm 0.021 \ \delta_B = \left(109^{+29}_{-31} \pm 4 \pm 7\right) (\text{mod } 180^\circ)$, $\delta_B^* = \left(-63^{+28}_{-30} \pm 5 \pm 4\right) (\text{mod } 180^\circ)$, and $\delta_S = \left(104^{+45}_{-41} \pm 17 \pm 5\right)$. The first error is statistical, the second is the experimental systematic uncertainty and the third reflects the uncertainty on the $D$ decay Dalitz models. The results for $\gamma$ and the ratios $r_B$, $r_B^*$ and $r_S$ are shown in fig. 3.

We also report on the update of the Dalitz analysis of $B^- \to D^0 K^-$ and $B^- \to D^{*0} K^-$ ($D^{*0} \to D^0 \pi^0$), with $D^0 \to K_S \pi^+ \pi^-$ using 635 $10^6$ $B\bar{B}$ pairs collected with the Belle detector. In this analysis, $M_{bc}, \Delta E$ and a combination of event shape variables are used in the maximum
Figure 3: [Babar Dalitz analysis] $\alpha = 1 - \text{CL}$ as a function of $\gamma$ (left plot) and of $r_B$, $r_B^*$ and $r_S$ (right plot) for $B^- \to D^0 K^-$, $B^- \to D^{*0} K^-$, and $B^- \to D^0 K^*$ decays separately, and their combination, including statistical and systematic uncertainties and their correlations. The dashed (upper) and dotted (lower) horizontal lines correspond to the one- and two-standard deviation intervals, respectively.

Using a frequentist analysis, the experimental results for $x_\pm^{(*)}$ and $y_\pm^{(*)}$ are interpreted in terms of the weak phase $\gamma$, the amplitude ratios $r_B$, $r_B^*$ and the strong phases $\delta_B$, $\delta_B^*$, giving $\gamma = \left(76^{+12}_{-13} \pm 4 \pm 9\right)^\circ \text{ (mod 180$^\circ$)}$, $\delta_B = \left(136^{+14}_{-16} \pm 4 \pm 23\right)^\circ \text{ (mod 180$^\circ$)}$, $\delta_B^* = \left(343^{+20}_{-22} \pm 4 \pm 23\right)^\circ \text{ (mod 180$^\circ$)}$, $r_B = 0.16 \pm 0.04 \pm 0.01 \pm 0.05$ and $r_B^* = 0.21 \pm 0.08 \pm 0.02 \pm 0.05$. The first error is statistical, the second is the experimental systematic uncertainty on independent data samples, is used as an input. The results are shown in tab. 2, where the second reflects the experimental systematic uncertainty. The uncertainty associated with the Dalitz model is not shown and it is assumed to be equal to the one evaluated in the previous analysis by Belle collaboration$^{18}$.

Using a frequentist analysis, the experimental results for $x_\pm^{(*)}$ and $y_\pm^{(*)}$ are interpreted in terms of the weak phase $\gamma$, the amplitude ratios $r_B$, $r_B^*$ and the strong phases $\delta_B$, $\delta_B^*$, giving $\gamma = \left(76^{+12}_{-13} \pm 4 \pm 9\right)^\circ \text{ (mod 180$^\circ$)}$, $\delta_B = \left(136^{+14}_{-16} \pm 4 \pm 23\right)^\circ \text{ (mod 180$^\circ$)}$, $\delta_B^* = \left(343^{+20}_{-22} \pm 4 \pm 23\right)^\circ \text{ (mod 180$^\circ$)}$, $r_B = 0.16 \pm 0.04 \pm 0.01 \pm 0.05$ and $r_B^* = 0.21 \pm 0.08 \pm 0.02 \pm 0.05$. The first error is statistical, the second is the experimental systematic uncertainty and the third reflects the uncertainty on the $D$ decay Dalitz model. It can be noticed that this analysis finds slightly higher values for the $r_B$ and $r_B^*$ ratios with respect to the Babar analysis, which explains the smaller statistical errors on $\gamma$, also if the precision on the cartesian coordinates is similar. The results for $\gamma$ and the ratios $r_B$ and $r_B^*$ are shown in fig. 4.

Table 2: $CP$-violating parameters $x_\pm^{(*)}$ and $y_\pm^{(*)}$, as obtained from the $CP$ fit.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>$B^- \to D^0 K^-$</th>
<th>$B^- \to D^{*0} K^-$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$x_-$</td>
<td>$+0.105 \pm 0.047 \pm 0.011$</td>
<td>$+0.024 \pm 0.140 \pm 0.018$</td>
</tr>
<tr>
<td>$y_-$</td>
<td>$+0.177 \pm 0.060 \pm 0.018$</td>
<td>$-0.243 \pm 0.137 \pm 0.022$</td>
</tr>
<tr>
<td>$x_+$</td>
<td>$-0.107 \pm 0.043 \pm 0.011$</td>
<td>$+0.133 \pm 0.083 \pm 0.018$</td>
</tr>
<tr>
<td>$y_\pm$</td>
<td>$-0.067 \pm 0.059 \pm 0.018$</td>
<td>$+0.130 \pm 0.120 \pm 0.022$</td>
</tr>
</tbody>
</table>

4 Experimental results on the neutral $B$ decays

Lately, within the Babar collaboration, there have been efforts to constrain $\gamma$ and related quantities from the study of neutral $B \to DK$ decays. As already discussed, the $r_B$ ratios in these channels are expected to be higher than in the charged ones, hence giving higher sensitivity to $\gamma$. 

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We first report on a new Dalitz analysis\(^{21}\) of \(B^0 \rightarrow D^0 K^{*0}\), with \(K^{*0} \rightarrow K^+\pi^-\) and \(D^0 \rightarrow K_S\pi^+\pi^-\) using 371 \(10^6\) \(B\bar{B}\) pairs collected with the Babar detector. In this analysis, \(m_{ES}\) and a combination of event shape variables are used in the maximum likelihood fit to extract the number of signal and background events and then a \(CP\) fit is performed. A likelihood scan in polar coordinates \((\gamma, \delta^0_S, r^0_S)\) is extracted from data and combined with an external information on \(r^0_S\)^{20}. The results obtained are shown in tab. 3, where the first error is statistical, the second is the experimental systematic uncertainty and the third reflects the uncertainty on the \(D\) decay Dalitz model.

\[
\begin{array}{|c|c|}
\hline
\text{Parameters} & \text{Value} \\
\hline
\gamma [^\circ] & 162 \pm 55 \pm 6.5 \text{ (mod } 180[^\circ]\text{)} \\
\delta^0_S [^\circ] & 62 \pm 55 \pm 3.1 \pm 15.8 \text{ (mod } 180[^\circ]\text{)} \\
\hline
\end{array}
\]

Table 3: Results for \(\gamma\), \(\delta^0_S\) and \(r^0_S\), as obtained from the \(CP\) fit.

We also report on a new time-dependent Dalitz plot analysis\(^{22}\) of \(B^0 \rightarrow D^- K^0\pi^+\) using 347 \(10^6\) \(B\bar{B}\) pairs collected with the Babar detector. This analysis studies the interference between \(b \rightarrow u\) and \(b \rightarrow c\) transitions through the \(B\) mesons mixing and hence gives sensitivity to the combination of CKM weak phases \(2\beta + \gamma\). In this analysis, \(m_{ES}\), \(\Delta E\) and a combination of event shape variables are used in the maximum likelihood fit to extract the number of signal and background events and then a time-dependent fit to the neutral \(B\) Dalitz distribution is performed to extract \(2\beta + \gamma\). In this fit, the ratio \(r^0_B\) is assumed to be \(r^0_B = 0.3\) and the effect of this assumption is taken into account in the systematics evaluation by varying this ratio of \(\pm 0.1\). The result obtained for \(2\beta + \gamma\) is the following:

\[
2\beta + \gamma = (83 \pm 53 \pm 20)^[^\circ] \text{ (mod } 180[^\circ]\text{)},
\]

where the first error is statistical and the second one is systematic.

5 Combined results and conclusions

From all the available measurements, including the new ones presented here, the knowledge of \(\gamma\), according to the combination performed by the UTfit collaboration, is \(\gamma = (80 \pm 13[^\circ]).\)
The combined results obtained for the other quantities are \( r_B = 0.10 \pm 0.02, \ r_B^* = 0.09 \pm 0.04, \ r_S = 0.13 \pm 0.09, \ r_S^* < 0.55 \ 95\% \) probability and \( 2\beta + \gamma = (88 \pm 29)\). In conclusion both the Babar and Belle collaboration have made enormous efforts to constraint the CKM angle \( \gamma \) and related quantities using many methods in different channels, leading to a precision in the determination that was not expected to be accessible at the B-factories experiments.

21. B. Aubert et al. [BABAR Collaboration], [arXiv:0805.2001 [hep-ex]].
R Value Measurements at 2.60, 3.07 and 3.65 GeV with BESII

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Using a data sample with a total integrated luminosity of 9998.3 nb$^{-1}$ collected at 2.6, 3.07 and 3.65 GeV with BESII, the cross section for $e^+e^-$ annihilating into hadronic final states ($R$ values) are measured. The statistical errors are smaller than 1%, and the systematic errors are about 3.5%. The running coupling constant of the strong interaction $\alpha_s^{(n)}(s)$ and $\alpha_s^{(5)}(M_Z^2)$ are determined from the measured $R$ values.

1 Instruction

The $R$ ratio is defined as the lowest level hadronic cross section normalized by theoretical $\mu^+\mu^-$ production cross section in $e^+e^-$ annihilation

$$R = \frac{\sigma_{\mu\mu}^0(e^+e^- \to \gamma^* \to \text{hadrons})}{\sigma_{\mu\mu}^0(e^+e^- \to \gamma^* \to \mu^+\mu^-)},$$

and is an important input parameter for precision tests of the Standard Model (SM). The errors on $R$ values measurements below 5 GeV have significant influence on the uncertainties of calculations of the QED running electromagnetic coupling constant $\alpha(s)$, muon anomalous magnetic moment $(\mu - 2)$ and global fits for the Higgs mass.$^{1,2,3}$ In addition, precision measurements of $R$ values between 2.0 - 3.7 GeV provide a test of perturbative QCD and QCD sum rule calculations.$^{4,5,6}$

$R$ value measurement is made at BESII$^7$ based on the expression$^8,9$

$$R_{exp} = \frac{N_{obs}^{had} - N_{bg}}{\sigma_{\mu\mu}^0 L_{exp} \epsilon_{\gamma}^0 \epsilon_{\mu\mu}^0 (1 + \delta_{obs})},$$

$^{a}$e-mail: haiming@ihep.ac.cn
where \(N_{\text{obs}}^{\text{had}}\) is number of the observed hadronic events, \(N_{\text{bg}}\) is number of the remnant QED backgrounds \((e^+e^-, \mu^+\mu^-, \tau^+\tau^-, \gamma\gamma, \text{etc.})\), \(L\) is the integrated luminosity, \(\varepsilon_{\text{trg}}\) is the trigger efficiency for hadronic events, \(\varepsilon_{\text{had}}^0\) is the hadronic efficiency without the simulation of the initial state radiation (ISR), and \((1+\delta_{\text{ISR}})\) is the effective factor of ISR in which the hadronic efficiencies for different bremsstrahlung energies are considered.

In 1998 and 1999, two series of \(R\) value measurements were made at 91 energy points between 2 - 5 GeV by the BESII experiments\(^8,9\). The average statistical errors are 2 - 4\%, and the systematical errors are 5 - 8\% depending on the energy points. In 2004, large-statistics data samples were accumulated at the center-of-mass energies of 2.6, 3.07 and 3.65 GeV; the total integrated luminosity was 9998.3 nb\(^{-1}\), and an additional 65.2 nb\(^{-1}\) data sample was accumulated at 2.2 GeV for the purpose of tuning the parameters of the hadronic event generator. Some improvements in the event selection, tuning of generator parameters and luminosity measurement are made in order to decrease the systematical errors. The previously used EGS-based detector simulation has been replaced by a GEANT3-based one. The consistency between data and Monte Carlo (MC) has been validated using many high purity physics channels\(^10\). With these improvements, the errors on the new measured \(R\) values are reduced to about 3.5\%.

2 Data analysis

The measurement scheme for this work is similar to that of used in previous one\(^9\). The strategy for selecting hadronic events is to subtract the QED backgrounds, cosmic ray and beam-associated backgrounds, the remaining events are then selected with specialized hadronic criteria. Two large sources of error in the measurement arise from the event selection and hadronic efficiency, and these have strong correlations between them.

2.1 Selection of hadronic events

In the BEPC energy region, the processes that originate in the beam-beam interaction region are \(e^+e^- \rightarrow e^+e^-, \mu^+\mu^-, \tau^+\tau^-, \gamma\gamma, e^+e^-X\) \((X\) means any possible final states\()\), hadrons and beam associated backgrounds. The observed final state charged particles are \(e, \mu, \pi, K\) and \(p\). Different types of backgrounds are identified using specialized criteria, and most of them can be rejected with good efficiency\(^8,9,11\).

Candidate hadronic events are classified by their number of charged tracks. The selection of hadronic events is done in two successive steps: one at the track level and the other at the event level\(^8,9\). In the BEPC energy region, the number of events with one observed/reconstructed charged track in BESII accounts for about 8 - 13\% of all hadronic events. In previous measurements, only hadronic events with two or more charged tracks \((n_{\text{ch}} \geq 2)\) were selected\(^8,9\). The omission of 0 and 1-track hadronic events introduces some uncertainty in the tuning of the hadronic event generator parameters; this in turn induces a sizable systematic error into the hadronic efficiency. However, for single-track events, contamination from beam-associated backgrounds is significant. Therefore, a more strict hadronic event selection is applied\(^11\). The event must have one charged track with good helicity fitting (the event can have any number of charged tracks with bad helicity fitting) and at least one reconstructed \(x^0\) are considered as single-track hadronic events.

Figure 1 shows the \(z\)-vertex distributions of the candidate hadronic events (including the residual beam-associated and QED backgrounds). The events produced by \(e^+e^-\) collisions originate near the collision point (in the neighborhood of \(z = 0\)), and the non-beam-beam backgrounds, such as those from beam-gas and beam-wall scattering, are distributed all along the beam-pipe. The number of observed hadronic events \(N_{\text{obs}}\) can be determined by fitting the distribution of event vertices along the beam direction with a Gaussian to describe the hadronic

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events and a \( n \)-th degree polynomial for the residual beam associated backgrounds.

The numbers of residual QED backgrounds, \( N_{bg} \), in Eq. (2), are deduced statistically from MC. The QED event generators with 1\% are used. The number of residual background events is determined as

\[
N_{bg} = L(\epsilon_{ee}\sigma_{ee} - \epsilon_{\mu\mu}\sigma_{\mu\mu} - \epsilon_{\tau\tau}\sigma_{\tau\tau} - \epsilon_{\gamma\gamma}\sigma_{\gamma\gamma}),
\]

where \( \sigma_{ee} \) is the production cross section for Bhabha events given by the corresponding generator, and \( \epsilon_{ee} \) is the residual ratio of Bhabha events that pass the hadronic event selection criteria. Other symbols have the similar meanings. The fraction of the remaining background \( e^+e^- \rightarrow e^+e^-X \) is much smaller than 1\% and is neglected. The values of \( \epsilon_{ee} \) and \( \epsilon_{\mu\mu} \) are about \( 5 \times 10^{-4} \), and \( \epsilon_{\tau\tau} \) is 36.45\% at 3.65 GeV. The errors on \( N_{bg} \) are given in Table 2.

2.2 Tuning the LUARLW parameters

The hadronic efficiency is determined using the LUARLW hadronic event generator. The physical basis of LUARLW is the solution of the Lund area law. The production of hadrons are described as the fragmentation of a semiclassical relativistic string, the quark components of the string and decays of unstable particles are handed by subroutines in JESTSET.

Both LUARLW and JETSET have some phenomenological parameters that have to be determined from the data. The basic method is to find a set of parameters that make various distributions (such as those sensitive to the efficiency) simulated by MC agree well with experimental data at all of the measured energies points. The distributions used in the data-MC comparison are: the multiplicities of charged and neutral tracks, the \( z \) vertices for charged tracks, the charged track momentum, the polar-angle between tracks and the beam direction, the deposited energy in the BSC, and fractions of \( \pi^\pm, K^\pm \), and some other short-lived particles (\( \pi^0, K_S, \phi, \Lambda \) etc.). With these distributions, the systematic errors corresponding to each criteria used in the hadronic event selection can be determined. Figure 2 shows some comparisons between data and the LUARLW MC at 3.07 GeV, where reasonable agreement is evident. More distributions at other energy points can be found in Refs. 11, 10.

2.3 Trigger

The trigger conditions are almost the same as those used for the \( R \) measurements reported in Refs. 8, 9. Since single-track hadronic events are also included in this measurement, the TOF back-to-back hit trigger requirement is not used, thereby making the trigger conditions somewhat lower than before. The trigger table used in data taking is given in Ref. 11. The trigger efficiencies
\( \epsilon_{\text{rej}} \) for hadronic events, listed in Table 1, are almost 99.8\%, and their associated errors are conservatively estimated to be 0.5\%.

2.4 Luminosity

The integrated luminosity is measured with wide angle Bhabha events. The measurement method is very similar to that described in Refs. \(^8,9\). In previous measurements, an EGS-based detector simulation was used. And for the measurement reported here, it has been replaced by a GEANT3-based package, which has better geometrical and material descriptions for the sub-detectors. In particular, the simulation of the BSC is significantly improved, and provides better consistency between MC and data. The Bhabha events are selected by using the BSC information. In order to decrease the uncertainty caused by events selection criteria, another Bhabha control sample only selected with MDC information is employed to correct the difference between data and MC. The efficiency correction factors of ranging from 0.994 to 1.026 are given in different energy points. In addition, the contribution from \( \ell^+\ell^- \rightarrow \gamma\gamma \) process is taken into account. As a result, the measurement precision of luminosity is significantly improved and their systematic uncertainties are smaller than 2\%.

2.5 Initial state radiative correction

An \( \mathcal{O}(\alpha^3) \) Feynman-diagram-based calculation for the initial state radiative (ISR) correction is used in both calculation of the ISR factor \((1 + \delta_{\text{obs}})\) and the simulation of radiative events by LUARLW. The detailed description on the ISR treatment can be found in Refs. \(^{20,21,22,23,24}\). In the ISR simulations and calculations, the contributions from both continuum and resonances are considered (quantities related to the narrow \( J/\psi \) and \( \psi' \) are treated analytically). For comparison, another approach based on structure functions \(^{25}\) is also used at all energy points. The differences between these two schemes for theoretical value of \((1 + \delta)\) are smaller than 1.1\%. The uncertainty in the effective ISR factor \((1 + \delta_{\text{obs}})\) due to errors of the hadronic cross sections at the different effective energies for radiative events are also considered (the errors on the hadronic cross section...
given in the PDG06 tables are used); these decrease with increasing of energy from 0.9% to 0.1%. The values of \( (1 + \delta_{\text{obs}}) \) and their errors are listed in Tables 1 and 2, respectively.

3 Error analysis

The Feynman-diagram-based ISR simulated angle and momentum distributions for the radiated photon is built into LUARLW, and the averaged hadronic efficiency \( \varepsilon_{\text{had}} \) with radiative effects can be obtained. The number of hadronic events \( N_{\text{had}}^{\text{obs}} \), the hadronic efficiency \( \varepsilon_{\text{had}} \) and their errors are correlated. The equivalent number of hadronic events, which corresponds to the number of hadronic events produced at the collision point, is defined as

\[
N_{\text{had}} = \frac{N_{\text{had}}^{\text{obs}}}{\varepsilon_{\text{had}}}.
\] (4)

The combined systematic error associated with the event selection and hadronic efficiency is denoted as \( \Delta N_{\text{had}} \). This error is caused by the discrepancy between data and MC samples for the selection hadronic criteria discussed in Section 2.1.

Except for the error mentioned above, an additional uncertainty of the parameters in the MC hadronization model is estimated to be about 1% by comparing different sets of tuned parameters, and is considered in error of hadronic event efficiency.

The error on \( N_{\text{had}}^{\text{obs}} \) due to the choice of degree for the polynomial used in the fitting is less than 0.7%. The fit errors for \( N_{\text{had}}^{\text{obs}} \) are 1.34% at 2.6 GeV, 1.11% at 3.07 GeV, and 0.73% at 3.65 GeV, which are calculated from the uncertainties in the fitted parameters of the Gaussian signal peaks. The total \( \Delta N_{\text{had}} \) is the quadratic sum of all fractional errors.

A conclusion of the KLN theorem is that the radiative corrections due to final state radiation (FSR) are negligible for a measurement of the inclusive hadronic cross sections that sums over all hadronic final states. At the present level of precision, the FSR correction factor in Eq. (2) can be neglected. However, the absence of final state radiation in the event generator introduces some error into the determination of the hadronic event detection efficiency. The masses of the hadrons produced in the final states are much greater than that of the initial state \( e^\pm \). As a result, the effect of FSR is much weaker than initial bremsstrahlung. Its influence is estimated to be 0.5% and is included in the error.

The 0-track events are not selected in this analysis, and the influence of 0-track events on the parameter tuning of LUARLW is not considered. This introduces some error into the hadronic event efficiency determination. Events with no charged tracks cannot be well separated from background. The fraction of 0-track events is estimated from the MC to be 3.4% at 2.6 GeV, 2.9% at 3.07 GeV, and 2.4% at 3.65 GeV. If the difference for 0-track events between MC and data is assumed to be 20%, the estimated errors for the lost/unobserved 0-track events are 0.7%, 0.9% and 0.5%, respectively. The error related to 0-track events is included into the error of \( N_{\text{had}} \) defined in Eq. (4).

In this analysis, hadronic events are classified according to their number of charged tracks. Therefore, errors in the tracking efficiency \( \sigma_{\text{trk}} \), the differences in the track reconstruction between data and MC, introduce some error into the classification and counting of the number of events. For an event with \( n_{\text{ch}} \) charged tracks, the probability that \( n_{\text{er}} \) of \( n_{\text{ch}} \) tracks are wrongly constructed roughly obeys a binomial distribution \( B(n_{\text{er}}; n_{\text{ch}}, \sigma_{\text{trk}}) \), where the parameter \( \sigma_{\text{trk}} \sim 2\% \) is the tracking error. Considering the distribution of charged multiplicity \( P(n_{\text{ch}}) \) for the inclusive hadronic sample (such as shown in Figure 2(a)), the effective error of tracking efficiency is

\[
\Delta \varepsilon_{\text{trk}} = \sum_{n_{\text{ch}} \geq 1} P(n_{\text{ch}}) B(n_{\text{er}}; n_{\text{ch}}, \sigma_{\text{trk}}).
\] (5)
The $R$ value measurement is, in fact, a counting of the number of hadronic events, so only those cases where all $n_{ch}$ tracks in an event are wrongly reconstructed ($n_{tr} = n_{ch}$) will cause an error in $\Delta N_{\text{had}}$. The values of $\Delta \epsilon_{trk}$ estimated from Eq. (5) are listed in Table 2. Since the fraction of 1-track events decreases with increasing center-of-mass energy, the error $\Delta \epsilon_{trk}$ also decreases with energy.

The numbers of errors on the $R$ value measurement are given in Table 2. As a cross check, the $R$ values are measured using the relation

$$R_{\text{exp}} = \frac{N_{\text{obs}}^\text{had} - N_{bg}}{\epsilon_\mu^0 \mu \epsilon_{\text{trg}}^\text{had}(1 + \delta)},$$

where $\epsilon_{\text{trg}}$ is the hadronic efficiency averaged over all of the ISR spectrum, and $(1 + \delta)$ is the corresponding theoretical ISR factor. The $R$ values determined with Eq. (6) are $2.17 \pm 0.01 \pm 0.07$ at 2.6 GeV, $2.13 \pm 0.01 \pm 0.07$ at 3.07 GeV, and $2.16 \pm 0.01 \pm 0.08$ at 3.65 GeV. The $R$ values measured with Eq. (2) and (6) are consistent to within 1%. Another cross check is also made by selecting hadronic events with $n_{ch} \geq 2$ as was done in Refs. 8, 9. In this case, the $R$ values measured at the three energy points are $2.20 \pm 0.02 \pm 0.08$, $2.13 \pm 0.01 \pm 0.07$ and $2.15 \pm 0.01 \pm 0.08$, respectively. The differences in $R$ values determined by selecting hadronic events with $n_{ch} \geq 1$ and $n_{ch} \geq 2$ are consistent within 1%.

Table 1: Items used in the determination of $R$ at each energy point.

<table>
<thead>
<tr>
<th>$E_{\text{cm}}$ (GeV)</th>
<th>$L$ (nb$^{-1}$)</th>
<th>$N_{\text{had}}$</th>
<th>$N_{bg}$</th>
<th>$\epsilon_{\text{trg}}$ (%)</th>
<th>$\epsilon_{\text{had}}^0$ (%)</th>
<th>$(1 + \delta_{\text{obs}})$</th>
<th>$R$</th>
<th>$\sigma_{\text{sys}}$</th>
<th>$\sigma_{\text{sys}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>2.60</td>
<td>1222.06</td>
<td>24026</td>
<td>193</td>
<td>99.85</td>
<td>63.81</td>
<td>1.08</td>
<td>2.18</td>
<td>0.02</td>
<td>0.08</td>
</tr>
<tr>
<td>3.07</td>
<td>2290.72</td>
<td>33333</td>
<td>208</td>
<td>99.79</td>
<td>67.63</td>
<td>1.11</td>
<td>2.13</td>
<td>0.02</td>
<td>0.07</td>
</tr>
<tr>
<td>3.65</td>
<td>6485.30</td>
<td>83707</td>
<td>4937</td>
<td>99.89</td>
<td>71.83</td>
<td>1.21</td>
<td>2.14</td>
<td>0.01</td>
<td>0.07</td>
</tr>
</tbody>
</table>

Table 2: Summary of the systematic errors in percent (%).

| $E_{\text{cm}}$ (GeV) | $L$ (nb$^{-1}$) | $N_{\text{had}}$ | $N_{bg}$ | $\Delta \epsilon_{trk}$ | $\epsilon_{trg}$ | $(1 + \delta_{\text{obs}})$ | Total |
|-----------------------|-----------------|-----------------|--------|-----------------|-----------------|-----------------|-----------------|--------|
| 2.60                  | 2.00            | 2.79            | 0.05   | 0.32            | 0.50            | 1.18            | 3.68            |
| 3.07                  | 1.96            | 2.53            | 0.05   | 0.29            | 0.50            | 1.15            | 3.45            |
| 3.65                  | 1.38            | 2.74            | 0.35   | 0.26            | 0.50            | 1.10            | 3.33            |

4 Results and discussions

Tables 1 and 2 list the quantities used in the measurement of $R$ with Eq. (2) and the contributions to the total error. The results are displayed in Fig. 3, together with previous measurements. The errors on the $R$ values reported here are about 3.5%. The measured $R$ values are consistent within errors with the prediction of perturbative QCD 4.

Compared with our previous results 8, 9, the measurement precision has been improved due to four main refinements to the analysis: (1) the simulation of BES with a GEANT3 based package that has a more detailed geometrical description and matter definition for the sub-detectors; (2) large data samples are taken at each energy point, with statistical errors that are smaller than 1%; (3) the selected hadronic event sample is expanded to include single-track events, which supplies more information to the tuning of LARLW with resulting improved parameters; (4) the distributions used for parameter tuning are those related to the hadronic selection criteria; the better agreement between MC and data reduces the error on the hadronic event efficiency.

Based upon the measured $R$ values and the perturbative QCD expansion that computes $R_{QCD}(\alpha_s)$ to an $O(\alpha_s^4)$ approximation 26, the strong running coupling constant $\alpha_s^{(3)}(s)$ can be
determined at each energy point\textsuperscript{28,29} by solving the equation \( R_{\text{exp}} \pm \sigma_{\text{sfs}} \pm \sigma_{\text{sus}} = R_{\text{QCD}}(\alpha_s^{(3)}) \). The obtained of \( \alpha_s^{(3)}(s) \) are evolved to 5 GeV, and the weighted average of \( \tilde{\alpha}_s^{(4)}(25\text{GeV}^2) \) is listed in Table 3. When evaluated at the \( M_Z \) scale, the resulting value is \( \alpha_s^{(5)}(M_Z^2) = 0.117_{-0.017}^{+0.012} \), which agrees with the world average value within the quoted errors\textsuperscript{4}.

![Graph showing R values reported here together with other measurements below 5 GeV.]

**Figure 3:** \( R \) values reported here together with other measurements below 5 GeV.

**Table 3:** \( \alpha_s(s) \) determined from \( R \) values measured at 2.600, 3.070 and 3.650 GeV, and evolved to 5 GeV and \( M_Z \). The first and second errors are statistical and systematic, respectively.

<table>
<thead>
<tr>
<th>( \sqrt{s}(\text{GeV}) )</th>
<th>( \alpha_s^{(3)}(s) )</th>
<th>( \alpha_s^{(4)}(25\text{GeV}^2) )</th>
<th>( \alpha_s^{(4)}(25\text{GeV}^2) )</th>
<th>( \alpha_s^{(5)}(M_Z^2) )</th>
</tr>
</thead>
<tbody>
<tr>
<td>2.60 ( \pm 0.03 )</td>
<td>0.266( \pm 0.030 \pm 0.125 )</td>
<td>0.212( \pm 0.019 \pm 0.086 )</td>
<td>0.200( \pm 0.044 )</td>
<td>0.117( \pm 0.017 )</td>
</tr>
<tr>
<td>3.07 ( \pm 0.029 \pm 0.103 )</td>
<td>0.192( \pm 0.097 \pm 0.104 )</td>
<td>0.160( \pm 0.023 \pm 0.086 )</td>
<td>0.200( \pm 0.050 )</td>
<td>0.117( \pm 0.012 )</td>
</tr>
<tr>
<td>3.65 ( \pm 0.015 \pm 0.104 )</td>
<td>0.207( \pm 0.015 \pm 0.104 )</td>
<td>0.189( \pm 0.013 \pm 0.069 )</td>
<td>0.200( \pm 0.050 )</td>
<td>0.117( \pm 0.017 )</td>
</tr>
</tbody>
</table>

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Light Hadrons and New Enhancements in $J/\psi$ Decays at BESII

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Based on 58 million $J/\psi$ samples collected by the BESII detector at the BEPC, many mesons, baryons, and new resonances have been reported. Here, I will review some recent results of glueball candidates and new enhancement.

1 Introduction

In this paper, some recent BESII results are reported based on 58 million $J/\psi$ events collected by the BESII detector at the BEPC. For much more detail, please see the references.

2 Scalars ($0^{++}$)

As we know that so many scalars are listed in PDG06 \cite{PDG06}, but according to the quark model no enough room for all of these scalar particles. On the other hand, the Lattice QCD predicted that the ground state glueball is $0^{++}$, and its mass is around 1.5~1.8 GeV. Theoretical physicists expect that glueballs will mix with nearby $q\bar{q}$ states of the same quantum numbers \cite{Chen:2004mr, Cheung:2005ux}, which makes the situation more difficult for the glueball identification. Although the identification of a glueball is very complicated, there are several glueball candidates, such as $f_{0}(1500)$ and $f_{0}(1710)$, considering the possible mix with the ordinary $q\bar{q}$ meson, $f_{0}(1370)$, $f_{0}(1500)$, $f_{0}(1710)$, and $f_{0}(1790)$ have been analyzed for more detail by using the partial wave analyzes (PWA) method in $J/\psi \rightarrow \gamma\pi\pi$, $\gamma K\bar{K}$, $J/\psi \rightarrow \omega K\bar{K}$, and $J/\psi \rightarrow \phi\pi\pi$, $\phi K\bar{K}$ channels.

2.1 The Analysis of $J/\psi \rightarrow \gamma\pi\pi$ and $\gamma K\bar{K}$ Channels

The partial wave analyzes of $J/\psi \rightarrow \gamma\pi^{+}\pi^{-}$ and $J/\psi \rightarrow \gamma\pi^{0}\pi^{0}$ show the evidence for two $0^{++}$ states around the 1.45 and 1.75 GeV/c$^2$ mass regions (Fig. 1, 2) \cite{Chen:2006iy}. The $f_{0}(1500)$ has
a mass of $1466 \pm 6 \pm 20 \text{ MeV}/c^2$, a width of $108_{-11}^{+14} \pm 25 \text{ MeV}/c^2$, and a branching fraction $B(J/\psi \rightarrow \gamma f_0(1500) \rightarrow \gamma \pi^+ \pi^-) = (0.67 \pm 0.02 \pm 0.30) \times 10^{-4}$. The $0^{++}$ state in the $\sim 1.75$ GeV/$c^2$ mass region has a mass of $1765_{-3}^{+4} \pm 13 \text{ MeV}/c^2$ and a width of $145 \pm 8 \pm 69 \text{ MeV}/c^2$.

The PWA of $J/\psi \rightarrow \gamma K^+ K^-$ and $J/\psi \rightarrow \gamma K^0_S K^0_S$ show strong production of the $f_2(1525)$ and the S-wave resonance $f_0(1710)$ (Fig. 3)\(^5\). The $f_0(1710)$ peaks at a mass of $1740 \pm 4^{+10}_{-25} \text{ MeV}$ with a width of $166^{+5}_{-8}^{+15}_{-10} \text{ MeV}$.

The Analysis of $J/\psi \rightarrow \omega K^+ K^-$ Channel

From Fig. 4, one can see that a dominant feature is $f_0(1710)$\(^6\). The fitted $f_0(1710)$ optimizes at $M = 1738 \pm 30 \text{ MeV}/c^2$, $\Gamma = 125 \pm 20 \text{ MeV}/c^2$.

The Analysis of $J/\psi \rightarrow \phi \pi^+ \pi^-$ and $\phi K^+ K^-$ Channels

After the partial wave analyzes for these $\phi \pi \pi$ and $\phi K K$ channels\(^7\), the data reported here have three important features. Firstly, the parameters of $f_0(980)$ are all well determined. Secondly,
there is the clearest signal to date of $f_0(1370) \to \pi^+\pi^-$; a resonant phase variation is required, from interference with $f_2(1270)$. Thirdly, there is a clear peak in $\pi\pi$ at 1775 MeV/c$^2$, consistent with $f_0(1790)$; spin 2 is less likely than spin 0.

In summary, (1) $f_0(1370)$ has been seen in $J/\psi \to \phi\pi\pi$, but not in $J/\psi \to \omega\pi\pi$. (2) No peak of the $f_0(1500)$ directly seen in $J/\psi \to \phi K K, \omega K K, \phi\pi\pi$, and $\omega\pi\pi$, but in proton-proton scattering is quite clear. (3) $f_0(1710)$ is observed clearly in both $J/\psi \to \phi K K$ and $J/\psi \to \omega K K$, but with $Br(J/\psi \to \omega f_0(1710) \to \omega K K)/Br(J/\psi \to \phi f_0(1710) \to \phi K K) \sim 6$, which is against a simple $s\bar{s}$ configuration for this state. (4) $f_0(1790)$ which is seen in $\pi\pi$ rather than $KK$.

Different models have different interpretations for these experimental results. One of the interpretations is from Cheng $^8$, he explained that (1) $f_0(1710)$ is composed primarily of the scalar glueball. (2) $f_0(1500)$ is close to an $SU(3)$ octet. The glueball content of $f_0(1500)$ is very tiny because an $SU(3)$ octet does not mix with the scalar glueball. (3) $f_0(1370)$ consists of an approximate $SU(3)$ singlet with some glueball component ($\sim 10\%$).
3 Pseudo-scalars (0−+)

The first observation of η(1440) was made in p ¯p annihilation at rest into η(1440)π+π−, η(1440) → K ¯Kπ. Nowadays, the existence of two overlapping pseudo-scalar states has been suggested to instead of the η(1440): one around 1405 MeV/c² decays mainly through a0(980)π (or direct K ¯Kπ), and the other around 1475 MeV/c² mainly to K*(892) ¯K. It is therefore conceivable that the higher mass state is the s ¯s member of the 21S0 nonet, while the lower mass state may contain a large gluonic content.

Figure 6: The γρ invariant mass distribution. The insert shows the full mass scale where the η(958) is clearly observed.

Figure 7: The invariant mass of γφ after sideband background subtraction.

In our J/ψ → γγV analysis, there is a resonance around 1424 MeV at the J/ψ → γγρ channel. Comparing our result on the branching ratio B(J/ψ → γX(1424) → γγρ) = (1.07 ± 0.17 ± 0.11) × 10−4, and the upper limit of B(J/ψ → γX(1424) → γγφ) < 0.82 × 10−4 (95% C.L.), we cannot draw a definite conclusion on whether the X(1424) is either a q ¯q state or a glueball state.

We also analyzed the η(1405)/η(1475) at J/ψ → {ω, φ}K ¯Kπ channels. In the invariant mass spectra of K0S ¯Kπ and K+K−π0 recoiling against the ω signal region, the resonance at 1.44 GeV/c² is observed, while in the invariant mass spectra of K0S ¯Kπ and K+K−π0 recoiling against the φ signal region, no significant structure near 1.44 GeV/c² is seen and an upper limit on the J/ψ decay branching fractions at the 90% C.L. are given in Table 1.

Table 1: The mass, width, and branching fractions of J/ψ decays into {ω, φ}X(1440).

<table>
<thead>
<tr>
<th>J/ψ → ωX(1440)</th>
<th>J/ψ → ωX(1440)</th>
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<tbody>
<tr>
<td>(X → K0S ¯K+π− + c.c.)</td>
<td>(X → K+K−π0)</td>
</tr>
<tr>
<td>M = 1437.6 ± 3.2 MeV/c²</td>
<td>M = 1445.9 ± 5.7 MeV/c²</td>
</tr>
<tr>
<td>Γ = 48.9 ± 9.0 MeV/c²</td>
<td>Γ = 34.2 ± 18.5 MeV/c²</td>
</tr>
<tr>
<td>B(J/ψ → ωX(1440) → ωK0S ¯K+π− + c.c.) = (4.86 ± 0.69 ± 0.81) × 10−4</td>
<td>B(J/ψ → ωX(1440) → ωK+K−π0) = (1.92 ± 0.57 ± 0.38) × 10−4</td>
</tr>
<tr>
<td>B(J/ψ → φX(1440) → φK0S ¯K+π− + c.c.) &lt; 1.93 × 10−5 (90% C.L.)</td>
<td>B(J/ψ → φX(1440) → φK+K−π0) &lt; 1.71 × 10−5 (90% C.L.)</td>
</tr>
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</table>
4 New Enhancements

A narrow enhancement is observed in $J/\psi \rightarrow \gamma p\bar{p}$\textsuperscript{17}. Assuming that the $p\bar{p}$ system is in an S-wave resulted in a resonance with mass $M = 1859^{+3+25}_{-10-25}$ MeV/c\(^2\), width $\Gamma \sim 30$ MeV/c\(^2\) (at the 90% C.L.) and product branching fraction $B(J/\psi \rightarrow \gamma X) \cdot B(X \rightarrow p\bar{p}) = (7.0 \pm 0.4(stat) \pm 1.9(syst)) \times 10^{-5}$. The data not precise enough to determine the angular distribution. According to the theoretical calculation\textsuperscript{18}, if the $X$ is a bound state of $(p\bar{p})$, the decay channel $(X \rightarrow \eta 4\pi)$ is favored over $(X \rightarrow \eta 2\pi, 3\eta)$.

The decay channel $J/\psi \rightarrow \gamma \pi^+\pi^-\eta'$ is analyzed using two $\eta'$ decay modes, $\eta' \rightarrow \pi^+\pi^-\eta$ and $\eta' \rightarrow \rho$\textsuperscript{19}. A resonance, the $X(1835)$, is observed with a high statistical significance of $7.7\sigma$ in the $\pi^+\pi^-\eta'$ invariant mass spectrum. From a fit with a Breit-Wigner function, the mass is determined to be $M = 1833.7 \pm 6.1(stat) \pm 2.7(syst)$ MeV/c\(^2\), the width is $\Gamma = 67.7 \pm 20.3(stat) \pm 7.7(syst)$ MeV/c\(^2\), and the product branching fraction is $B(J/\psi \rightarrow \gamma X) \cdot B(X \rightarrow \pi^+\pi^-\eta') = (2.2 \pm 0.4(stat) \pm 0.4(syst)) \times 10^{-4}$. The mass and width of the $X(1835)$ are not compatible with any known meson resonance\textsuperscript{1}. If we redoing the S-wave BW fit to the $p\bar{p}$ invariant mass spectrum\textsuperscript{17} including the zero Isospin, S-wave final-state-interactions (FSI) factor\textsuperscript{20}, yields a mass $M = 1831 \pm 7$ MeV/c\(^2\) and a width $\Gamma < 153$ MeV/c\(^2\) (at the 90% C.L.), these values are in good agreement with the mass and width of $X(1835)$ reported here.

In the analysis of $J/\psi \rightarrow \omega p\bar{p}$\textsuperscript{21}, no significant enhancement near the $p\bar{p}$ mass threshold is observed, and an upper limit of $B(J/\psi \rightarrow \omega X)B(X \rightarrow p\bar{p}) < 1.5 \times 10^{-5}$ is determined at the 95% confidence level.

5 Summary

Using the 58 M $J/\psi$ events sample taken with the BESII detector at the BEPC storage ring, BES experiment provided many interesting results, especially for the study of the lowest glueball candidates, the structure of $\eta(1440)$, and the new enhancement of $X(1835)$, but since the limit of the statistics, the better results (with higher statistics and better accuracy) will be needed for well understanding. The upgraded BEPCII/BESIII will provide a huge $J/\psi$ decay samples for the further analysis.

References

NA48 Results

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Measured decay rates of $K^\pm \rightarrow e^\pm \pi^0 \nu_e$ and $K^\pm \rightarrow \mu^\pm \pi^0 \nu_\mu$, normalized to $K^\pm \rightarrow \pi^\pm \pi^0$ are presented. These measurements are based on $K^\pm$ decays collected in a dedicated run in 2003 by the NA48/2 experiment at CERN. Using the PDG 2006 average for the $K^\pm \rightarrow \pi^\pm \pi^0$ normalization mode, the results are found to be larger than the current values given by the PDG 2006 and lead to a larger magnitude of the $|V_{us}|$ CKM element than previously accepted. When combined with the latest PDG 2006 value of $|V_{ud}|$, the result is in agreement with unitarity of the CKM matrix.

The ratio $R_K = \Gamma(K^\pm \rightarrow e^\pm \nu) / \Gamma(K^\pm \rightarrow \mu^\pm \nu)$ is calculated with very high precision within the Standard Model (SM), but corrections due to the presence of New Physics could be as high as 3%. The data obtained by the NA48/2 experiment in two years of data taking at the CERN SPS accelerator has been analyzed. The obtained result for $R_K$ is two times more precise than the world average but is still insufficient to probe the existence of physics beyond the Standard Model. The status of the analysis of the data taken in 2007, aimed for a sub-percent precision of $R_K$, will be summarized.

1 Introduction

The NA48 experiment at CERN SPS is a fixed target experiment devoted to kaon physics operating since 1997. Until 2001 the experiment studied the neutral kaon decays and provided the final measurement of $\epsilon'/\epsilon$ 1. A charged kaon physics program (NA48/2) took place in 2003 and 2004: it was mainly devoted to the search for direct CP violation in the $K^\pm$ decays into three pions\(^2\). Beside this main topic, also semileptonic and rare charged kaon decays were studied. To this end dedicated runs with reduced intensity were taken in 2003 and 2004. The present work describes the final result of the measurement of the branching ratio of $K^\pm \rightarrow l^\pm \pi^0 \nu_l$ ($l = e, \mu$)\(^3\) using the 2003 data and the preliminary results of the measurement of the ratio $R_K = \Gamma(K^\pm \rightarrow e^\pm \nu_e) / \Gamma(K^\pm \rightarrow \mu^\pm \nu_\mu)$ based on the 2003 and 2004 data. The NA62 collaboration is currently carrying on the kaon physics program at CERN SPS. The first phase of this experiment aims
for a sub-percent precision measurement of $R_K$, for which data were taken in 2007 with the NA48/2 apparatus.

2 NA48/2 Experimental Setup

The experiment used simultaneous $K^\pm$ beams produced by 400 GeV/c protons delivered by the SPS and impinging on a Be target with a duty cycle of 4.8 s spill over a 16.8 s accelerator period. The proton intensity on target was about $7 \times 10^{11}$ proton per spill during the 2003 and 2004 normal runs. It was reduced during the special runs to allow data taking with a minimum bias trigger, while it was increased up to more than $10^{12}$ protons per spill during the 2007 run. A 100 m long beam line selected charged beams with $60 \pm 3$ GeV/c average momentum in 2003 and 2004 and $75 \pm 2$ GeV/c in 2007. The detector sits about 100 meter downstream to the end of the beam line and detected the products of the kaon decays happening in the evacuated region between the end of the beam line and the beginning of the detector. A detailed description of the NA48 apparatus can be found elsewhere\(^4\). The most relevant devices for the measurements described here were: the magnetic spectrometer, consisting of 4 drift chambers and one magnet and the high resolution liquid krypton electromagnetic calorimeter. The spectrometer worked with a reduced magnetic field in 2003 and 2004 and with full magnetic field in 2007 allowing a better momentum resolution. Other devices were the hodoscope for charged particle triggering and precise time measurement and a muon detector.

3 Measurement of the $K_{l3}$ Branching Ratio

3.1 Theoretical aspects

The following master formula describes the branching ratio of the semileptonic charged kaon decays\(^5\):

$$BR(K_{l3}) = \tau_K G_F^2 \frac{m_K^5}{384\pi^3} S_{EW} |V_{us}|^2 |f_+(0)|^2 I_K^l(1 + \delta_{SU2}^K + \delta_{EM}^K)^2.$$  \hspace{1cm} (1)

Here $K_{l3}$ is a short-hand notation for $K^\pm \to l^\pm \pi^0 \nu_l$ with $l$ equal to $e$ or $\mu$. $\tau_K$ is the average life time of $K^\pm$, $G_F$ the Fermi constant and $m_K$ the mass of the charged kaon. $S_{EW}$ is the short distance radiative correction, $\delta_{SU2}^K$ and $\delta_{EM}^K$ are the model dependent long distance corrections due to isospin breaking in strong and electromagnetic interactions. Two form factors, $f_+(t)$ and $f_0(t)$, describe the dynamic of the semileptonic decays. Their $t$ dependence can be approximated as:

$$f_+(t) = f_+(0) \left( 1 + \lambda_+ \frac{t}{m_+^2} + \lambda''_+ \frac{t^2}{m_+^4} \right), \quad f_0(t) = f_+(0) \left( 1 + \lambda_0 \frac{t}{m_0^2} \right).$$ \hspace{1cm} (2)

$f_+(0)$ is the form factor at zero momentum transfer. The parameters $\lambda_+$, $\lambda''_+$ and $\lambda_0$ are measured\(^6\). $I_K^l$ is the result over the phase space integration after factorizing out the $f_+(0)$ and depends on $\lambda_+$, $\lambda''_+$ and $\lambda_0$, using the above approximation \(^5\). Finally $V_{us}$ is the element of the CKM matrix which describes the $u$-$s$ transitions.

It turns out that the measurement of the branching ratio of the charged $K_{l3}$ decays allows a clean test of the $u$-$s$ quark transitions. Moreover the ratio between the branching ratios of the $K_{e3}$ and $K_{\mu3}$ provides also an experimental test of the $\mu$--$e$ universality.

3.2 Data taking and Analysis Strategy

Because of the impossibility to measure precisely the absolute kaon flux, NA48 measured the semileptonic branching ratios normalized to $K^\pm \to \pi^\pm \pi^0$, that is the ratios $R_{K_{l3}/K_{2\pi}} \equiv$
\(\Gamma(K_{e3})/\Gamma(K^{\pm} \to \pi^{\pm}\pi^{0})\). It is relevant that the single track topology for both the signal and the normalization channel allows a first order cancellation of the systematics.

Hits in the hodoscope compatible with a one track decay were the only input of the trigger. The trigger efficiency was measured on data to be greater than 99.8\%. An offline one track selection using the spectrometer informations and a \(\pi^{0}\) identification based on the calorimeter data, defined a sample of \(K_{e3}, K_{\mu3}\) and \(K^{\pm} \to \pi^{\pm}\pi^{0}\) decays. Extra activity in the calorimeter was allowed to select inclusively also the corresponding radiative decays. Kinematical cuts exploiting the missing energy and the decay topology separated the semileptonic from the two pions decays. The particle identification was used to distinguish the electron from the muon channel. In particular the requirement \(E_{LKR}/P > 0.95\) identified an electron, where \(E_{LKR}\) is the energy released by the particle in the calorimeter and \(P\) is the particle momentum measured by the spectrometer; the cut \(E_{LKR}/P < 0.8\) defined a pion. Finally, the presence of a hit in the muon detector, matching in space and time with the track, tagged a muon. The total number of selected events per decay mode was: \(87 \times 10^{3}\ K_{e3}, 77 \times 10^{3}\ K_{\mu3}\) and \(729 \times 10^{3}\ K^{\pm} \to \pi^{\pm}\pi^{0}\).

The acceptance was computed using a GEANT\(^7\) based Monte Carlo simulation. The event generation made use of the previously described parametrization for the form factors, with \(\lambda_{e}^{0}, \lambda_{\mu}^{0}\) and \(\lambda_{0}\) taken from reference \(^6\). The phase space was corrected according to the Ginsberg prescription \(^8\) to account for radiative corrections. The PHOTOS package \(^9\) provided the generation of real bremsstrahlung photons. The acceptance varied between 7\% and 14\% depending on the decay mode. Different expressions of the form factors were also considered \(^10\) and the corresponding variation of the final result quoted as systematic uncertainty. The particle identification was a source of inefficiency not canceled in the single ratio. It was measured on data and varied between 98.5\% and 99.5\%, depending on the particle type. The corresponding error was quoted as systematic uncertainty. The Monte Carlo simulation pointed out a background contamination below 0.1\% for \(K_{e2}\) and at the level of 0.2\% and 0.3\% for \(K_{\mu2}\) and \(K^{\pm} \to \pi^{\pm}\pi^{0}\), respectively.

### 3.3 Results

The results are:

\[
\begin{align*}
\mathcal{R}_{K_{e3}/K^{2}\pi} &= 0.2470 \pm 0.0009_{\text{stat}} \pm 0.0004_{\text{syst}} \\
\mathcal{R}_{K_{\mu3}/K^{2}\pi} &= 0.1636 \pm 0.0006_{\text{stat}} \pm 0.0003_{\text{syst}} \\
\mathcal{R}_{K_{\mu3}/K_{e3}} &= 0.663 \pm 0.003_{\text{stat}} \pm 0.001_{\text{syst}}
\end{align*}
\]

(3)

Analysis of these results as a function of their basic distributions shown stability.

Taking the branching ratio of \(K^{\pm} \to \pi^{\pm}\pi^{0}\) from \(^6\) the branching ratio for the semileptonic decays are:

\[
\begin{align*}
BR(K_{e3}) &= 0.05168 \pm 0.00019_{\text{stat}} \pm 0.00008_{\text{syst}} \pm 0.00030_{\text{norm}} \\
BR(K_{\mu3}) &= 0.03425 \pm 0.00013_{\text{stat}} \pm 0.00006_{\text{syst}} \pm 0.00020_{\text{norm}}
\end{align*}
\]

(4)

The uncertainty is dominated by the error on the measurement of the branching ratio of the \(K^{\pm} \to \pi^{\pm}\pi^{0}\). Both the values are significantly above the PDG 2006 values. The \(BR(K_{e3})\) agrees with the BNL E865 \(^{11}\) and the ISTRA+ ’07 \(^{12}\) measurements. Both the NA48 measurements, however, do not agree with the values measured by KLOE \(^{13}\) which are in agreement with \(^6\). The recent KLOE measurement of the \(BR(K^{\pm} \to \pi^{\pm}\pi^{0})\) \(^{14}\), significantly lower than the PDG 2006 one, partially recover the difference between NA48 and KLOE.

The measurements 4 allow the extraction of \(V_{us}\). To this end the following values were used: \(S_{ew} = 1.023^{15}\), \(I_{K} = 0.1591\) and \(I_{K}^{c} = 0.1066\) (\(\lambda_{e}^{0}, \lambda_{\mu}^{0}\) and \(\lambda_{0}\) from \(^6\)), \(\delta_{SU2} = 2.31\%\), \(\delta_{em} = 0.03\%\) and \(\delta_{K_{em}} = 0.2\%\) from \(^{16,17}\), \(G_{F} = 1.16637 \times 10^{-5}\) GeV\(^{-2}\) \(^{18}\) and \(m_{K}\) and \(\tau_{K}\) from \(^6\). The result is

\[
|V_{us}|f_{+}(0) = 0.2188 \pm 0.0012
\]

(5)
combined for $K_{e3}$ and $K_{\mu3}$. The values obtained for the two decay modes separately are in agreement among themselves. The result is in agreement with the expected value computed using $V_{ud} = 0.9738 \pm 0.0003^{15}$, $|V_{ub}| = (3.6 \pm 0.7) \times 10^{-3}^{6}$, $f_+(0) = 0.961 \pm 0.008^{5}$ and assuming unitarity, as shown in figure 1. The results are compatible with the unitarity of the CKM matrix. Finally the measured value of $R_{K_{\mu3}/K_{e3}}$ implies the $\mu$-e universality violating quantity $g_{\mu}f_{+}^{\mu}(0)/g_{e}f_{+}^{e}(0) = 0.99 \pm 0.01$, consistent with one within the experimental errors.

4 Measurement of $R_{K}$

4.1 Theoretical aspects and experimental status

The measurement of $R_{K} \equiv R(K_{e2})/R(K_{\mu2})$ provides an accurate test of the lepton universality predicted in the SM. Here $K_{I2}$ is a short-hand notation for $K^{\pm} \rightarrow l^{\pm}\nu_{l}$. Thanks to the cancellation in the ratio of the hadronic uncertainties, the SM predicts $R_{K}$ with a sub-permille accuracy $^{20}$:

$$R_{K} = \frac{m_{K}^{2}}{m_{\mu}^{2}} \left( \frac{m_{K}^{2} - m_{e}^{2}}{m_{K}^{2} - m_{\mu}^{2}} \right)^{2} (1 + \delta R_{QED}) = (2.477 \pm 0.001) \times 10^{-5}. \quad (6)$$

Here $m_{K,e,\mu}$ are the masses of the kaon, electron and muon and $\delta R_{QED}$ is the correction for virtual photon processes and inner bremsstrahlung photon emission.

The helicity suppression makes $R_{K}$ sensitive to new physics. A theoretical study $^{21}$ suggests the possibility of up to some percent deviation from the SM value induced by lepton flavor violating effects, as those arising in supersymmetry extensions of SM. As a consequence a sub-percent precision measurement of $R_{K}$ could probe physics beyond SM.

The PDG 2006 value, $R_{K} = (2.45 \pm 0.11) \times 10^{-5}$, is far from the accuracy needed. NA48 provided preliminary measurements at 2% precision using 2003 and 2004 data. More recently KLOE $^{22}$ measured this quantity with 2% level accuracy. NA62 took data for 4 months in 2007 and collected more than $10^{5}$ $K_{e2}$ aiming for a 0.5% precision.

4.2 Analysis Strategy

The signal signature is one track in the final state compatible with a two body kinematics. Both kinematics and particle identification discriminate between the electron and the muon channel. The requirement $E_{L,K_{\ell}}/P > 0.95$ identifies an electron, like in the $K_{\ell3}$ analysis previously described. Once data are collected using similar triggers for the two channels, systematics
4.3 Preliminary results from 2003-2004 run

The number of $K_{e2}$ collected by NA48 in 2003 and 2004 after background subtraction was $(4670 \pm 77_{\text{stat}} \pm 20_{\text{syst}})$ and $(3407 \pm 63_{\text{stat}} \pm 54_{\text{syst}})$, respectively. The systematic uncertainty refers to the background subtraction procedure. In particular the muon background in the $K_{e2}$ sample was estimated at the level of 14%, using a pure $K_{e2}$ sample at low momentum. The results are \(^{23,24}\)

$$ R_K = (2.416 \pm 0.043 \pm 0.024) \times 10^{-5} \quad (2003) $$

$$ R_K = (2.455 \pm 0.045 \pm 0.041) \times 10^{-5} \quad (2004) $$

The 2003 data suffered from kinematical requests at trigger level which induced a large trigger efficiency correction. The choice of a minimum bias trigger for $K_{\mu2}$ and and the minimum bias plus a further requirement on the total energy in the electromagnetic calorimeter for $K_{e2}$, avoided that problem in 2004. The systematics of both the measurements are largely dominated by the uncertainty in the background subtraction. The other systematics are below 0.2%.

4.4 NA62 run: data collected and status of the analysis

NA62 took data in 2007. In comparison to the 2003-2004 run, the increase of the average beam momentum from 60 to 75 GeV/c and the shrink of the momentum bite from 3 to 2 GeV/c allowed a better background rejection. For the same purpose the spectrometer worked with a stronger magnetic field. The trigger was the same as in 2004. During the run an important accidental background appeared in the $K^-$ data. For that reason only $K^+$ were taken for most of the period. The statistics collected matched the goal of the run: the total number of $K_{e2}$ selected on-line was, in fact, $1.1 \times 10^5$. Figure 2 a) shows the squared invariant missing mass distribution, $m_{\text{miss}}^2$, for selected $K_{e2}$-like events, where $m_{\text{miss}}^2$ is defined as the square of the difference between the kaon and the measured track four momenta. The number of good $K_{e2}$ refers to the events under the peak.

Part of the data were taken with a lead bar 18.0 cm wide and 9 $X_0$ thick in front of the liquid kripton calorimeter to measure the probability of muon catastrophic energy loss. The presence of the bar induced about 18% loss in $K_{e2}$ acceptance. The lead acted as a muon filter selecting a pure sample of muons without electron contamination. More precisely this bar was placed just in front of six scintillator counters of the hodoscope used to disentangle muons not interacting in lead. The normal data taking provided more than 2000 $\mu$ with momentum greater than 35 GeV/c faking an electron. Other 2000 $\mu$ of that type came from special muon runs. The preliminary result of the muon catastrophic energy loss probability as a function of momentum measured using un-calibrated data from the special runs only is shown figure 2...
Figure 2: (a) $m_{miss}^2$ in GeV$^2/c^2$ for $K_{e2}$ events collected during the 2007 run. The prediction for the $K_{e2}$ and the $K_{e2\gamma}$ structure dependent contamination are also shown. (b) Probability that a muon releases in the liquid krypton calorimeter more than 95% of its energy as a function of muon momentum in GeV/c.

(b). It corresponds to a $K_{\mu2}$ contamination in the $K_{e2}$ sample of $7.5 \pm 0.1\%$. The background level, therefore, can be controlled with the requested accuracy. Special runs with the kaon beam dumped and with $K^-$ only were also taken to study the residual accidental background in $K^+$ data. Finally a measurement of the electron identification efficiency on the overall $K_{e2}$ momentum spectrum required also special runs with $K_L$ beam, which allow the selection of a pure sample of electron through $K_L \rightarrow e^+\pi^-\nu$ decays.

The analysis of the 2007 data is already started and preliminary results are expected soon.

References

7. GEANT Detector Description and Simulation Tool, CERN Program Library Long Write-up W5013 1994.
\[ \Sigma^+ \rightarrow p\mu^+\mu^- : \text{Standard Model or New Particle?} \]

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The HyperCP collaboration observed three events for the decay \( \Sigma^+ \rightarrow p\mu^+\mu^- \). They suggested that new physics may be required to understand the implied decay rate and the observed \( m_{\mu\mu} \) distribution. Motivated by this result, we re-examine this mode. First within the standard model, and then assuming there is a new particle. Within the SM we find that \( \Sigma^+ \rightarrow p\mu^+\mu^- \) is long-distance dominated and its rate falls within the range suggested by the HyperCP measurement. We then examine the conditions under which the observation is consistent with a light Higgs boson and find an explicit example that satisfies all the constraints: the light pseudoscalar Higgs boson in the next-to-minimal supersymmetric standard model (NMSSM).

1 Introduction

The HyperCP collaboration has observed three events for the mode \( \Sigma^+ \rightarrow p\mu^+\mu^- \). A striking feature of the result is that the three events have the same muon pair invariant mass, 214.3 MeV. HyperCP estimates the probability for this clustering at 0.8\% using a “form factor” distribution for the standard model expectations. This observation invites two calculations and we report on the results in this talk. First we present the best possible prediction for the Standard Model expectation. Since there are no known particles of mass 214 MeV, we do not expect a peak at that muon pair invariant mass. However, we need to know whether the SM distribution is narrower or wider than the form used by HyperCP to assess the significance of the clustering. Even if the three events represent new physics, it is necessary to know the SM level in order to determine if HyperCP should have seen events at other values of \( m_{\mu\mu} \).

The second calculation involves assuming that the observed events are indeed evidence for a new particle and confronting this observation with existing constraints from kaon and B physics. In particular we study the conditions under which the observation is consistent with a light Higgs boson and find an explicit candidate for the new particle: the lightest CP-odd Higgs boson in the NMSSM, the \( A_1^0 \).

2 Standard Model Calculation

We first present the ingredients that enter the calculation within the SM. The short distance contribution is too small to explain these events by four orders of magnitude, this decay is long distance dominated as is the case in similar kaon modes.

The long distance contributions to \( \Sigma^+ \rightarrow p\mu^+\mu^- \) can be pictured schematically as arising from the \( \Sigma^+ \rightarrow p\gamma^* \) process. There are four independent form factors allowed by electromagnetic
gauge invariance,
\[ \mathcal{M}(B_i \to B_f \gamma^*) = -eG_F \bar{B}_f \left[ i\sigma^\mu\nu q_\mu (a + b\gamma_5) + (q^2\gamma^\nu - q^\nu q)(c + d\gamma_5) \right] B_i \varepsilon_\nu. \tag{1} \]

Two of the form factors, \(a(q^2)\) and \(c(q^2)\), are parity conserving whereas \(b(q^2)\) and \(d(q^2)\) are parity violating. In addition, two of the form factors are non-zero at \(q^2 = 0\) and contribute to the radiative decay \(\Sigma^+ \to p\gamma\): \(a(0)\) and \(b(0)\). All four form factors are complex and receive imaginary parts from \(N\pi\) intermediate states.

We estimate these imaginary parts by taking the weak vertex \(\Sigma^+ \to p\mu^+\mu^-\) from experiment and using the \(N\pi \to p\gamma\) scattering at lowest order in \(\chi PT\) (both conventional and heavy baryon). We check that our calculations agree with the existing ones at \(q^2 = 0\).

To estimate the real part of the form factors we use \(a(0)\) and \(b(0)\), as determined from the width and decay distribution of the radiative decay \(\Sigma \to p\gamma\) up to a discrete ambiguity. We then assume that value for the range of \(q^2\) needed. This is consistent with our finding that the imaginary parts of the form factors are smooth and slowly varying over the \(q^2\) range of interest.

Finally, the real parts of \(c(q^2)\) and \(d(q^2)\) are obtained using a vector meson dominance model.

There is some uncertainty in the calculation, but the resulting range, \(1.6 \times 10^{-8} \leq B(\Sigma^+ \to p\mu^+\mu^-)_{SM} \leq 9.0 \times 10^{-8}\), is in good agreement with the measured rate, \(B(\Sigma^+ \to p\mu^+\mu^-) = (8.6^{+6.6}_{-5.3} \pm 5.5) \times 10^{-8}\). The predicted \(m_{\mu\mu}\) distribution shows no peaks near 214 MeV (or elsewhere) and is slightly flatter than the form factor used by HyperCP. This leads us to conclude that the probability of having the three events at the same invariant mass is about 0.5%.

Furthermore, the lower end of the range predicted for the rate is consistent with no events for HyperCP, allowing for the possibility of all three events being consistent with new physics.

## 3 A new Particle with mass 214 MeV?

We now turn to the interpretation of the 3 HyperCP events as a new particle \(^1\) with \(M_{P^0} = 214.3\ MeV\) and \(B(\Sigma^+ \to p\mu^+\mu^-)_{P^0} = (3.1^{+2.3}_{-1.9} \pm 1.5) \times 10^{-8}\). The observation implies that this hypothetical new light state, \(P^0\), is short lived, does not interact strongly, is narrow and decays only into \(\mu^+\mu^-\), \(e^+e^-\) or \(\gamma\gamma\), and has a \(\Delta S = 1, \Delta I = 1/2\) coupling to \(s\bar{d}\) quarks. There are three questions to be answered and we address them in order. Why hasn’t it been seen before? Is there a candidate for such a state? Where else could it be observed?

### 3.1 Why hasn’t it been seen before?

The most stringent constraint on a possible new particle \(P^0\) is its non-observation in kaon decay. After all, the modes \(K \to \pi\mu^+\mu^-\) proceed via the same quark level transition as \(\Sigma^+ \to p\mu^+\mu^-\): \(s \to d\mu^+\mu^-\). Of the three experiments that have studied these modes: BNL865 \(^4\), HyperCP \(^5\) and NA48 \(^6\) the one with most statistics was BNL865 \(^4\) with 430 events, 30 of which were in their lowest bin \(2m_{\mu\mu} \lesssim m_{\mu\mu} \lesssim 225\ MeV\) where the signal would have been observed. Their observation shows no peaks in the \(m_{\mu\mu}\) distribution, which is consistent with long distance SM physics. On that basis, the most optimistic scenario for the new physics hypothesis is to assume that all the 30 events in the first bin were due to \(P^0\) which leads to a 95% confidence limit bound \(B(K^+ \to \pi^+P^0) \leq 8.7 \times 10^{-9}\) (assuming that statistical errors dominate). This translates into a rate for \(\Sigma^+ \to pP^0\) some 25 times too small to explain the HyperCP events. Similar results are obtained from the other kaon experiments, none of which saw a peak in their \(m_{\mu\mu}\) distribution.

Another constraint arises from the non-observation of the hypothetical new particle in \(b \to s\mu^+\mu^-\). In this case both Belle and BaBar \(^8\) have results that can be interpreted as a 95% confidence level bound \(B(B \to X_sP^0) \leq 8 \times 10^{-8}\).

In Figure 1, we can see schematically how it is possible for the new state to be observed in \(\Sigma\) decay while not in \(K^+\) decay: the kaon decay modes with only one pion in the final state only

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constrain the effective $|\Delta S| = 1$ scalar coupling of the new state whereas the $\Sigma$ decay is sensitive also to the effective $|\Delta S| = 1$ pseudoscalar coupling. Any viable model for $P^0$ will then have an effective scalar coupling about 25 times smaller than the corresponding pseudoscalar coupling.

In a similar manner, the constraints from $B$ decay require that the effective $bs$ coupling of $P^0$ be about an order of magnitude smaller than the corresponding $sd$ coupling scaled by $m_b/m_s$ and $(V_{ts}V_{tb}^\ast)/(V_{ts}V_{td}^\ast)$. The latter scaling is the appropriate one for one-loop Higgs penguins dominated by a top-quark and a $W$ boson in the intermediate state. A successful model for $P^0$ can not have these penguin diagrams dominating the effective FCNC of $P^0$ to down-type quarks.

We have also considered additional processes that can, in principle, constrain the interactions of the hypothetical $P^0$. $K - \bar{K}$ mixing allows an effective pseudoscalar coupling up to 50 times as large as required to explain the 3 HyperCP events. $K_L \rightarrow \mu^+\mu^-$ combined with the muon $g - 2$ allows a $P^0$ coupling to muons $g_{P\mu} \lesssim 5 \times 10^{-4}$ which interestingly is about $m_{\mu}/v^{9,10}$.

3.2 Is there a candidate for $P^0$?

The possibility that $P^0$ is a light sgoldstino has been explored to some extent in the literature. Here, we pursue the possibility that $P^0$ is a light Higgs boson. For detailed phenomenology of kaon and hyperon decays involving a light Higgs particle it is necessary to recall that there are two types of contributions that are generally of similar size. There are two-quark “Higgs penguin” contributions that arise at one loop order and depend on the details of the flavor changing sector of the model. There are also “four-quark” contributions arising from a tree-level, SM $W$ mediated $|\Delta S| = 1$ decay, in which the light Higgs is radiated from any of the $u,d,s$ quarks or the $W$ boson via the tree-level flavor diagonal couplings of the Higgs. Both of these contributions can be calculated in chiral perturbation theory, and we do so at leading order. Given our discussion in the previous section we concentrate on CP-odd or pseudoscalar Higgs bosons.

One possible candidate for $P^0$ is the $A^0_1$ of the NMSSM. The Higgs sector of the NMSSM contains the usual two Higgs doublets $H_D$ and $H_U$ that appear in the MSSM plus the Higgs singlet $N$. In the physical spectrum there are two CP-odd scalars, of which the $A^0_1$ is the lightest. It has been proposed in the literature that this $A^0_1$ can be naturally light due to a global $U(1)$ symmetry.

The main features of the couplings of the $A^0_1$ to SM fields are as follows. Its coupling to $Zh$ ($h$ being the lightest CP even Higgs) is suppressed by $\tan\beta$ with respect to the MSSM $ZhA$ coupling allowing an evasion of LEP bounds in the large $\tan\beta$ regime. Its couplings to quarks are also suppressed by $\tan\beta$ with respect to those of the $A$ in the MSSM. This results, for large $\tan\beta$, in negligible couplings to up-type quarks. The couplings to down-type quarks are
Figure 2: Parameter space for $m_{\tilde{u}} - m_{\tilde{c}}$ and $m_{\tilde{u}}/(\lambda x)$ where $A_0^1$ can explain the HyperCP events (gray regions) and simultaneously satisfy the kaon bounds (black regions). The horizontal axis corresponds to parameters in the chargino mass matrix.

independent of tan $\beta$ and can be written in terms of one parameter, $l_d$, which can be of order one $^{14}$: $\mathcal{L} = -l_dm_d d\gamma_5 d(iA_0^1)/v - l_dm_\ell \ell\gamma_5 (iA_0^1)/v + \cdots$.

The four-quark contributions to $A_0^1$ production in light meson and hyperon decay are thus proportional to $l_d$ and independent of other parameters in the model. It is then straightforward to compute these contributions to the HyperCP case. We find $^{15}$, $\mathcal{B}_{4q}(\Sigma^+ \rightarrow p A_0^1) = 1.7 \times 10^{-7}|l_d|^2$, which matches the central value of the HyperCP result for $l_d \sim 0.4$. The bad news is that this then leads to $\mathcal{B}_{4q}(K^+ \rightarrow \pi^+ A_0^1) \sim 10^{-6}$, two orders of magnitude larger than the limit from BNL E865. The conclusion illustrated by this calculation is that it is relatively easy to have a light Higgs that matches the HyperCP observation but it is very hard to avoid seeing it in kaon decay as well.

However, there are also the two-quark contributions to the amplitudes and it is possible to arrange a cancellation between amplitudes that satisfies the kaon bounds. The two-quark contributions are much more model dependent than the four-quark contributions, but also suffer from additional constraints due to non-observation of $P^0$ in $B$ decay. We have not performed a full parameter scan, but rather illustrated that it is possible to satisfy all constraints. To this effect we start with the specific model considered by Hiller $^{14}$ and modify it accordingly. To suppress the FCNC in $B$ decay we consider $m_{\tilde{t}} = m_{\tilde{\tau}}$ and negligible squark mixing. The strength of the two-quark contribution to kaon decay is then tuned with $m_{\tilde{u}} - m_{\tilde{\tau}}$. We further consider (large) tan $\beta = 30$, $m_{\tilde{t}} \sim 2.5$ TeV and $-\lambda x = 150$ GeV to obtain neutralino masses in the 100-1500 GeV range $^{15}$. In Figure 2 we show our results $^{15}$: the light shaded region corresponds to parameters that reproduce the HyperCP observation. The dark shaded region corresponds to those points that also satisfy the kaon bounds. As mentioned before the overlapping region is significantly smaller due to the cancellation required to satisfy the kaon bounds.

3.3 Where else can $P^0$ be observed?

Finally, we explore other processes that can test the new particle hypothesis for the HyperCP result. We begin by considering only the effect of two-quark operators, assuming that the existing kaon bounds are bypassed because the effective $sd$ coupling is pseudoscalar. In this case the new state would show up in kaon decay modes with two pions in the final state and we can easily derive from the HyperCP measurement that (the errors reflect the experimental error only)$^9$

$$\mathcal{B}(K_L \rightarrow \pi^+ \pi^- P^0) \approx (1.8^{+1.6}_{-1.4}) \times 10^{-9}$$
$$\mathcal{B}(K_L \rightarrow \pi^0 \pi^0 P^0) \approx (8.3^{+7.5}_{-6.6}) \times 10^{-9}.$$ (2)
Both of these represent very significant enhancements over the corresponding SM rates and may be accessible to KTeV or NA48. In a similar manner this scenario results in $^9$,$^7$:

$$B(\Omega^- \rightarrow \Xi^- P^0) \approx (2.0^{+1.6}_{-1.2}) \times 10^{-6}. \quad (3)$$

The best upper bound for this mode, also from HyperCP $^{16}$, is $6.1 \times 10^{-6}$.

If the new state $P^0$ is a light Higgs, then there are other processes that are sensitive only to its flavor diagonal couplings $^{19}$ (or four-quark operators). For example the modes $V \rightarrow \gamma A_1^0$ have been proposed in the literature $^{17}$. The results are that $B(\Upsilon_{1S} \rightarrow \gamma A_1^0)$ can reach about $1 \times 10^{-4} l_d^2$ and may be accessible to the B factories. Similarly $B(\phi \rightarrow \gamma A_1^0)$ can reach $1.4 \times 10^{-8} l_d^2$ and may be accessible to DAΦNE $^{17}$. In a similar spirit we have proposed the modes $\eta \rightarrow \pi \pi A_1^0$ where we can predict $^{18} B(\eta \rightarrow \pi^+ \pi^- A_1^0) = 5.4 \times 10^{-7} l_d^2$, again possibly accessible to DAΦNE.

When the four-quark contributions are added to the two-quark contributions in the NMSSM (using parameters as in Hiller $^{14}$ and Xiandong $^{20}$) the results of Eq. 2 are modified. An example of the resulting predictions for the rate of the kaon modes is shown in Fig. 3. Full details can be found in the paper $^{18}$, but the $x$-axis is related to the strength of the two-quark contribution though an effective $g_P$ and the strength of the four-quark contribution is kept fixed. The dotted curves result from the two-quark contributions alone. The shaded (pink) bands indicate the allowed ranges of $C_L-C_R$ when the two and four-quark contributions have the same sign $^{18}$. Each vertical (green) dashed line corresponds to the special case $^{15}$ of chargino dominated penguins.

4 Conclusions

The decay $\Sigma^+ \rightarrow p\mu^+\mu^-$ within the SM is long distance dominated and the predicted rate is in the right range to explain the HyperCP observation. However, the predicted $m_{\mu\mu}$ distribution makes it unlikely to find the three events at the same mass ($P \lesssim 0.8\%$). Existing constraints from kaon and B physics allow a new particle interpretation of the HyperCP result provided that the FCNC couplings of the new particle are mostly pseudoscalar and smaller for $b \rightarrow s$ transitions than naive scaling with CKM angles would predict.

The NMSSM has a CP-odd Higgs boson, the $A_1^0$ that could be as light as the required 214 MeV. Its diagonal couplings to quarks and muons in the large $\tan\beta$ limit can have the right size as well. There are several modes that can test this hypothesis independently from the details of the flavor changing sector of the model: $\Upsilon_{1S} \rightarrow \gamma A_1^0$, $\phi \rightarrow \gamma A_1^0$ and $\eta \rightarrow \pi \pi A_1^0$.

It is harder to suppress the effective scalar $sd$ coupling that appears in this model to the level required to satisfy the existing kaon bounds, but it is possible for certain values of the relevant
parameters. The measurement of one of the modes $K_L \rightarrow \pi\pi\mu^+\mu^-$ can confirm or refute this scenario.

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References

MEASUREMENTS FROM KTeV OF RARE DECAYS OF THE $K^0_L$ AND $\pi^0$

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The KTeV collaboration at Fermi National Accelerator Laboratory has recently completed searches for and measurements of several decay modes of the neutral kaon and pion. These include new searches for lepton flavor violating decays (which have not been seen), and a new study of the parity properties of the decay $\pi^0 \rightarrow e^+e^-e^+e^-$. 

1 The KTeV Detector

Fermilab’s KTeV detector (Fig. 1) was constructed for Experiments 799 and 832. The two experiments were designed to concentrate on different aspects of neutral kaon physics: E799 on rare decays of the $K_L$ and E832 on measurement of Re($\epsilon'/\epsilon$). A primary proton beam with energy 800 GeV struck a BeO target at a targeting angle of 4.8 mrad, and collimation and sweeping magnets produced two parallel neutral hadron beams. The beams entered a 60 m long vacuum decay region, which ended at a Mylar-Kevlar vacuum window. Decay products were tracked with a series of drift chambers surrounding a dipole analysis magnet. Downstream of the drift chambers were a series of transition radiation detectors (TRD) (in E799 only) and a pure CsI electromagnetic calorimeter, an active hadron beam absorber, and a set of muon detectors behind steel shielding. Photon veto detectors surrounded the fiducial volume in the transverse directions. The detector is described in more detail in Ref. 1.

2 The decay $\pi^0 \rightarrow e^+e^-e^+e^-$ and the parity of the $\pi^0$

The neutral pion’s parity has historically been studied in two ways: indirectly via the cross-section of $\pi^-$/capture on deuterons$^{2,3}$, or directly via the double Dalitz decay $\pi^0 \rightarrow e^+e^-e^+e^-$. While both sets of results are consistent with the negative parity, the direct measurement has only 3.6 $\sigma$ significance. KTeV has now reported results$^5$ that conclusively confirm the negative $\pi^0$
Figure 1: The KTeV spectrometer as configured for E799.

Figure 2: Lowest order Feynman diagram for $\pi^0 \rightarrow e^+e^-e^+e^-$. The direct contribution is shown; a second diagram exists with $e_1^+$ and $e_2^-$ exchanged.

parity as well as the first-ever searches for parity and $CPT$ violaton, and the first measurements of the electromagnetic form factor, in this mode.

The $\pi^0 \rightarrow e^+e^-e^+e^-$ decay proceeds through a two-photon intermediate state (Fig. 2). The most general interaction Lagrangian for the $\pi^0 \rightarrow \gamma^*\gamma^*$ transition can be written:

$$L \propto C_{\mu\nu\rho\sigma} F^{\mu\nu} F^{\rho\sigma} \Phi$$ (1)

where $F^{\mu\nu}$ and $F^{\rho\sigma}$ are the photon fields, $\Phi$ is the pion field, and the coupling has the form

$$C_{\mu\nu\rho\sigma} \propto f(x_1, x_2)[\cos \zeta e_{\mu\nu\rho\sigma} + \sin \zeta e^{i\delta} (g_{\mu\rho}g_{\nu\sigma} - g_{\mu\sigma}g_{\nu\rho})].$$ (2)

The first term in $C_{\mu\nu\rho\sigma}$ is the expected pseudoscalar coupling and the second term introduces a scalar coupling with a mixing angle $\zeta$ and a phase difference $\delta$. Nuclear parity violation would introduce a nonzero $\zeta$, while $CPT$ violation would cause the phase $\delta$ to be nonzero. We assume the standard parity-conserving form for the $\gamma^* \rightarrow e^+e^-$ conversion.

The form factor $f(x_1, x_2)$ is expressed in terms of the momentum transfer of each of the virtual photons, or equivalently the invariant masses of the two Dalitz pairs: $x_1 \equiv (m_{e_1^+e_2^-}/M_{\pi^0})^2$, $x_2 \equiv (m_{e_1^+e_2^-}/M_{\pi^0})^2$. In calculating the phase space variables for an individual event, there is an intrinsic ambiguity in assigning each electron to a positron to form a Dalitz pair. KTeV’s analysis uses a matrix element model that includes the exchange diagrams and therefore avoids
the need to enforce a pairing choice. The form factor is parametrized using a model based on that of D’Ambrosio, Isidori, and Portolés (DIP) \(^7\), but with an additional constraint that ensures the coupling vanishes at large momenta \(^8\). In terms of the remaining free parameters, the form factor is:

\[
 f_{\text{DIP}}(x_1, x_2; \alpha) = \frac{1 - \mu(1 + \alpha)(x_1 + x_2)}{(1 - \mu x_1)(1 - \mu x_2)},
\]

where \(\mu = M_p^2/M_\pi^2 \approx 0.032\).

The parity properties of the decay can be extracted from the angle \(\phi\) between the planes of the two Dalitz pairs in Fig. 2, where pair 1 is defined as having the smaller invariant mass. The distribution of this angle from the dominant direct contribution has the form \(d\Gamma/d\phi \sim 1 - A \cos(2\phi) + B \sin(2\phi)\), where \(A \approx 0.2 \cos(2\xi)\) and \(B \approx 0.2 \sin(2\xi) \cos \delta\). A pure pseudoscalar coupling, therefore, would produce a negative \(\cos(2\phi)\) dependence.

The branching ratio measurement, which we describe here first, makes use of a normalization mode in which two pions decay via \(\pi^0 \rightarrow e^+e^-\gamma\) and the third \(\pi^0 \rightarrow \gamma\gamma\). This “double single-Dalitz” mode, denoted \(K_L \rightarrow \pi^0\pi^0\pi^0\pi^0\) where \(\pi^0\) refers to \(\pi^0 \rightarrow e^+e^-\gamma\), has the same final state particles as the signal mode. Both modes are fully reconstructed in the detector and the total invariant mass is required to match the kaon’s. The two modes are distinguished by a \(\chi^2\) formed of the three reconstructed \(\pi^0\) masses. This serves to identify the best pairing of particles for a given decay hypothesis, as well as to select the more likely hypothesis of the two. The similarity of these modes allows cancellation of most detector-related systematic effects in the branching ratio measurement, but also allows each mode to be a background to the other.

Radiative corrections complicate the definition of the Dalitz decays in general. We define the signal mode \(\pi^0 \rightarrow e^+e^-\gamma\) to be inclusive of radiative final states where the squared ratio of the invariant mass of the four electrons to the neutral pion mass \(x_{4e} \equiv (M_{4e}/M_\pi^2)^2\) is greater than 0.9, while events with \(x_{4e} < 0.9\) (approximately 6% of the total rate) are treated as \(\pi^0 \rightarrow e^+e^-\pi^0\). For normalization, the decay \(\pi^0 \rightarrow e^+e^-\gamma\) is understood to include all radiative final states, for consistency with previous measurements of this decay \(^9\). Radiative corrections in this analysis are taken from an analytic calculation to order \(O(\alpha^2)\) \(^6\).

Radiative corrections complicate the definition of the Dalitz decays in general. The signal mode \(\pi^0 \rightarrow e^+e^-\gamma\) is defined to be inclusive of radiative final states where the squared ratio of the invariant mass of the four electrons to the neutral pion mass \(x_{4e} \equiv (M_{4e}/M^2)^2\) is greater than 0.9, while events with \(x_{4e} < 0.9\) (approximately 6% of the total rate) are treated as \(\pi^0 \rightarrow e^+e^-\gamma\). Radiative corrections in this analysis are taken from an analytic calculation to order \(O(\alpha^2)\) \(^6\).

The final event sample contains 30 511 signal candidates with 0.6% residual background and 141 251 normalization mode candidates with 0.5% background (determined from the Monte Carlo simulation). The background in the signal sample is dominated by mistagged events from the normalization mode. v KTeV finds the following the ratio of decay rates:

\[
\frac{B_{\pi^0 eee \gamma}^{>0.9} \cdot B_{ee\gamma}}{B_{ee\gamma}} = 0.2245 \pm 0.0014(\text{stat}) \pm 0.0009(\text{syst}).
\]

The \(\pi^0 \rightarrow e^+e^-\gamma\) branching ratio can be calculated from the double ratio using the known values \(B_{ee\gamma} = 0.9980 \pm 0.0003\) and \(B_{ee\gamma} = (1.198 \pm 0.032) \times 10^{-2}\) \(^{10}\). This yields \(B_{ee\gamma}^{>0.9} = (3.26 \pm 0.18) \times 10^{-5}\), where the error is dominated by the uncertainty in the \(\pi^0 \rightarrow e^+e^-\gamma\) branching ratio. KTeV uses the radiative corrections model \(^6\) to extrapolate to all radiative final states, finding:

\[
\frac{B_{ee\gamma} \cdot B_{ee\gamma}}{B_{ee\gamma}^{>0.9}} = 0.2383 \pm 0.0015(\text{stat}) \pm 0.0010(\text{syst}),
\]

and \(B_{ee\gamma} = (3.46 \pm 0.19) \times 10^{-5}\). This branching ratio result is in good agreement with previous measurements \(^4\).
The parameters of the $\pi^0\gamma^*\gamma^*$ coupling are found by maximizing an unbinned likelihood function composed of the differential decay rate in terms of ten phase-space variables. The first five are $(x_1, x_2, y_1, y_2, \phi)$, where $x_1$, $x_2$, and $\phi$ are described above and the remaining variables $y_1$ and $y_2$ describe the energy asymmetry between the electrons in each Dalitz pair in the $\pi^0$ center of mass. The remaining five are the same variables, but calculated with the opposite choice of $e^+e^-$ pairings. The likelihood is calculated from the full matrix element including the exchange diagrams and $\mathcal{O}(\alpha^2)$ radiative corrections.

The fit yields the DIP $\alpha$ parameter and the (complex) ratio of the scalar to the pseudoscalar coupling. For reasons of fit performance, the parity properties are fit to the equivalent parameters $\kappa$ and $\eta$, where $\kappa + i\eta \equiv \tan \zeta e^{i\delta}$. The shape of the minimum of the likelihood function indicates that the three parameters $\alpha$, $\kappa$, and $\eta$ are uncorrelated. Acceptance-dependent effects are included as a normalization factor calculated from Monte Carlo simulations.

Systematic error sources on $\alpha$ and $\kappa$ are similar to those for the branching ratio measurement. The dominant systematic error is due to variation of cuts, resulting in a total systematic error of 0.9 and 0.011 on $\alpha$ and $\kappa$ respectively. For the $\eta$ parameter, the primary uncertainty results from the resolution on the angle $\phi$ between the two lepton pairs. This behavior was studied with Monte Carlo simulation and a correction was calculated. The uncertainty on this correction results in a systematic error of 0.031.

The $\phi$ distribution is shown in Fig. 3. For plotting the data a unique pairing of the four electrons is chosen such that $x_1 < x_2$ and the product $x_1x_2$ is minimized: this choice represents the dominant contribution to the matrix element. It is clear that the pseudoscalar coupling dominates, as expected, with no evidence for a scalar component. The distributions of all five phase space variables agree well with the Monte Carlo simulation.

The $\phi$ distribution is shown in Fig. 3. For plotting the data a unique pairing of the four electrons is chosen such that $x_1 < x_2$ and the product $x_1x_2$ is minimized: this choice represents the dominant contribution to the matrix element. It is clear that the pseudoscalar coupling dominates, as expected, with no evidence for a scalar component. The distributions of all five phase space variables agree well with the Monte Carlo simulation.

The parameters $\kappa$ and $\eta$ are transformed into limits on the pseudoscalar-scalar mixing angle $\zeta$ under two hypotheses. If $CPT$ violation is allowed, then the limit is set by the uncertainties in $\eta$, resulting in $\zeta < 6.9^\circ$ at the 90% confidence level. If instead, $CPT$ conservation is enforced, $\eta$ must be zero, and the limit derives from the uncertainties on $\kappa$, resulting in $\zeta < 1.9^\circ$, at the same confidence level. These limits on $\zeta$ limit the magnitude of the scalar component of the decay amplitude, relative to the pseudoscalar component, to less than 12.1% in the presence of $CPT$ violation, and less than 3.3% if $CPT$ is assumed conserved. The limits on scalar contributions apply to all $\pi^0$ decays with two-photon intermediate or final states.

This analysis confirms the negative parity of the neutral pion with much higher statistical significance than the previous result, and places tight limits on nonstandard scalar and $CPT$-violating contributions to the $\pi^0 \rightarrow e^+e^-e^+e^-$ decay.
3 Lepton Flavor Violation

Lepton Flavor Violation (LFV) in weak decays is a key signature of several beyond-Standard Model physics scenarios. Supersymmetry\textsuperscript{11}, new massive gauge bosons\textsuperscript{12,13}, and technicolor\textsuperscript{14} all can lead to LFV decays which might be within reach of current experiments. Searches in $K_L$ decays are complementary to searches in the charged lepton sector, since $K_L$ decays probe the $s \rightarrow d\mu e$ transition\textsuperscript{12}. KTeV-E799 has searched for the decays $K_L \rightarrow \pi^0\mu^\pm e^\mp$ and $\pi^0 \rightarrow \mu^\pm e^\mp$, and has made the first reported search for $K_L \rightarrow \pi^0\pi^0\mu^\pm e^\mp$\textsuperscript{15}.

In each case, the analysis required two charged tracks, one of which was identified as a muon and the other an electron. The key detector elements for particle identification were $E/p$ in the CsI calorimeter, response of the TRD, and muon hodoscopes downstream of the muon filter steel. Clusters in the CsI with no tracks pointing to them were considered photons.

3.1 $K_L \rightarrow \pi^0\mu^\pm e^\mp$

The dominant background for $K_L \rightarrow \pi^0\mu^\pm e^\mp$ was the decay $K_L \rightarrow \pi^\pm e^\mp\nu_e$ ($K_{e3}$), with a $\pi^\pm$ decay or punch through to the muon hodoscopes, accompanied by two accidental photons faking a $\pi^0$. Since accidental photons were often accompanied by other accidental activity, we removed events with evidence of additional in-time activity in the detector. Additionally, the two photons were required to form a good $\pi^0$ mass, and the square of the $\pi^0$ momentum in the $K_L$ rest frame was required to be positive and therefore physical.

The signal and control regions were defined using a likelihood variable $L$ derived from $p_t^2$, the sum of the momentum components of all final-state particles perpendicular to the kaon flight line, and $M_{\pi^0\mu e}$, the invariant mass of the $\pi^0\mu e$ system. The signal (control) region was defined by a cut on $L$ chosen to retain 95% (99%) of signal Monte Carlo events after all other cuts. Expected background levels were 0.66 ±0.23 events in the signal region and 4.21 ±0.53 events in the control region. Both the signal and control regions were blind during the analysis. Figure 4 shows the $p_t^2 - M_{\pi^0\mu e}$ plane after all cuts: five events were found in the control region and zero in the signal. The resulting limit is $B(K_L \rightarrow \pi^0\mu^\pm e^\mp) < 7.56 \times 10^{-11}$ at 90% CL, a factor of 82 improvement over the previous best limit for this mode.\textsuperscript{16}

![Figure 4: Surviving events in the $p_t^2 - M_{\pi^0\mu e}$ plane for the $K_L \rightarrow \pi^0\mu^\pm e^\mp$ search data. The signal and control regions are shown as the inner and outer solid contours.](image-url)
3.2 Other lepton flavor violating modes

KTeV has also searched for the decay $K_L \to \pi^0\pi^0\mu^\pm e^\mp$. Reconstructing a second $\pi^0$ greatly reduces the backgrounds, so some particle identification and anti-accidental cuts were relaxed to improve the signal acceptance. A similar analysis, including a cut on a kinematic likelihood variable, yielded no events in either the control region or the signal region. This resulted in a limit $B(K_L \to \pi^0\pi^0\mu^\pm e^\mp) < 1.64 \times 10^{-10}$. This is the first limit reported for this decay.

The decay chain $K_L \to \pi^0\pi^0\pi^0$, $\pi^0 \to \mu^\pm e^\mp$ gives the same final state particles as $K_L \to \pi^0\pi^0\mu^\pm e^\mp$, and therefore the same analysis procedure applies with the additional requirement that the invariant mass $M_{\mu e} \approx M_{\pi^0}$. Since no events were found, the limit is $B(\pi^0 \to \mu^\pm e^\mp) < 3.59 \times 10^{-10}$. This limit on $\pi^0 \to \mu^\pm e^\mp$ is equally sensitive to both charge modes, while the previous best limits were not.

Assuming equal contributions from both charge combinations, KTeV’s result is about a factor of two better than the previous best limit on $\pi^0 \to \mu^- e^+$ and about a factor of 10 greater than the previous best limit on $\pi^0 \to \mu^+ e^-$. The measurement of $\pi^0 \to e^+ e^- e^+ e^-$ represents the best direct determination of the parity of the $\pi^0$ and the first searches for nonstandard parity and CPT violation in this mode. It also yields the best branching ratio and the first measurement of the form factor in this mode. The limits on lepton flavor violation are now the most stringent in the world for these decay modes.

4 Conclusion

KTeV has completed several measurements recently on the decays of neutral $K$ and $\pi$ mesons. The measurement of $\pi^0 \to e^+ e^- e^+ e^-$ represents the best direct determination of the parity of the $\pi^0$ and the first searches for nonstandard parity and CPT violation in this mode. It also yields the best branching ratio and the first measurement of the form factor in this mode. The limits on lepton flavor violation are now the most stringent in the world for these decay modes.

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References

Recent results from KLOE

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In this report I will present the recent results on $K$ mesons from the KLOE experiment at the DAFNE $e^+e^-$ collider working at the center of mass energy $W \sim m_\phi \sim 1.02$ GeV. The $\phi$ mesons are produced essentially at rest and decay to $K_SK_L (K^+K^-) \sim 34\%$ ($\sim 49\%$) of the times. The $K$ mesons are produced in a pure $J^{PC} = 1^{--}$ coherent quantum state, so that observation of a $K_S (K^+)$ in an event signals (tags) the presence of a $K_L (K^-)$ and vice-versa: highly pure, almost monochromatic, back-to-back $K_S (K^+)$ and $K_L (K^-)$ beams can be obtained. Moreover $K_S$ and $K_L$ are distinguishable on the basis of their decay length: $\lambda_S \sim 0.6$ cm and $\lambda_L \sim 340$ cm.

1 The KLOE experiment

The KLOE detector operates at DAΦNE, an $e^+e^-$ collider working at the center of mass energy $W \sim m_\phi \sim 1.02$ GeV. The $\phi$ mesons are produced essentially at rest and decay to $K_SK_L (K^+K^-) \sim 34\%$ ($\sim 49\%$) of the times. The $K$ mesons are produced in a pure $J^{PC} = 1^{--}$ coherent quantum state, so that observation of a $K_S (K^+)$ in an event signals (tags) the presence of a $K_L (K^-)$ and vice-versa: highly pure, almost monochromatic, back-to-back $K_S (K^+)$ and $K_L (K^-)$ beams can be obtained. Moreover $K_S$ and $K_L$ are distinguishable on the basis of their decay length: $\lambda_S \sim 0.6$ cm and $\lambda_L \sim 340$ cm.
The KLOE detector consists essentially of a drift chamber (DC), surrounded by an electromagnetic calorimeter (EMC). The DC is a cylinder of 4 m diameter and 3.3 m in length which constitutes a large fiducial volume for $K_L$ decays ($\sim 1/2 \lambda_L$). The momentum resolution for tracks at large polar angle is $\sigma_p/p \leq 0.4\%$. The EMC is a lead-scintillating fiber calorimeter consisting of a barrel and two endcaps, which cover 98% of the solid angle. The energy resolution is $\sigma_E/E \sim 5.7\%/\sqrt{E/(\text{GeV})}$. The intrinsic time resolution is $\sigma_T = 54\text{ps}/\sqrt{E/(\text{GeV})} \pm 50\text{ps}$. A superconducting coil surrounding the barrel provides a 0.52 T magnetic field.

The present report is based on a first data sample of $\sim 500\text{ pb}^{-1}$, except for quantum coherence, CPT and Lorentz symmetry tests; at present KLOE has about 2.2 $\text{ fb}^{-1}$ on disk.

### 2 $V_{us}$ determination

In the Standard Model, the coupling of the $W$ boson to the weak charged current is written as

$$\frac{g}{\sqrt{2}} W^\pm_T (\bar{U} L V_{CKM}^\alpha D_L + \bar{e}_L \gamma^\alpha \nu_e L + \bar{\mu}_L \gamma^\alpha \nu_\mu L + \bar{\tau}_L \gamma^\alpha \nu_\tau L) + \text{h.c.,}$$

where $U^T = (u, c, t)$, $D^T = (d, s, b)$ and $L$ is for lefthanded. In the coupling above there is only one coupling constant for leptons and quarks. Quarks are mixed by the Cabibbo-Kobayashi-Maskawa matrix, $V_{CKM}$, which must be unitary.

The most precise check on the unitarity of the $V_{CKM}$ matrix is provided by measurements of $|V_{us}|$ and $|V_{ud}|$, the contribution of $V_{ub}$ being at the level of $10^{-5}$. $|V_{us}|$ may be extracted by the measurements of the semileptonic decay rates, fully inclusive of radiation, which are given by:

$$\Gamma(K_{\ell 3(\gamma)}) = \frac{C_K^2 G_F^2 M_K^5}{192\pi^3} S_{\text{EW}} |V_{us}|^2 \left| f_+(0) \right|^2 I_{K \ell} \left( 1 + \delta_{SU}^{(2)} + \delta_{EM}^{(2)} \right)^2.$$  

(2)

In the above expression, the index $K$ denotes $K^0 \to \pi^\pm$ and $K^\pm \to \pi^0$ transitions, for which $C_K^2 = 1$ and $1/2$, respectively. $M_K$ is the appropriate kaon mass, $S_{\text{EW}}$ is the universal short-distance electroweak correction \(^3\) and $\ell = e, \mu$. Following a common convention, $f_+(0) \equiv f_+^{K^0\pi^-}(0)$. The mode dependence is contained in the $\delta$-terms: the long-distance electromagnetic (EM) corrections, which depend on the meson charges and lepton masses and the SU(2)-breaking corrections, which depend on the kaon species \(^4\). $I_{K \ell}$ is the integral of the dimensionless Dalitz-plot density over the physical region for non radiative decays and includes $|f_{+,0}(t)|^2$, the reduced form factor, defined below.

$|V_{us}|$ can be also extracted from $K \to \mu\nu$ decays using the relation

$$\frac{\Gamma(K_{\mu 2(\gamma)})}{\Gamma(\pi_{\mu 2(\gamma)})} = \frac{|V_{us}|^2}{|V_{ud}|^2} \frac{f_\pi^2}{f_K^2} \frac{m_K}{m_\pi} \frac{1-m_\mu^2/m_K^2}{1-m_\mu^2/m_\pi^2} \times (0.9930 \pm 0.0035),$$

(3)

where $f_\pi$ and $f_K$ are the pion- and kaon-decay constants and the uncertainty in the numerical factor is dominated from structure-dependent radiative corrections. This ratio can be combined with direct measurements of $|V_{ud}|$ to obtain $|V_{us}|$.

The measurement of $V_{us}$ from leptonic and semileptonic kaon decays allows both the test the unitarity of the CKM matrix and and the leptonic quark universality. Moreover the universality of electron and muon interactions can be tested by measuring the ratio $\Gamma(K \to \pi \mu \nu)$/$\Gamma(K \to \pi e \nu)$ and the comparison between the measurement of $V_{us}$ from leptonic decays and that from semileptonic decays allows to put bounds on new physics.

The experimental inputs to eq. 2 and 3 are the semileptonic and leptonic decay rates, fully inclusive of radiation, i.e. branching ratios (BR) and lifetimes, and the reduced form factors $f_+(t)$ and $f_0(t)$, whose behaviour as a function of $t$, the 4-momentum transfer squared $(P_K - p_\pi)^2$,
is obtained from the decay pion spectra. Details on the measurements and the treatment of correlations can be found in ref. 5. In this report I will present the recent measurement of the $K_{\mu 3}$ form factors, the charged kaon life time, the $\text{BR}(K^+ \to \pi^+ \pi^0)$ and the $\text{BR}(K^+ \to \pi^+ \pi^0)$.

### 3 $K_{\mu 3}$ from factors

The largest uncertainty in calculating $|V_{us}|$ from the decays rates is due to the difficulties in computing the matrix element $\langle \pi | J^{\text{had}}_\alpha | K \rangle$ which has the form:

$$\langle \pi | J^{\text{had}}_\alpha | K \rangle = f_+(0) \times ((P + p)_\alpha f_+(t) + (P - p)_\alpha (f_0(t) - f_+(t))\Delta_{K\pi}/t)$$

where $P(p)$ is the $K(\pi)$ momentum, $t = (P - p)^2$ and $\Delta_{K\pi} = M^2_K - m^2_\pi$. The above equation defines the vector and scalar form factors (FF) $f_+(t) = f_+(0)f_+(t)$ and $f_0(t) = f_+(0)f_0(t)$, which take into account the non point-like structure of the pions and kaons. The term $f_+(0)$ has been factored out, since the FFs must have the same value at $t = 0$. If the FFs are expanded in powers of $t$ up to $t^2$ as $\hat{f}_+(t) = 1 + \lambda' + |t/m^2| + 2 \lambda_+ |t/m^2|^2$, four parameters ($\lambda'_+, \lambda'^0_+, \lambda_0'$ and $\lambda^0_0$) need to be determined from the decay pion spectrum in order to be able to compute the phase-space integral. However, this parametrization of the form factors is problematic, because the values for the $\lambda$s obtained from fits to the experimental decay spectrum are strongly correlated 6. It is therefore necessary to obtain a form for $\hat{f}_0(t)$ and $\hat{f}_+(t)$ with at least $t$ and $t^2$ terms but with only one parameter. The Callan-Treiman relation 7 fixes the value of scalar FF at $t = \Delta_{K\pi}$ (the so-called Callan-Treiman point) to the ratio of the pseudoscalar decay constants $\hat{f}_K/\hat{f}_\pi$. $\hat{f}_0(\Delta_{K\pi}) = \frac{\hat{f}_K}{\hat{f}_\pi} \frac{1}{f_+(0)} + \Delta_{\text{CT}}$, where $\Delta_{\text{CT}}, \text{SU}(2)$-breaking correction 8, is of $O(10^{-3})$. A recent dispersive parametrization for the scalar form factor 9, $\hat{f}_0(t) = \exp\left[\frac{\Delta_{K\pi}}{m^2}(\ln C - G(t))\right]$, allows the constraint given by the Callan-Treiman relation to be exploited, such that $C = \hat{f}_0(\Delta_{K\pi})$ and $\hat{f}_0(0) = 1$. $G(t)$ is derived from $K\pi$ scattering data. As suggested in ref. 9, a good approximation to the dispersive parametrization is $\hat{f}_0(t) = 1 + \lambda_0 |t/m^2| + \lambda_0^2 + p_2 \left(\frac{t}{m^2}\right)^2 + \lambda_0^3 + 3p_2\lambda_0 + p_3 \left(\frac{t}{m^2}\right)^3$ with $p_2$ and $p_3$ given in ref. 9. Also for the vector FF we make use of a dispersive parameterization 10, twice subtracted at $t = 0$, $\hat{f}_+(t) = \exp\left[\frac{t}{m^2}(\Lambda_+ + H(t))\right]$, where $H(t)$ is obtained from $K\pi$ scattering data and $\Lambda_+$ has to be determined from the fit to experimental data. At KLOE energies clean and efficient

![Figure 1: Residuals of the fit (top plots) and $E_\nu$ distribution for data events superimposed on the fit result (bottom plot).](image)

$\pi/\mu$ separation, required to measure the $t$ spectrum, is difficult. The FF parameters have
been therefore obtained from fits to the distribution of the neutrino energy $E_\nu$ after integration over the pion energy. About 1.8 Million of $K_{\mu3}$ are selected by means of kinematic cuts, time of flight (TOF) measurements and calorimetric information. Details on the analysis can be found in ref. \textsuperscript{11}. Using the dispersive parameterizations for the vector and scalar FF’s and combing the $K_{\mu3}$ and $K_{e3}$ data, we find $\lambda_+ = (25.7 \pm 0.4 \pm 0.2 \text{param}) \times 10^{-3}$ and $\lambda_0 = (14.0 \pm 1.6 \pm 1.3 \pm 0.2 \text{param}) \times 10^{-3}$ with $\chi^2/\text{dof} = 2.6/3$ and a correlation coefficient of $-0.26$. The result of the fit on $K_{\mu3}$ data is shown in figure 1. Preliminary results based on $1 \text{fb}^{-1}$ have also been obtained and averaged with that presented above: $\lambda_+ = (26.0 \pm 0.5_{\text{stat+syst}}) \times 10^{-3}$ and $\lambda_0 = (15.1 \pm 1.4_{\text{stat+syst}}) \times 10^{-3}$.

4 \quad \tau(K^\pm), BR(K_{e3}^\pm) and BR(K^+ \rightarrow \pi^+\pi^0)

We have combined the recent published measurements of the semileptonic BRs and the charged kaon lifetime to use them in the evaluation of $|V_{us}|$.

At KLOE, two methods are used to reconstruct the proper decay time distribution for charged kaons. The first is to obtain the decay time from the kaon path length in the DC, accounting for the continuous change in the kaon velocity due to ionization energy losses. A fit to the proper-time distribution in the interval from $15\text{–}35\ \text{ns}$ $(1.6\tau_\pm)$ gives the result $\tau_\pm = 12.364 \pm 0.031_{\text{stat}} \pm 0.033_{\text{syst}}\ \text{ns}$. Alternately, the decay time can be obtained from the precise measurement of the arrival times of the photons from $K^+ \rightarrow \pi^+\pi^0$ decays. In this case, a fit to the proper-time distribution in the interval from $13\text{–}42\ \text{ns}$ $(2.3\tau_\pm)$ gives the result $\tau_\pm = 12.337 \pm 0.030_{\text{stat}} \pm 0.020_{\text{syst}}\ \text{ns}$. Taking into account the statistical correlation between these two measurements ($\rho = 0.307$), we obtain the average value $\tau_\pm = 12.347 \pm 0.030\ \text{ns}$, see \textsuperscript{12}.

To measure BR($K_{e3}^\pm$) and BR($K_{\mu3}^\pm$), we use both $K \rightarrow \mu\nu$ and $K \rightarrow \pi\pi^0$ decays as tags. We measure the semileptonic BRs separately for $K^+$ and $K^-$. Therefore, BR($K_{e3}$) and BR($K_{\mu3}$) are each determined from four independent measurements ($K^+$ and $K^-$ decays; $\mu$ and $\pi\pi^0$ tags). Two-body decays are removed by kinematics and the photons from the $\pi^0$ are reconstructed to reconstruct the $K^\pm$ decay point. From the TOF and momentum measurement for the lepton tracks, we obtain the $m_{\ell^2}$ distribution shown in figure 2. Further details are given in \textsuperscript{13}. Using the above result for $\tau_\pm$ to estimate the fiducial volume acceptance, we obtain BR($K_{e3}$) = $0.04972 \pm 0.00053$ and BR($K_{\mu3}$) = $0.03273 \pm 0.00039$, which we use in our evaluation of $|V_{us}|$.

We have also obtained a preliminary result on the BR($K^+ \rightarrow \pi^+\pi^0$), which is crucial to perform the fit of all $K^\pm$ BRs and for the $|V_{us}|$ determination of several experiments (NA48, ISTRA+, E865) in the normalization of the BRs ($K_{e3}^\pm$). About 800000 $K^+ \rightarrow \pi^+\pi^0$ have been select with kinematic cuts. Our preliminary result, BR($K^+ \rightarrow \pi^+\pi^0$) = $(20.658 \pm 0.065 \pm 0.090)\%$, is lower than the PDG value \textsuperscript{14} of about 1.3\%. Further details can be found in ref. \textsuperscript{15}.

![Figure 2: Distribution of $m_{\ell^2}$, from TOF information, for $K_{e3}^\pm$ events.](image-url)
5 $|f_+(0)V_{us}|$ and lepton universality

Using the BR($K_{f3}^{0,\pm}$), $\tau(K_L)$, $\tau(K^\pm)$ and the FFs from the KLOE results and $\tau(K_S)$ from the PDG, the values of $|f_+(0)V_{us}|$ has been evaluated for $K_{L\pi3}$, $K_{L\mu3}$, $K_{S\pi3}$, $K_{e3}$ and $K_{\mu3}$ decay modes. The inputs from theory, according to eq. 2, are the SU(2)-breaking correction evaluated with ChPT to $O(p^4)$, as described in $^{16}$, the long distance EM corrections to the full inclusive decay rate evaluated with ChPT to $O(\epsilon^2p^2)$ $^{16}$ using low-energy constants from ref. $^{17}$.

The average on the five different determination obtained taking into account all correlations is: $|f_+(0)V_{us}| = 0.2157 \pm 0.0006$ with $\chi^2$/dof $= 7.0/4$.

Comparison of the values of $|f_+(0)V_{us}|$ for $K_{e3}$ and $K_{\mu3}$ modes provides a test of lepton universality. We calculate the following quantity

$$r_{\mu e} = \frac{|f_+(0)V_{us}|_{\mu3, exp}}{|f_+(0)V_{us}|_{e3, exp}} = \frac{\Gamma_{\mu3}}{\Gamma_{e3}} \frac{I_{e3} (1 + \delta_{K e})^2}{I_{\mu3} (1 + \delta_{K \mu})^2},$$

where $\delta_{K \ell}$ stands for $\delta_{K}^{SU(2)} + \delta_{K \ell}^{EM}$. In the SM $r_{\mu e} = 1$. Averaging between charged and neutral modes, we find $r_{\mu e} = 1.000 \pm 0.008$. The sensitivity of this result is competitive with that obtained for $\pi \to l\nu$ and $\tau \to l\nu$ decays $^{18,19}$ whose accuracy is $\sim 0.4\%$.

6 Test of CKM unitarity

To get the value of $|V_{us}|$ we have used the recent determination of $f_+(0) = 0.9644 \pm 0.0049$ from BRC and UKQCD Collaborations obtained from a lattice calculation with $2 + 1$ flavors of dynamical domain-wall fermions $^{22}$. Using their value for $f_+(0)$, our $K_{f3}$ results give $|V_{us}| = 0.2237 \pm 0.0013$. Additional information is provided by the determination of the ratio $|V_{us}/V_{ud}|$, using eq. 3. From our measurements of BR($K_{f2}$) and $\Gamma(\tau_{\mu, e})$ from ref. $^{14}$ and the recent lattice determination of $f_K/f_\pi$ from the HPQCD/UKQCD collaboration, $f_K/f_\pi = 1.189 \pm 0.007$ $^{21}$, we obtain $|V_{us}/V_{ud}|^2 = 0.0541 \pm 0.0007$. We perform a fit to the above ratio and our result $|V_{us}|^2 = 0.05002 \pm 0.00057$ together with the result $|V_{ud}|^2 = 0.9490 \pm 0.0005$ from superallowed $\beta$-decays $^{20}$. We find $1 - |V_{us}|^2 - |V_{ud}|^2 = 0.0004 \pm 0.0007$ ($\sim 0.6\sigma$) and confirm the unitarity of the CKM quark mixing matrix as applied to the first row. The result of the fit is shown in figure 3.

![Figure 3: KLOE results for $|V_{us}|^2$, $|V_{us}/V_{ud}|^2$ and $|V_{ud}|^2$ from $\beta$-decay measurements, shown as 2$\sigma$ wide grey bands. The ellipse is the 1 $\sigma$ contour from the fit. The unitarity constraint is illustrated by the dashed line.](image)

7 Bounds on new physics from $K_{f2}$ decays

The comparison between the values for $|V_{us}|$ obtained from helicity-suppressed $K_{f2}$ decays and helicity-allowed $K_{f3}$ decays allows to put bounds on new physics. We study the quantity $R_{f23}$ =
\[
\left| \frac{V_{us}(K_{\mu 2})}{V_{us}(K_{\ell 3})} \right| \times \frac{V_{ud}(0^+ \rightarrow 0^+)}{V_{ud}(\pi^0 \rightarrow \mu^+)} ,
\]
which is unity in the SM, but would be affected only in \( V_{us}(K_{\mu 2}) \)
by the presence of non-vanishing scalar or right-handed currents. A scalar current due to a
charged Higgs exchange is expected to lower the value of \( R_{\ell 23} \), which becomes (see 23):
\[
R_{\ell 23} = \left[ 1 - \frac{m_{K^+}^2}{m_{H^+}^2} \left( 1 - \frac{m_{u^+}^2}{m_{K^+}^2} \right) \tan^2 \beta \right] \text{tan} \theta \text{ with tan} \beta \text{ the ratio of the two Higgs vacuum expectation values in the MSSM and } \epsilon_0 \approx 0.01.\]
Using our result on \( K_{\mu 2} \) and \( K_{\ell 3} \) decays, the lattice determinations of \( f_+(0) \) and \( f_K/f_\pi \) and the value of \( |V_{ud}| \) discussed above, we obtain \( R_{\ell 23} = 1.008 \pm 0.008 \). Fig. 4 shows the region in the \( \{ m_{H^+}, \tan \beta \} \) plane excluded at 95% CL by our result for \( R_{\ell 23} \).

![Figure 4: Region in the \( m_{H^+}, \tan \beta \) plane excluded by our result for \( R_{\ell 23} \); the region excluded by measurements of \( \text{BR}(B \rightarrow \tau \nu) \) is also shown.](image)

The ratio \( R_K = \frac{BR(K_{\mu 2})}{BR(K_{\ell 2})} \) is extremely well known in the SM, being almost free on hadronic uncertainties. Since the electron channel is helicity suppressed \( R_K \) is sensitive to contributions from physics beyond the SM. Deviations up to few percent on \( R_K \) are expected in minimal supersymmetric extensions of the SM and should be dominated by lepton-flavour violating contributions with tauonic neutrinos emitted 25. KLOE has selected about 8000 \( K_{\mu 2} \) events on 1.7 \( \text{pb}^{-1} \) by performing a direct search without the tag of the other kaon. Background from \( K_{\mu 2} \) has been reduced by means of kinematic cuts and calorimeter particle identification. Our preliminary result, \( R_K = (2.55 \pm 0.05 \pm 0.5) \times 10^{-5} \), allows to put bounds on the charged Higgs mass and tan \( \beta \) for different slepton mass matrix off-diagonal elements \( \Delta_{1,3} \). An accuracy of \( \sim 1\% \) is expected increasing the data sample analyzed, the control sample and Monte Carlo statistics.

8 Test of quantum coherence, CPT and Lorentz symmetry with the neutral kaons

Test of quantum mechanics (QM) can be performed by studying the time evolution of the
quantum correlated \( K_S K_L \) system, in particular studying the interference pattern of the decay
\( K_L K_S \rightarrow \pi^+ \pi^- \pi^+ \pi^- \). The distribution of the difference decay times is given by:
\[
I(|\Delta t|) \propto e^{-|\Delta t|/\Gamma_L} + e^{-|\Delta t|/\Gamma_S} - 2\cos(\Delta m|\Delta t|)e^{-\Gamma_L/2|\Delta t|} .
\]
One of the most direct ways to search for deviations from QM is to introduce a decoherence parameter \( \zeta \), i.e. multiplying by a factor \( (1 - \zeta) \) the interference term in the last equation.
The definition of \( \zeta \) depends on the basis chosen for the initial state 27 \( |i \rangle \propto |K_S(+\bar{p})\rangle|K_L(-\bar{p})\rangle - |K_L(+\bar{p})\rangle|K_S(-\bar{p})\rangle \) or \( |i \rangle \propto |K^0(+\bar{p})\rangle|K^0(-\bar{p})\rangle - |K^0(+\bar{p})\rangle|K^0(-\bar{p})\rangle \).
The case \( \zeta = 1 \) (i.e. total decoherence) corresponds to the spontaneous factorization of states (known as Furry’s hypothesis)\(^{28}\). Selecting a pure sample of \( K_L K_S \rightarrow \pi^+ \pi^- \pi^+ \pi^- \) and fitting eq. 6 to data, KLOE has obtained the following preliminary result based on \( 1 \text{ fb}^{-1} \): 
\[
\zeta SL = 0.009 \pm 0.022_{\text{stat}} \text{ and } \zeta_{00} = (0.03 \pm 0.12_{\text{stat}}) \times 10^{-5}
\]
consistent with QM predictions.

In a quantum gravity framework, space-time fluctuations at the Planck scale \( (\sim 10^{-33} \text{ cm}) \), might induce a pure state to evolve into a mixed one\(^{29}\). This decoherence, in turn, necessarily implies CPT violation\(^{30}\). In this context the CPT operator may be “ill-defined” and CPT violation effects might also induce a breakdown of the correlation in the initial state\(^{31,32}\) which can be parametrized in general as:
\[
|i⟩ \propto |K_S(\mp p)|K_L(\mp p)⟩ - |K_L(\mp p)|K_S(\mp p)⟩ + ω(|K_S(\mp p)|K_S(\mp p)⟩ - |K_L(\mp p)|K_L(\mp p)⟩) \quad \text{where } ω \text{ is a complex parameter describing CPT violation.}
\]
Its order of magnitude might be at most \( |ω| \sim \sqrt{(M_K^2/M_{\text{Planck}})/ΔΓ} \sim 10^{-3} \), with \( ΔΓ = Γ_S - Γ_L \). KLOE has improved its limit on the \( ο \) parameter using about \( 1 \text{ fb}^{-1} \). The preliminary results, obtained by fitting the \( I(Δt; \pi^+ \pi^- \pi^+ \pi^-) \) distribution, are \( \text{Re}ω = (-2.5_{-3.1}^{+3.4}) \times 10^{-4} \) and \( \text{Im}ω = \left(-2.2_{-3.4}^{+4.1}\right) \times 10^{-4} \), consistent with quantum coherence and CPT symmetry. The accuracy reaches the interesting region of the Planck’s scale.

![Figure 5: Fit of the difference \( t_1 - t_2 \) of the decay times of \( K_S \rightarrow \pi^+ \pi^- \) and \( K_L \rightarrow \pi^+ \pi^- \), where \( t_1 \) is the time of the kaon having cos \( θ > 0 \), in the range \( 0 < t_{\text{sid}} < 4h \). The black points are the experimental data, the histogram is the fit results and the hatched area is the uncertainty arising from the efficiency, the resolution and the background evaluation.](image)

Another possibility for CPT violation is based on spontaneous breaking of Lorentz symmetry in the context of the Standard Model Extension (SME)\(^{33,34}\). In the SME CPT violation manifests to lowest order only in the \( δ \) parameter, describing CPT violation in the time evolution, which exhibits a dependence on the kaon 4-momentum:

\[
\delta(p, θ, t_{\text{sid}}) = \frac{1}{2π} \int_0^{2π} δ(\vec{p}, t_{\text{sid}})dΦ = \frac{i sinΦSW e^{iΦSW}γ}{Δm} (Δa_0 + βΔa_Z cos θ + βΔa_Y sin θ sin Ωt_{\text{sid}} + βΔa_X sin χ cos θ cos Ωt_{\text{sid}})
\]

(7)

after integration on \( Φ \), where \( ∆a_0, ∆a_Z, ∆a_Y, ∆a_X, Φ, χ, Ω \) are the conventional polar and azimuthal angles defined in the laboratory frame around the \( z \) axis. ∆\( a_\mu \) are four CPT and Lorentz symmetry violating coefficients for the two valence quarks, \( β \) is the kaon velocity, \( γ = 1/\sqrt{1 - β^2} \), \( φ_{SW} \) is the superweak angle, \( χ \) is the angle between the \( z \) laboratory axis and the Earth’s rotation axis and \( Ω \) is Earth’s sidereal frequency. The sidereal time \( (t_{\text{sid}}) \) dependence arises from the rotation.

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of the Earth. KLOE has measured the $\Delta a_{X,Y,Z}$ parameters by using the channel $K_S K_L \to \pi^+ \pi^- \pi^+ \pi^-$ and performing an analysis on the polar angle $\theta$ and the sidereal time $t_{sid}$. Fitting the distribution of the decay times difference $f(t_1 - t_2; \pi^+ \pi^- (\cos \theta_1 > 0) \pi^+ \pi^- (\cos \theta_2 < 0); t_{sid})$ we obtain the preliminary results based on $1 fb^{-1}$: $\Delta a_X = (-6.3 \pm 6.0) \times 10^{-18} \text{GeV}$, $\Delta a_Y = (-2.8 \pm 5.9) \times 10^{-18} \text{GeV}$ and $\Delta a_Z = (-2.4 \pm 9.7) \times 10^{-18} \text{GeV}$. The result of the fit is shown in fig. 5. A limit on the $\Delta a_0$ parameter has been obtained through the difference on the $K_S$ and $K_L$ semileptonic charge asymmetry integrated on $t_{sid}$ and on a symmetrical polar angle region. Our preliminary result is $\Delta a_0 = (0.4 \pm 1.8) \times 10^{-17} \text{GeV}$.

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MEASUREMENT OF DIRECT CP VIOLATION PARAMETER $Re(\epsilon'/\epsilon)$ IN THE NEUTRAL KAON SYSTEM

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The final measurement of the direct CP violation parameter $Re(\epsilon'/\epsilon)$ performed by the KTeV collaboration is presented. The new result, $Re(\epsilon'/\epsilon) = [19.2 \pm 1.1 _{\text{stat}} \pm 1.8 _{\text{syst}}]$, improves precision of the previous measurement$^1$ and is consistent with it. Along with the measurement of $Re(\epsilon'/\epsilon)$, new measurements of the $K_L - K_S$ mass difference, $\Delta m$, the $K_S$ lifetime, $\tau_S$, the phase $\phi_\epsilon = \arg(\epsilon)$ and the phase difference $\Delta \phi$ are performed. The data are consistent with CPT symmetry, the value of $Re(\epsilon'/\epsilon)$ is consistent with the NA48 result$^2$.

1 Introduction

Violation of CP symmetry in weak interactions was first discovered in 1964 when the decay $K_L \rightarrow \pi^+\pi^-$ was observed. It was realized in the following experiments that the main reason for the effect is a small difference between $K^0 \rightarrow \bar{K}^0$ and $\bar{K}^0 \rightarrow K^0$ transition rates, which is termed as indirect CP violation. CP can be also violated directly in a decay amplitude, a search for this process has been performed by experiments at CERN$^{3,2}$ and Fermilab$^{4,1}$. In this letter, the final measurement of direct CP violation by the KTeV experiment at Fermilab is reported.

Direct CP violation manifests itself as a difference in the level of CP violation for different decay modes. For neutral kaons, $K \rightarrow \pi^+\pi^-$ and $K \rightarrow \pi^0\pi^0$ decay amplitudes can be compared:

$$
\begin{align*}
\eta_{+-} &= \frac{A(K_L \rightarrow \pi^+\pi^-)}{A(K_S \rightarrow \pi^+\pi^-)} = \epsilon + \epsilon' \\
\eta_{00} &= \frac{A(K_L \rightarrow \pi^0\pi^0)}{A(K_S \rightarrow \pi^0\pi^0)} = \epsilon - 2\epsilon'.
\end{align*}
$$

Here $\epsilon$ quantifies common indirect CP violation while $\epsilon'$ parameterizes a difference between the two modes and thus is a direct CP violation parameter.
CPT invariance imposes additional constraints on the complex parameters $\epsilon$ and $\epsilon'$. In particular phase of $\epsilon$ must be equal to the "superweak" phase, $\phi_\epsilon = \phi_{SW} \equiv \arctan(2\Delta m/\Delta \Gamma)$, where $\Delta m \equiv m_L - m_S$ is the $K_L - K_S$ mass difference and $\Delta \Gamma \equiv \Gamma_S - \Gamma_L$ is the difference in the decay widths. CPT invariance together with measurements of the strong phase shifts also requires that $\phi_\epsilon \approx \phi_{\epsilon'}$. Therefore, $Re(\epsilon'/\epsilon)$ is a measure of direct CP violation while $Im(\epsilon'/\epsilon)$ is a measure of CPT violation. Experimentally, $Re(\epsilon'/\epsilon)$ is determined using double ratio of the decay rates:

$$\frac{\Gamma(K_L \rightarrow \pi^+\pi^-)/\Gamma(K_S \rightarrow \pi^+\pi^-)}{\Gamma(K_L \rightarrow \pi^0\pi^0)/\Gamma(K_S \rightarrow \pi^0\pi^0)} \approx 1 + 6Re(\epsilon'/\epsilon), \quad (2)$$

while $Im(\epsilon'/\epsilon)$ can be determined from the phase difference of the decay amplitudes:

$$\Delta \phi \equiv \phi_{00} - \phi_{+} \approx -3Im(\epsilon'/\epsilon). \quad (3)$$

Previous measurements of $Re(\epsilon'/\epsilon)$ have established that it has small non-zero value. This letter presents the final KTeV measurement of $Re(\epsilon'/\epsilon)$ which is based on complete data sample, including new 1999 data period that about doubles the statistics of the previous KTeV publication, and significantly improved experimental procedure.

2 KTeV Detector and Data Analysis

The KTeV apparatus (see Fig. 1) uses double beam technique to simultaneously collect the four decay modes $K_{L,S} \rightarrow \pi^+\pi^- (\pi^0\pi^0)$. The two neutral beams are formed from secondary particles produced by 800 GeV/c protons colliding on a beryllium oxide target using a system of collimators, absorbers and sweeping magnets. The neutral kaon decays are detected in $110-158$ m range from the production target (for the KTeV coordinate system this corresponds to a positive Z direction). The kaon energies used in this analysis are in $40-160$ GeV range. At
125 m from the production target one of the beams passes through a plastic regenerator which produces coherent mixture of $K_L$ and $K_S$ states, for $K \rightarrow \pi\pi$ decays the $K_S$ state dominates. The regenerator alternates between the two neutral beams during the periods with no proton collisions on target, at about once per minute rate, in order to reduce systematic differences between $K_L$ and $K_S$ decays. The kaon beam with the regenerator is termed in the following as the regenerator beam while the other beam is termed as the vacuum beam.

The charged decay products are detected in a drift chamber spectrometer. The spectrometer is equipped with two chambers before and two after an analyzing magnet. Each chamber measures charged particle tracks in horizontal and vertical views. The neutral decay products are measured in a CsI crystal calorimeter, located after the spectrometer at 186 m from the production target. The crystals of the calorimeter have transverse dimensions of 2.5 $\times$ 2.5 cm$^2$ for the central region surrounded by 5 $\times$ 5 cm crystals in the outer range, there are 3100 crystals in total.

An extensive veto system rejects background events coming from interactions in the regenerator, semileptonic and $K_L \rightarrow \pi^0\pi^0\pi^0$ decays. The background levels, which include non-$K \rightarrow \pi\pi$ decays as well as $K \rightarrow \pi\pi$ decays in which the kaon scatters in the regenerator, after all selection cuts do not exceed 0.1% for the $\pi^+\pi^-$ ("charged") and 1.2% for the $\pi^0\pi^0$ ("neutral") mode.

The reconstruction of $K \rightarrow \pi^+\pi^-$ mode starts from selecting events with two track measured in the spectrometer. Each track is matched to a cluster in CsI calorimeter and $E/p < 0.85$ is required to reject $K \rightarrow \pi^+\pi^-\nu\overline{\nu}$ events. No signal is allowed in the muon veto system, located behind the CsI calorimeter, to reject $K \rightarrow \pi^+\mu^+\nu$ events. A high efficiency of the muon system is ensured by imposing $p > 8$ GeV/c condition for momentum of each track. The invariant mass of the two tracks, assuming the tracks are charged pions, is selected in 488 MeV/c$^2 < m_{\pi^+\pi^-} < 508$ MeV/c$^2$ range. The transverse momentum squared of the kaon is required to be $p_T^2 < 250$ MeV/c$^2$ in order to reject events in which the kaon undergoes scattering in the regenerator or in an upstream collimator.

To measure $K \rightarrow \pi^0\pi^0$ decays four photon clusters of energy are detected in the CsI calorimeter. The clusters are paired together to reconstruct $\pi^0 \rightarrow \gamma\gamma$ decays. For each pairing the $Z$ coordinate of the decay point with respect to the calorimeter surface is calculated as $Z_{12} = r_{12}\sqrt{E_1E_2/m_{\pi^0}}$, where $E_{1,2}$ are the photon energies, $r_{12}$ is the distance between the photons and $m_{\pi^0}$ is the nominal $\pi^0$ mass. All six pairings are considered and the one which leads to the most consistent $Z_{12}$ determination is used. The decay $Z$ vertex position is estimated using an error weighted average of $Z_{12}$. The kaon transverse vertex position is reconstructed by using a center of energy of the clusters, it is required to be situated inside the beam profile in order to reduce scattering background. The kaon energy is measured as a sum of the cluster energies. A cut on total invariant mass is imposed 488 MeV/c$^2 < m_{\pi^0\pi^0} < 508$ MeV/c$^2$ which rejects $K \rightarrow \pi^0\pi^0\pi^0\pi^0$ events.

Distributions of the $Z$ coordinate of $K_S \rightarrow \pi\pi$ and $K_L \rightarrow \pi\pi$ decay vertices have very different shape because of the difference in the lifetimes. To take this into account, KTeV uses a detailed Monte Carlo simulation (MC). Quality of this simulation can be tested by comparing the $Z$ vertex distribution in the vacuum beam, see Fig 2. A linear slope in the ratio of the data to MC distributions can be directly translated into uncertainty of $Re(e'/e)$ using a difference of an average $Z$ position of the decay vertex for $K_S$ and $K_L$ decays. The systematic uncertainty is derived based on $K_L \rightarrow \pi^+\pi^-$ decays for the charged and $K_L \rightarrow \pi^0\pi^0\pi^0\pi^0$ decays for the neutral mode.

Compared to the previous KTeV publication\cite{1}, several significant improvements of the measurement procedure were introduced. These include improvements for the 1999 data taking (i.e. better duty cycle for the proton extraction and repaired electrons of CsI calorimeter), for the data analysis (i.e. better model for drift chamber resolution which lead to ~ 15% increase of $m_{\pi^+\pi^-}$ resolution), while the main improvements were made for the detector simulation. The

\cite{1}
updates in MC include new charged particle tracing in the detector, which were also used for the KTeV measurement of the parameter $V_{us}$ and better description of the photon showers, using a new GEANT-based shower library. The new simulation of the photon showers leads to significant reduction of the energy scale uncertainty, which is the main source of the error for $Re(\epsilon'/\epsilon)$, this error is reduced from $1.3 \times 10^{-4}$ to $0.65 \times 10^{-4}$.

### 3 Results

For the full combined dataset, the result of the analysis is

$$Re(\epsilon'/\epsilon) = [19.2 \pm 1.1_{\text{stat}} \pm 1.8_{\text{syst}}] \times 10^{-4} = [19.2 \pm 2.1] \times 10^{-4}. \quad (4)$$

The result is in a good agreement with the previous KTeV publication:\ $Re(\epsilon'/\epsilon) = [20.7 \pm 1.5_{\text{stat}} \pm 2.4_{\text{syst}}] \times 10^{-4}$. A comparison of the KTeV measurement with other experiments is presented in Fig. 3. A good agreement between different results is observed; the world average, $Re(\epsilon'/\epsilon) = [16.8 \pm 1.4] \times 10^{-4}$, corresponds to a measurement of the direct CP violation parameter with 8% precision.

Decays in the regenerator beam are sensitive to $K_L - K_S$ interference and thus allow to measure $\Delta m$, $\phi_e$ and $Im(\epsilon'/\epsilon)$. Measurements of $\Delta m$ and $\phi_e$ depend strongly on the properties of the kaon regeneration and transmission in the regenerator beam. The transmission in the regenerator beam has been re-measured using a high statistics sample of $K \rightarrow \pi^+\pi^-\pi^0$ events collected in 1999. A dedicated study of the screening corrections allowed to significantly reduce uncertainty arising from the kaon regeneration. As a result, the measurement of $\phi_e$ is significantly improved compared to previous KTeV publication providing a better CPT symmetry test. For
an analysis without CPT constraints, KTeV obtains:

\begin{align*}
\tau_S &= [89.589 \pm 0.070] \times 10^{-12} \text{ s}, \\
\Delta m &= [5279.7 \pm 19.5] \times 10^6 \text{ h/s}, \\
\phi_e &= [43.86 \pm 0.63]^{\circ}, \\
\text{Im}(\epsilon' / \epsilon) &= [-17.20 \pm 20.20] \times 10^{-4}. 
\end{align*}

The measured \( \text{Im}(\epsilon' / \epsilon) \) corresponds to \( \Delta \phi = [0.30 \pm 0.35]^{\circ} \). The data are consistent with CPT symmetry: \( \text{Im}(\epsilon' / \epsilon) \) and \( \delta \phi = \phi_e - \phi_{SW} = [0.40 \pm 0.56]^{\circ} \) are consistent with zero. Imposing the CPT conservation as an additional constraint allows to reduce uncertainties on \( \tau_S \) and \( \Delta m \). This is illustrated in Fig. 4 which shows correlations of \( \tau_S, \Delta m \) and \( \phi_e \) together with a band derived from \( \delta \phi = 0 \) condition. The resulting \( \tau_S \) and \( \Delta m \) are:

\begin{align*}
\tau_S &= [89.623 \pm 0.047] \times 10^{-12} \text{ s}, \\
\Delta m &= [5269.9 \pm 12.3] \times 10^6 \text{ h/s}. 
\end{align*}

Using these values KTeV determines \( \phi_{SW|_{opt}} = [43.419 \pm 0.058]^{\circ} \).

### 4 Conclusions

The final measurement of \( Re(\epsilon' / \epsilon) \) and other kaon system parameters by the KTeV collaboration based on complete dataset is presented. Increase of the data sample and improvements of the analysis techniques allow to reduce the total uncertainties compared to the previous publication\(^1\). The world measurements of \( Re(\epsilon' / \epsilon) \) are consistent with each other and establish firmly the presence of direct CP violation in the kaon decays. With improved precision, the data do not show any indication of CPT symmetry violation.
Figure 4: $\chi^2 = 1$ contours of total uncertainty for (a) $\Delta m$ and $\tau_S$, (b) $\phi_\epsilon$ and $\Delta m$, (c) $\tau_S$ and $\phi_\epsilon$. Dashed lines show $\phi_\epsilon = \phi_{SW}$ CPT constraint. Larger (smaller) ellipses correspond to results without (with) assumption of CPT symmetry conservation.

References

Status of the CKM matrix

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I review the status of Cabibbo-Kobayashi-Maskawa matrix within the Standard Model, with a focus on exclusive $b \to (d,s)\gamma$ transitions and on charm and strange physics.

In the Standard Model (SM), the weak charged-current transitions mix quarks of different generations, which is encoded in the unitary Cabibbo-Kobayashi-Maskawa (CKM) matrix. In the case of three generations of quarks, the physical content of this matrix reduces to four real parameters, among which one phase, the only source of CP violation in the Standard Model (the lepton sector can also exhibit similar sources of CP violation once masses, provided by New Physics (NP), are considered). One can define these four real parameters as:

\[
\begin{align*}
\lambda^2 &= \frac{|V_{us}|^2}{|V_{ud}|^2 + |V_{us}|^2} \\
A^2\lambda^4 &= \frac{|V_{cb}|^2}{|V_{ud}|^2 + |V_{us}|^2} \\
\bar{\rho} + i\bar{\eta} &= -\frac{V_{ub}V_{cb}^*}{V_{cd}V_{cb}^*}. 
\end{align*}
\]  

(1)

This parametrisation is exact, unitary to all orders in $\lambda$ and independent of phase conventions. A Wolfenstein-like parametrisation of the CKM matrix can be derived up to an arbitrary power in the Cabibbo angle $\lambda = \sin(\theta_C)$, using the unitarity of the matrix to determine all its elements. A challenge for both experimentalists and theorists consist in extracting information on the underlying mechanism of CP violation from the wealth of data currently available, in the presence of the strong interaction that binds quarks into hadrons. Does the above CKM mechanism describe accurately the data? If yes, what are the values of $\lambda$, $A$, $\bar{\rho}$ and $\bar{\eta}$? If no, what is (are) the source(s) of CP violation beyond the Standard Model?

The CKMfitter group follows this program within the Rfit frequentist approach\(^1\). The likelihood function $L$ is defined as the product $L(y_{\text{mod}}) = L_{\text{exp}}(x_{\text{exp}} - x_{\text{the}}(y_{\text{mod}})) \cdot L_{\text{the}}(y_{QCD})$ where $x_{\text{exp}}$ denote experimental measurements and $x_{\text{the}}$ the corresponding theoretical predictions. $x_{\text{the}}$ depends on $y_{\text{mod}}$ which are either free parameters of the theory (e.g., the CKM matrix parameters) or hadronic quantities (e.g., form factors, decay constants... denoted $y_{QCD}$). Each
eral slightly incompatible solutions, the frequentist statistical treatment treats all the solutions on the same footing, leading to a broadening of the confidence intervals for $\gamma$ (a Bayesian analysis would integrate over hadronic parameters, so that different incompatible solutions sharing the same value of $\gamma$ yield an increase degree of belief in this value, reducing the uncertainty in the posterior p.d.f of $\gamma$).

The outcome of the global fit is shown in Fig. 1 in the usual $(\bar{\rho}, \bar{\eta})$ plane

$$A = 0.795^{+0.025}_{-0.015}, \quad \lambda = 0.2252^{+0.0008}_{-0.0008}, \quad \bar{\rho} = 0.135^{+0.033}_{-0.016}, \quad \bar{\eta} = 0.345^{+0.015}_{-0.018}$$

but also in the $(\bar{\rho}_s, \bar{\eta}_s)$ plane defined as $\bar{\rho}_s + i \bar{\eta}_s = -(V_{us}V_{ub}^*)/(V_{cs}V_{cb}^*)$ and more suitable to discuss the CKM mechanism for the $B_s$ sector. The corresponding triangle $(V_{us}V_{ub}^*)/(V_{cs}V_{cb}^*) + 1 + (V_{ts}V_{tb}^*)/(V_{cs}V_{cb}^*) = 0$ is squashed, with 2 sides of $O(\lambda^0)$ and 1 side of $O(\lambda^2)$. $\beta_s = \arg[-V_{cs}V_{cb}^*/(V_{ts}V_{tb}^*)]$, the angle opposite to the small side, is related to $B_s$ mixing in the SM. The global fit yields a small and well-predicted value $\beta_s = -0.0183^{+0.0009}_{-0.0008}$ rad, with which recent flavour-tagged $B_s^0 \rightarrow J/\psi\phi$ analysis from CDF and D0 present some tension.

The two experiments used different assumptions for their analyses (strong phases, width of the $B_s$ meson) and obtained nontrivial likelihoods. It seems sensible to wait for a combined analysis within a common framework and for a larger data sample before claiming a hint of NP in the $B_s$ sector.

## 2 $B \rightarrow V\gamma$

It has been known for a long time that the loop processes $b \rightarrow (d, s)\gamma$ can give an access to $|V_{t(d,s)}|$ which complement $\Delta m_{d,s}$ in an interesting fashion: we can test penguin versus box diagrams, so that an inconsistency between the two determinations, and with the global fit, would teach us in which direction to look for NP. Inclusive $B \rightarrow X_s \gamma$ decays have been computed with a high accuracy, but one can also consider exclusive $B \rightarrow V\gamma$ decays. The first attempts to compute the corresponding amplitudes used a factorisation approach. It was in particular used to determine

$$R_{\rho/\omega} = \frac{\mathcal{B}(\rho^\pm\gamma) + \frac{\tau_{\rho\pm}}{\tau_{\rho0}} \mathcal{B}(\rho^0\gamma) + \mathcal{B}(\omega\gamma)}{\mathcal{B}(K^{\pm\gamma}) + \frac{\tau_{K^{\pm\gamma}}}{\tau_{K^{0\gamma}}} \mathcal{B}(K^{0\gamma})} = \left|\frac{V_{td}}{V_{ts}}\right|^2 \left(\frac{1 - m_{\rho}^2/m_B^2}{1 - m_{K^\pm}^2/m_B^2}\right)^3 \frac{1}{\xi^2} [1 + \Delta R]$$

where $\xi$ is a ratio of form factors and $\Delta R$ is a correction from hadronic physics estimated as $\Delta R = 0.1 \pm 0.1$. This important step led many questions open. What is the dependence of $\Delta R$ on $V_{td}/V_{ts}$?
on the CKM matrix elements? Can one estimate and exploit isospin breaking? How to estimate weak annihilation processes, which for \((\rho,\omega)\gamma\) occurs at tree level and can be large, despite a formal \(1/m_b\) suppression. A further step was proposed by estimating \(1/m_b\)-suppressed terms, missed in QCD factorisation or in SCET, through light-cone sum rules. For each final state, all contributions can be expressed as a factor to the leading amplitude, i.e., the magnetic operator \(Q_7 = (e/8\pi^2)m_b\bar{D}\sigma_{\mu\nu}(1 + \gamma_5)F_{\mu\nu}b\):

\[
\bar{A} \equiv \frac{G_F}{\sqrt{2}} \left( \lambda_6^D a_7^U(V) + \lambda_6^C a_7^C(V) \right) \langle V_\gamma|Q_7|\bar{B} \rangle \quad \lambda_6^D = V_{UD}^* V_{Ub},
\]

where \(D = d, s\) and the coefficient \(a_7^U(V) = a_7^{U,\text{QCDF}}(V) + a_7^{U,\text{ann}}(V) + a_7^{U,\text{soft}}(V)\) is the sum of three terms. QCDF denotes the result from QCD factorisation at leading-order in \(1/m_b\) and up to \(O(\alpha_s)\) corrections, whereas ann and soft correspond to weak-annihilation and soft-gluon contributions. The latter are \(1/m_b\)-suppressed contributions which can be computed within QCD factorisation, but can be estimated through light-cone sum rules.

In this approach, each decay is described individually and the short- and long-distance contributions of \(u\) and \(c\) internal loops can be identified (they are note combined in a single correction \(\Delta R\)). For the evaluation, we followed refs. 6,8 and for the expressions of \(a_7\)'s and hadronic inputs (form factors, distribution amplitudes), using leading-order Wilson coefficients and the HFAG averages for the branching ratios (in units of \(10^{-6}\)):

\[
K^{*-\gamma} : 40.3 \pm 2.6, \quad K^{*0}\gamma : 40.1 \pm 2.0, \quad \rho^+\gamma : 0.88^{+0.28}_{-0.26}, \quad \rho^0\gamma : 0.93^{+0.19}_{-0.18}, \quad \omega\gamma : 0.46^{+0.20}_{-0.17},
\]

Together with the Belle value \(B(B_s \to \phi\gamma) = (57^{+18+12}_{-15-11}) \cdot 10^{-6}\). Fig. 3 shows the improvement from the previous treatment. The constraint is not a perfectly circular ring, due to the (previously neglected) sensitivity to other CKM matrix elements in the decay amplitude. The constraints from \(B \to V\gamma\) and from neutral \(B\) meson mixing have been superimposed to illustrate the compatibility of the two determinations, and their complementarity (we compare box and penguin processes with different theory sources). The study of CP asymmetries should provide further information on the apex of the \(B\)-meson unitarity triangle.

Figure 3: Constraints on the unitarity triangle from \(B \to V\gamma\) using the simplified expression of \(R_{\rho/\omega}\) (left) and considering the available branching ratios and their theoretical expressions with \(1/m_b\) corrections (right).
3 Lighter quarks and the lattice

The above constraints derived from $b$ transitions can be translated into values of CKM matrix elements involving lighter quark and they can be compared to direct measurements which have recently improved. Indeed, some lattice simulations with three dynamical light quarks (unquenched) are available with astoundingly small systematics, thus reducing QCD uncertainties.

As a first example, $|V_{ud}|$ has benefited from an improved analysis of super-allowed $\beta$ decays of nuclei, whereas $|V_{us}|$ has a shrinking uncertainty due to recent experimental results on $K_{\ell 3}$ and an improved lattice estimate of the relevant form factor $f_+(0) = 0.964(5)$ (domain-wall fermions, UKQCD+RBC)\textsuperscript{10,11}. Both values are used in the global fit, but an interesting cross-check consists in comparing the value of $|V_{us}/V_{ud}|$ from the fit with the value obtained by combining the measured ratio of leptonic decays $K \to \ell\nu/\pi \to \ell\nu$ with the lattice ratio of decay constant $f_K/f_\pi = 1.189(7)$ (staggered fermions, HPQCD+UKQCD)\textsuperscript{12}. The agreement shown on the left of Fig. 4 is remarkable, $f_K/f_\pi$ being notoriously very difficult to compute on the lattice (it involves only light quarks and the chiral extrapolation can yield large uncertainties).

A second example is the charm sector, which has always been thought of as a favourite place to test lattice QCD, since $m_c$ is close to the typical hadronic scale of 1 GeV. Lattice computations of form factors and decay constants should pin down $|V_{cd}|$ and $|V_{cs}|$ to a high accuracy. We illustrate the current improvement in the field in Fig. 5. The constraints on the nucleon and the kaon provide only a mild constraint, since $|V_{ud}| \simeq |V_{cs}|$ and $|V_{cd}| \simeq |V_{us}|$ only at first non trivial order in $\lambda$ (one needs an input from another sector to fix higher orders). The $B$ sector alone constrains $|V_{cd}|$ and $|V_{cs}|$ tightly and the combination of all indirect constraints turns out to be very powerful. We have also represented the direct constraints for $|V_{cd}|$, from $\nu N$ scattering, and for $|V_{cs}|$ from charmed-tagged $W$ decays (left) and from CLEO-c results on $D \to K\ell\nu$ (right)\textsuperscript{13}. The distorted shape of these regions comes from $|V_{cd}|^2 + |V_{cs}|^2 \leq 1$.

These results for lighter quarks seem to confirm both the consistency of the CKM picture and the high accuracy advocated by lattice results. However, a recent result has shattered this beautiful convergence. Indeed, CLEO-c and Belle have both measured the leptonic decay $D_s \to \ell\nu$, whereas a related unquenched lattice result $f_{D_s} = 241 \pm 3$ MeV (staggered, HPQCD+UKQCD)\textsuperscript{12}. This yields to $|V_{cs}| = 1.076 \pm 0.041$ in flat disagreement with unitarity and with the fit value $|V_{cs}| = 0.97351^{+0.00020}_{-0.00022}$; as shown on Fig. 4 (right)\textsuperscript{14}. This result is quite unsettling since the $D_s$ involves only strange and charm valence quarks and should be an ideal place for lattice simulations, whereas NP is not supposed to play a major role for such mesons. Paradoxically, the much more complicated $f_K/f_\pi$ led to an impressive agreement of experiment...
and theory, while $f_{D_s}$ points towards either uncontrolled systematics in unquenched lattice simulations (due to dynamical quarks?), overlooked systematics in the experimental measurements (radiative corrections?), or NP $^{15}$. In any case, interesting news should come from this sector.

Acknowledgments

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References

2. V. Sordini, this conference.
Figure 1: Constraints on the unitarity triangles corresponding to the $B_d$ (left) and $B_s$ (right) mesons.

individual measurement entering $\mathcal{L}_{\text{exp}}$ is considered as Gaussian by default (in the case of a non-Gaussian experimental measurement, the exact description of the associated likelihood is directly used in the fit) and correlations, if known, are taken into account. The uncertainties on the theoretical parameters $y_{QCD}$ define the allowed range of values for each parameter: $\mathcal{L}_{\text{the}}(y_{QCD}(i))$ is one within the allowed range and zero outside. The fit is performed on all the parameters $y_{\text{mod}}$ by minimizing $\chi^2(y_{\text{mod}}) \equiv -2\ln(\mathcal{L}(y_{\text{mod}}))$. For metrology (assuming a good agreement between data and theory), one splits $y_{\text{mod}} = (a, \mu)$, where $a$ are the parameters of interest (e.g., $\bar{\rho}, \bar{\eta}$) and $\mu$ are the remaining parameters. The minimum value $\chi^2_{\text{min};\mu}(a)$ is computed for a set of fixed values $a$ while $\mu$ is allowed to vary. The Confidence Level represented on the plots is obtained from the $\Delta\chi^2 = \chi^2_{\text{min};\mu}(a) - \chi^2_{\text{min}}$.

1. **The global fit**

The global fit involves a large set of constraints. At the time of the conference, recent and significant changes occurred for $|V_{ud}|$ and $|V_{us}|$, which will be discussed below. In addition, BABAR and Belle have presented new determinations of $\gamma$, based on the interference between the colour-allowed $B^- \to D^0K^-$ and colour-suppressed $B^- \to \bar{D}^0K^-$ decays. The accuracy of the method is driven by the size of $r_B = |A_{\text{suppr}}|/|A_{\text{favours}}| \simeq |V_{ub}V_{cs}^*|/|V_{cb}V_{us}^*| \times O(1/N_c)$ typically of order 0.1-0.2, and the different methods try to improve on this ratio by different choices of $D$ decay channels (GLW: $D$ into CP eigenstates, ADS: $D^{(*)}$ into doubly Cabibbo-suppressed states, GGSZ: $D^{(*)}$ into 3-body state and Dalitz analysis). For the GGSZ analysis, BABAR and Belle have increased their statistics and BABAR includes neutral $D$ into $K_S^0K^+K^-$. There has also been a $DK$ update from BABAR for GLW, and a similar update from Belle for ADS$^2$.

Combining these results yields $\gamma = (72^{+34}_{-30})^0$ (68% CL) which shows a rather mild improvement at 2 and 3 $\sigma$ with respect to combinations showed at previous conferences. The various methods provide values for $\gamma$, but also for the hadronic quantities such as $r_B$ or the relative strong phase $\delta$ between the two amplitudes. The current values for these quantities are not completely consistent among the methods, as illustrated in Fig. 2: the methods yield similar ranges for $\gamma$ but rather different values of the hadronic parameters. In such a situation with sev-
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Electroweak Interactions and Unified Theories

IV - Neutrino Physics and Astrophysics
CURRENT NUMI/MINOS OSCILLATION RESULTS

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The MINOS experiment is now making precise measurements of the $\nu_\mu$ disappearance oscillations seen in atmospheric neutrinos, and will extend our reach towards the so far unseen $\theta_{13}$ by looking for $\nu_e$ appearance in the $\nu_\mu$ beam. It does so by using the intense, well-understood NuMI neutrino beam created at Fermilab and observing it 735km away at the Soudan Mine in Northeast Minnesota. Results from MINOS’ first two years of operations will be presented.

1 Introduction

Results from the Super-Kamiokande experiment used neutrinos produced by cosmic ray interactions with the upper atmosphere to show that muon neutrinos ($\nu_\mu$) of energies from a few hundred MeV through TeV oscillate to tau neutrinos ($\nu_\tau$) as they travel the tens to thousands of kilometers through the earth to the detector [1]. This implies that neutrinos have mass, a finding of fundamental importance to both particle physics and astrophysics. The K2K experiment used a beam of neutrinos shot across Japan to the Super-K detector to confirm this result in a controlled fashion [2]. The MINOS (Main Injector Neutrino Oscillation Search) experiment has unambiguously confirmed this result. MINOS will precisely measure the oscillation parameters using the intense, well-calibrated NuMI (Neutrinos at the Main Injector) beam of neutrinos generated at Fermilab. This neutrino beam was commissioned in early 2005 and is aimed toward the Soudan Underground Physics Laboratory in northeastern Minnesota. The neutrinos are observed by similar magnetized steel/scintillator calorimeters near their origin in Fermilab and after traveling 735 km to Soudan.

Differences in signals between the two detectors have already provided the best measurement yet of $\nu_\mu \leftrightarrow \nu_e$ flavor oscillations in a long-baseline accelerator experiment, using the first two years operation of the NuMI neutrino beam [3]. With more data, MINOS will reach its projected
sensitivity to this mixing, improved sensitivity to any sub-dominant $\nu_e$ modes (a probe of $\theta_{13}$) and high statistics neutrino cross section studies. This paper presents the current result on $\nu_\mu \leftrightarrow \nu_\tau$ oscillations, the first look at the spectrum of neutral current ("NC") events in the MINOS near detector, the methods which will be used to search for $\nu_e$ appearance, and new data-driven sensitivities to $\theta_{13}$.

1.1 The NuMI Beam

The NuMI neutrino beam [4] uses 120 GeV protons from the Main Injector synchrotron at Fermilab incident upon a graphite target. 90% of the primary protons interact over the two interaction-length long target, producing showers of $\pi$ and $K$ mesons. These showers are focused by a pair of parabolic aluminum “horns”, pulsed electromagnets carrying current sheaths which focus the mesons into a beam. This beam is sent down a 1 m radius, 675 m long decay pipe. While in this pipe the mesons have a chance to decay into muons and muon neutrinos, but few of the muons have enough time to further decay before they are absorbed at the end of the pipe, a decay which would produce electron and anti-muon neutrinos. The resulting neutrino beam is thus composed of approximately 92.9% $\nu_\mu$, 5.8% $\bar{\nu}_\mu$, 1.2% $\nu_e$ and 0.1% $\bar{\nu}_e$ for the low-energy ("LE") beam configuration.

The target and horns are movable with respect to each other, allowing different focusing optics. The result is a beam which is configurable in energy, as seen in Fig. 1. The LE configuration produces a spectral peak closest to the first oscillation minima, given the oscillation parameters measured by Super-K and the 735 km baseline to the far detector. Moving the target with respect to the horns produces the “pME” and “pHE” beams peaked at medium and higher energies. While not at ideal energies for the $\nu_\mu$ disappearance analysis, these beams are much more intense (~970 and 1340 neutrino events at the far detector per 10$^{20}$ protons on target, compared to ~390 for the LE beam) and provide extra handles when using the near detector data to model the beam’s properties. The MINOS near detector is only a km away from the target, so even the LE beam produces around 10$^7$ neutrino interactions per 10$^{20}$ protons on target, a very high statistics sample of this weakly interacting particle. The beam currently delivers 3.1×10$^{12}$ protons over a 12$\mu$s spill every 2.2 s for an average power of 270 kW. The NuMI beam has been operational since march of 2005, and to date (of this conference, March 2008) has delivered more than 4×10$^{20}$ protons on target.

1.2 The MINOS Detectors

The MINOS experiment observes the NuMI beam with two detectors, “near” and “far”. A third “calibration” detector was exposed to beams of protons, pions, electrons and muons from the CERN PS [6] to determine detector response. The near detector at Fermilab is used to characterize the neutrino beam with high statistics and is 1 km downstream from the NuMI target. The far detector is an additional 734 km downstream. This experiment compares the spectra of different types of neutrino interactions at these two detectors to test oscillation hypotheses.

All three MINOS detectors are steel-scintillator sampling calorimeters [5] made of alternate planes of 4.1×1 cm cross section plastic scintillator strips and 2.54 cm thick steel plates. The near and far detectors have magnetized steel planes. The calibration detector was not magnetized as the incoming particle momenta were known. The extruded polystyrene scintillator strips are read out with wavelength-shifting fibers and multi-anode photomultiplier tubes. The far detector is 705 m underground in Soudan, MN, in a disused iron mine currently operated as a State Park by the Minnesota Department of Natural Resources. The 5,400 metric ton far detector consists of 486 8 m-wide octagonal steel planes interleaved with planes of plastic scintillator strips.

The 282 plane, 980 metric ton MINOS near detector is located at the end of the NuMI beam facility at Fermilab in a 100 m deep underground cavern. While the NuMI beam has diverged to
Figure 1: The measured energy spectrum of neutrinos from the NuMI beam observed by the MINOS near detector (top) and the ratios of data and expectations (bottom). Data points are the black dots, the untuned MC predictions are the blue curves, and MC predictions after tuning on hadronic $x_F$ and $p_t$ simultaneously across many different beam configurations are the red curves.

a mile wide at Soudan, at the near detector it is mostly contained in a meter-wide area, allowing a smaller detector and a factor of $10^6$ higher neutrino rate.

The much smaller calibration detector was used to measure the detailed responses of the MINOS detectors in a charged-particle test beam. This 12 ton detector consisted of 60 planes of unmagnetized steel and scintillator, each $1 \times 1 \text{m}^2$ [6]. It measured the energy and topological responses expected in the the near and far detectors, including the different electronics used in both larger devices. The energy responses of the three MINOS detectors were normalized to each other by calibrating with cosmic-ray muons.

2 Data Analysis

MINOS beam-based data is analyzed using a “blind analysis”. This method avoids looking at the actual data containing the physics being studied until the very end, removing potential biases and increasing confidence in the final result. Monte Carlo (“MC”) predictions are tuned and verified using data not sensitive to the physics in question (e.g. near detector data which is at too short a baseline to have experienced oscillations), and analysis cuts and techniques developed solely using simulated data. Only after these techniques are optimized and set are the sensitive data (in this example, the far detector oscillated data) revealed. All three of the results discussed in this paper are blind analyses, and are at different stages in the process.

The first step, common to all beam-based analyses, is to understand the beam itself. A detailed MC tracks simulated particles through the proton-meson-neutrino chain described in Sec. 1.1, to create an expected neutrino spectrum at the near detector. This MC is developed and crosschecked with information from the NuMI beam monitoring system, including a hadron monitor in the absorber at the end of the decay pipe and three muon monitors further downstream. As can be seen in the blue curves in Fig. 1, this does a decent but not perfect job of predicting the observed neutrino spectra in the near detector. Further tuning is done by reweighting hadronic $x_F$ and $p_t$ in the MC simultaneously across seven different beam con-
figurations and comparing to real near detector data, as the hadronic models have the most theoretical uncertainty. Four additional beam configurations (with different horn focusing currents) beyond those shown are included in this fit, and the resulting tuned predictions are the red line in Fig. 1. With the MC truth information in hand, a far detector prediction can be made by applying changes due to mundane things like geometrical and kinematic factors, or more exciting things like neutrino oscillations.

With a beam MC prediction in hand, topological features in the near detector data can be examined. Fitters to find tracks and neutrino interaction vertices, shower-finding algorithms, and particle identification ("PID") routines can be developed, tested, and calibrated using near detector data, the beam MC, and cosmic ray data at both detectors. Once an analysis can correctly matches the real data and the MC data, efficiencies and purities of the resulting sample can be extracted from the MC truth information, systematic uncertainties estimated, and expected sensitivity curves to the final physics parameters calculated. Only at this point is the "box opened", the far detector data run through the analysis, and the hypotheses tested to see what Mother Nature is really doing..

2.1 Atmospheric sector neutrino oscillations

The main goal of the MINOS experiment is a precision measurement of the $\nu_\mu$ disappearance oscillations first observed in atmospheric neutrinos. In the Standard Model, neutrinos are assumed to be massless and direct neutrino mass measurements have been able to establish only upper limits to their masses. Quantum mechanics predicts that if neutrinos do indeed possess a non-zero mass, then although the neutrinos are created and interact via the weak force as flavor eigenstates (corresponding to the flavors of leptons: electrons, muons and taus - $\nu_e, \nu_\mu, \nu_\tau$) they propagate through space as mass eigenstates ($\nu_1, \nu_2, \nu_3$). The flavor eigenstates are simple superpositions of the mass eigenstates [7]. If the neutrinos have differing masses, then the flavor of the neutrino varies as these states drift into and out of phase with each other while propagating through space, thus "oscillating" in flavor. For the case of two-flavor oscillations (e.g. $\nu_\mu \leftrightarrow \nu_\tau$) the probability that a neutrino produced via the weak interaction in the muon flavor state has oscillated to, or will be detected as, the tau state by the time it interacts is:

$$P_{\nu_\mu \rightarrow \nu_\tau} = \sin^2 2\theta_{23} \sin^2 \left( \frac{\Delta m^2_{32} L}{4E_{\nu}} \right), \quad (1)$$

where the properties of nature being probed are the amplitude or mixing angle $\theta_{23}$ and $\Delta m^2_{32} = m_3^2 - m_2^2$. The observable quantities are the energy of the neutrino $E_\nu$ and the distance the neutrino has traveled, also called the "baseline" $L$. Observation of neutrino flavor oscillations which vary as $L/E$ implies that both the terms $\Delta m^2_{32}$ and $\sin^2 2\theta_{23}$ are non-zero, and that at least one of the participating neutrino flavors has mass.

The analysis techniques discussed above were applied to data from the start of the NuMI beam through March 2007, totaling $2.947 \times 10^{20}$ “LE” beam protons on target ("pot"). This includes the previously published [3] $1.27 \times 10^{20}$ pot, although the analysis has been improved for both old and new data. A 3% larger fiducial volume was used, the data reconstruction was improved and retained 4% more good neutrinos, and the PID algorithm was revamped to provide both better purity and efficiency. The resulting sample of 563 $\nu_\mu$ charged current ("CC") neutrino interactions is plotted as a function of reconstructed neutrino energy on the left of Fig. 2, and a ratio with expectations (right) shows an energy dependent deficit. Equation 1 was applied on a two-dimensional ($\Delta m^2, \sin^2 2\theta$) grid to the MC predictions, and a $\chi^2$ formed compared to the data. Estimated systematic errors are less than the current statistical errors and applied as penalty terms to the $\chi^2$. The best fit value for the oscillation
Figure 2: (Left) The observed $\nu_\mu$ energy spectrum seen in the MINOS beam at the far detector for an exposure of $2.947 \times 10^{20}$ pot. Black crosses are the data with statistical error bars, the black line the null hypothesis, the red line the expectations of the best fit oscillation scenario of $|\Delta m_{32}^2| = 2.38^{+0.20}_{-0.16} \times 10^{-3}$ eV$^2$, $\sin^2 2\theta_{23} = 1.00_{-0.08}$, and the blue line (barely visible in the first few energy bins) the expected NC contamination. (Right) The same quantities expressed as a ratio of observed over expected null hypothesis.

Figure 3: The allowed regions in the oscillation parameter space of Eq. 1, obtained by fitting reweighted MC predictions to the MINOS data in Fig. 2. MINOS results (red) at 68% and 90% c.l. are compared to Super-K results (green) [1,8] and K2K results (blue) [2] at 90% c.l.

parameters to the MINOS data are $|\Delta m_{32}^2| = 2.38^{+0.20}_{-0.16} \times 10^{-3}$ eV$^2$ and $\sin^2 2\theta_{23} = 1.00_{-0.08}$, and the resulting 68% and 90% confidence limit contours are shown in Fig. 3.

2.2 Neutral Current Interactions

The $\nu_\mu$ disappearance results discussed in the previous section (2.1) use topological information to form a PID to select a sample of CC $\nu_\mu$ neutrinos, on the assumption that the flavor they are disappearing to is $\nu_\tau$, an active flavor of neutrino, unobserved in MINOS since the bulk of the NuMI neutrino flux is at energies below $\tau$ production threshold. However, if the second flavor of neutrino is a non-standard model sterile neutrino (one which experiences no weak interactions), the disappearance signature could look the same with very different underlying physics.

NC neutrino interactions hold the key to separating these two scenarios in MINOS. Active neutrinos of any flavor can experience a NC $Z^0$ exchange with a nucleon in the detector and produce a diffuse electromagnetic shower from the resulting $\pi^0$ decay to $\gamma\gamma$. A hypothetical sterile neutrino would not, so if some fraction of the $\nu_\mu$ signal is changing to $\nu_s$, the NC spectrum
would be distorted and NC flux reduced. A simple set of cuts has been applied to the near
detector data to select a NC-rich sample of neutrino interactions for further study. Short
tracks (<20 planes) are selected, then events with either no track at all or no track beyond five planes
from the shower are chosen. The resulting spectrum of NC near detector neutrino interactions
is shown in the left of Fig. 4, with projected limits on $f_s$ (the fraction of disappearance to $\nu_s$
shown on the right if no NC disappearance is observed.

This analysis is in the middle of the “blind” analysis scheme discussed above. Having chosen
a set of topological cuts to select a NC signal, data and MC comparisons are being made with
near detector data to verify that they are well understood before looking at the far detector
data to see what the potentially oscillated signal might look like. This analysis is expected to
be complete the summer of 2008.

2.3 Sensitivity to $\nu_e$ appearance

A third possibility for which particle the $\nu_\mu$’s are disappearing to is $\nu_e$. We know that there
could be some natural mixing between all three active flavors of neutrinos, and the amplitude
of this is parametrized as $\theta_{13}$. The Chooz reactor experiment saw no evidence of the converse $\nu_e$
disappearance at short baselines to establish an upper limit on $\theta_{13}$ [9]. However, the presence
of $\nu_e$ in the MINOS far detector beyond the low (1.3%) level inherent to the NuMI beam
could provide evidence for a non-zero $\theta_{13}$ below the Chooz limit, if the background of hadronic
showers masquerading as electromagnetic showers can be overcome. The PID algorithm used
for $\nu_e$ selection uses a neural net technique to pick out $\nu_e$-induced showers. At the near detector
the baseline is far too short for $\nu_\mu$ to have oscillated to $\nu_e$, so any observed $\nu_e$ must either be
inherent in the beam or the mis-reconstructed hadronic showers in question. To improve the
MC estimations of what the levels of these backgrounds might really be in the MINOS detectors,
data-driven sensitivity studies have been performed by close examination of near detector data
tagged as $\nu_e$ events.

Two methods are used in these studies. The first is to take the well-understood class of
CC events and subtract out those parts of the event associated with the muon track, leaving
only any hadronic component near the interaction vertex caused by the nucleon’s share of the
interaction energy. These “Muon-Removed Charged Current” (MRCC) events which are mis-
classified as $\nu_\mu$ interactions are exactly the sort of electromagnetic-dominated hadronic showers
that form a large part of the background for a $\nu_e$ appearance search. There is a 20% discrepancy
between the data and the MC predictions in both the standard $\nu_e$ and MRCC samples with the
MC overestimating the background. Comparisons of standard data and MC shower topological
distributions disagree in the same way as does MRCC data with MRCC MC, confirming that hadronic shower modeling is a major component of the disagreement. The MRCC sample is thus used to make and ad-hoc correction to the model to NC events per bin, taking the beam $\nu_e$ from the well-understood beam MC.

A second method for estimating the $\nu_e$ background from hadronic showers uses comparisons between the neutrino beam produced when the focusing horn’s current is turned off and the standard LE beam. The actual composition of the selected $\nu_e$ events is quite different in the two cases, allowing for the algebraic deconvolution of the different background components by expressing the total number as a sum of the different parts in the case of each beam:

$$N_{on} = N_{NC} + N_{CC} + N_e$$
$$N_{off} = r_{NC}N_{NC} + r_{CC}N_{CC} + r_eN_e$$

where $N_{NC}$ and $N_{CC}$ are the numbers of background events present originating from CC or NC interactions, $N_e$ the inherent beam $\nu_e$ taken from the beam MC, and the $r$’s the ratios that hold the differences between the two equations, $r_{NC(CC,e)} = N_{off}^{NC(CC,e)} / N_{on}^{NC(CC,e)}$. The horn on/off ratios are extracted bin-by-bin in energy from the MC, are independent of hadronic modeling, and match well between data and MC. These fractions can then be applied to the data itself to extract the components of the background, indicating that there is 24% too much CC and 28% too much NC backgrounds in the MC. Checks with a third (pHE) beam produce similar results, and both are compatible with the corrections from the MRCC method outlined above.

These data-drive backgrounds can then be extrapolated to the far detector for use in establishing the sensitivity expected when using a $\nu_e$ appearance search to try and measure $\theta_{13}$. These sensitivities are presented in Fig. 5 for three different exposures, the current $3.25 \times 10^{20}$ pot as well as those expected for next two years. The systematic errors for the current background estimation are found to be 10%, and with more data and study it is projected to fall to 5% for future years. The unknown variable of CP-violating $\delta$ contributes to $\nu_e$ appearance through the matter effects on the beam between Fermilab and Soudan, so the y-axis of these plots shows the effect of this $\delta$. The actual sign of $\Delta m^2_{32}$ also enters in, making this analysis less sensitive for the “inverted” mass hierarchy. However, after two more years of exposure MINOS will be sensitive to $\theta_{13}$ below the Chooz limit for most combinations of $\delta$ and mass hierarchy. The next step in this blind analysis is to examine far detector data in “sidebands” that allow verification of techniques without being sensitive to actual $\nu_e$ appearance.

3 Summary

The MINOS long-baseline neutrino experiment has been receiving 735 km baseline neutrinos from the NuMI neutrino beam since early 2005. The primary experimental goal of a precision measurement of the $\nu_\mu \leftrightarrow \nu_\tau$ disappearance oscillation parameters has been achieved. With the first $2.5 \times 10^{20}$ protons on target, $|\Delta m^2_{32}| = 2.38^{+0.20}_{-0.16} \times 10^{-3}$ eV$^2$ and $\sin^2 2\theta_{23} = 1.00_{-0.08}$. This is about a quarter of the expected final exposure, which will allow fine distinction between alternative disappearance hypotheses such as decoherence and neutrino decay to be made in the future. The first measurement of the spectrum of neutral current neutrino interactions has been made in the high-statistics near detector data. When the blind analysis of the corresponding far detector is complete later this year, it will be sensitive to a sterile neutrino fraction $f_s \leq 0.5$ at 90% c.l. Again using the near detector data, a data-driven background estimate to the $\nu_e$ appearance analysis has been made. This yields a sensitivity estimates comparable to the Chooz limit for the currently available exposure of $3.25 \times 10^{20}$ protons on target, reaching several times lower than this limit as soon as next year.

The NuMI beam and the MINOS experiment are going strong, the data and beam are well understood, and quality results are being produced. The next year should see the completion
Figure 5: Projected 90% c.l. limits on $\theta_{13}$ as a function of CP-violating $\delta$ from the MINOS experiment in the absence of a $\nu_e$ appearance signal. These limits use data-driven background estimates from the near detector. The vertical dashed line is the Chooz limit [9]. The rightmost (red) curve is the limit using the current exposure in the case of an inverted mass hierarchy; the neighboring (blue) curve is the limit from the same exposure if nature has a normal neutrino mass hierarchy. The two curves on the left are the progressively more sensitive normal hierarchy limits for the increased NuMI beam exposure over the next two years. The corresponding inverted hierarchy curves for these two scenarios are not shown for ease of viewing, but improve over the current exposure limits in a corresponding manner to the normal curves.

of initial analyses on all major experimental goals and the continued refinement of the precision parameter measurement of neutrino oscillations in the atmospheric neutrino sector.

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References

FIRST NEUTRINO EVENTS IN THE OPERA EMULSION TARGET

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OPERA is a long baseline neutrino oscillation experiment designed to observe $\nu_\mu \to \nu_\tau$ oscillations by searching for the appearance of $\nu_\tau$'s in an almost pure $\nu_\mu$ beam. The beam is produced at CERN and sent towards the Gran Sasso INFN laboratories where the experiment is running. OPERA started its data taking in October 2007, when the first 38 neutrino interactions where successfully located and reconstructed. This paper reviews the status of the experiment discussing its physics potential and performances for neutrino oscillation studies.

1 Introduction

OPERA\textsuperscript{1} is a long baseline experiment at the Gran Sasso underground laboratories (LNGS) and is part of the CERN Neutrino to Gran Sasso (CNGS)\textsuperscript{3} project. The detector has been designed to observe the $\nu_\mu \to \nu_\tau$ oscillations in the parameter region indicated by Super-Kamiokande\textsuperscript{2} through direct observation of $\nu_\tau$ charged current interactions. The detector is based on a massive lead/nuclear emulsion target complemented by electronic detectors that allow the location of the event and drive the scanning of the emulsions. A magnetic spectrometer follows the instrumented target and measures charge and momentum of penetrating tracks.

The CNGS beam is designed to provide $45 \cdot 10^{18}$ proton-on-target/year (p.o.t./y) with a running time of 200 days per year. The first CNGS technical run occurred in August 2006 with a delivered luminosity of $0.76 \cdot 10^{18}$ p.o.t. At that time only the electronic detectors were installed and under commissioning.

The first physics run occurred in October 2007, when OPERA had 40\% of the target mass installed. Due to technical problems, only $0.79 \cdot 10^{18}$ p.o.t. were delivered. A new physics run is going to start in summer 2008 with a planned luminosity of $\sim 30 \cdot 10^{18}$ p.o.t.

2 The OPERA detector

OPERA is a large detector ($10 \text{ m} \times 10 \text{ m} \times 20 \text{ m}$) located in the underground experimental Hall C of LNGS. As shown in Figure 1, the detector is made of two identical super-modules, aligned along the CNGS beam direction, each one consisting of a target and a muon spectrometer. The target section combines passive elements, the lead-emulsion bricks, and electronic detectors. Each target section consists of a multi-layer array of 31 target walls followed by pairs of planes of plastic scintillator strips ('Target Tracker'). A magnetic spectrometer follows the instrumented target and measures charge and momentum of penetrating tracks.
2.1 The Emulsion Target

The development of automatized scanning systems during the last two decades has made possible the use of large nuclear emulsion detectors. Indeed, nuclear emulsion are still successfully used nowadays, especially in neutrino experiments. The realization of a new scanning system has been carried out by two different R&D programs in the Nagoya University (Japan) and in several European laboratories belonging to the OPERA collaboration. These scanning systems were designed to take into account the requests of high scanning speed (about 20 cm$^2$/h) while keeping the extremely good accuracy provided by nuclear emulsions. For the European system, the Berne group took in particular the responsibility to develop an automatic emulsion film changer and to implement an innovative technique of nuclear emulsion scanning with the use of dry lenses instead of oil immersion ones (as it was always in the past), in order to simplify the emulsion handling. About 40 automatic microscopes are installed in the various scanning laboratories of the OPERA experiment.

The total number of emulsion films in the OPERA detector will be about 9 millions, for an area of about 110000 m$^2$. These quantities are orders of magnitude larger than the ones used by previous experiments. That made necessary an industrial production of the emulsion films, performed by the Fuji Film company, in Japan, after an R&D program conducted jointly with the OPERA group of the Nagoya University.

The OPERA emulsions are made up of two emulsion layers 44 µm thick coated on both sides of a 205 µm triacetate base. The AgBr crystal diameter is rather uniform, around 0.2 µm, and the sensitivity is about 35 grains/100 µm for minimum ionizing particles.

The main constituent of the OPERA target is the brick. It is a pile of 57 nuclear emulsion sheets interleaved by 1 mm thick lead plates. The brick combines the high precision tracking capabilities provided by the emulsions with the large mass given by the lead. The OPERA brick is a detector itself. In addition to the vertex identification and τ decay detection, shower reconstruction and momentum measurements using the Multiple Coulomb Scattering can be performed, being the total brick thickness of 7.6 cm equivalent to 10 $X_0$. Bricks are hosted in the walls of the target.

The occurrence of a neutrino interaction inside the target is triggered by the electronic detectors. Muons are reconstructed in the spectrometers and all the charged particles in the target tracker. The brick finding algorithm indicates the brick where the interaction is supposed to be occurred. The trigger is confirmed in the Changeable Sheet Doublet (CSD), a pair of...
emulsion films hosted in a box placed outside the brick, as interface between the latter and the target tracker. Before detaching the CSD from the brick, they are exposed to an XRay spot, in order to define a common reference system for the two CS and the first emulsion in the brick (with a precision of a few tens of µm). Afterwards the CS are developed and the predictions from target tracker are searched for within a few cm area. If these are confirmed the brick is brought outside the Gran Sasso laboratory and exposed to cosmic ray before development.

The mechanical accuracy obtained during the brick piling is in the range of 50-100 µm. The reconstruction of cosmic rays passing through the whole brick allows to improve the definition of a global reference frame, leading to a precision of 1-2 µm.

All the tracks located in the CSD are subsequently followed inside the brick, starting from the most downstream film, until they stop. Then a general scanning around the stopping point(s) is performed, tracks and vertices are reconstructed, the primary vertex is located and the kinematic analysis defines the event topology.

2.2 The Target Tracker

The main role of the Target Tracker is to provide a trigger and identify the right bricks where the event vertex should be located. Each wall is composed by orthogonal planes of plastic scintillator strips (680 cm × 2.6 cm × 1 cm). The strips are made of extruded polystyrene with 2% p-terphenyl and 0.02% POPOP, coated with a thin diffusing white layer of TiO$_2$. Charged particle crossing the strips will create a blue scintillation light which is collected by wavelength-shifting fibers which propagate light at both extremities of the strip. All fibers are connected at both ends to multianode Hamamatsu PMTs. The detection efficiency of each plane is at 99%. A detailed description of the Target Tracker design can be found in $^{12}$

2.3 The Spectrometer

The spectrometer allows to suppress the background coming from charm production through the identification of wrong-charged muons and contributes to the kinematic reconstruction of the event performed in the target section. The magnet $^{13}$ is made of two vertical walls of rectangular cross section and of a top and bottom flux return path. The walls are built lining twelve iron layers (5 cm thickness) interleaved with 2 cm of air gap, allocated for the housing of the Inner Tracker detectors, Resistive Plate Chambers, RPCs. Each iron layer is made of seven slabs, with dimensions 50 × 1250 × 8200 mm$^3$, precisely milled along the two 1250 mm long sides connected to the return yokes to minimize the air gaps along the magnetic circuit. The slabs are bolted together to increase the compactness and the mechanical stability of the magnet which acts as a base for the emulsion target support. The nuts holding the bolts serve as spacers between two slabs and fix the 20 mm air gap where the RPCs are mounted.

The precision tracker is made of drift tubes planes located in front, behind and between the two magnet walls: in total 12 drift tube planes covering an area of 8 m × 8 m. The tubes are 8 m long and have an outer diameter of 38 mm. The trackers allows to reconstruct the muon momentum with a resolution $\Delta p/p \leq 0.25$. A particle entering the spectrometer is measured by layers of vertical drift tube planes located before and after the magnet walls. Left-right ambiguities are resolved by the two dimensional measurement of the spectrometer RPCs and by two additional RPC planes, equipped with pickup strips inclined of ±42.6° with respect to the horizon (XPC). The Inner Tracker RPCs, eleven planes per spectrometer arm, give a coarse measurement of the tracks and perform pattern recognition and track matching between the precision trackers. The OPERA RPCs$^{14}$ are “standard” bakelite RPCs, similar to those used in the LHC experiments: two electrodes, made of 2 mm plastic laminate (HPL) are kept 2 mm apart by means of polycarbonate spacers in a 10 cm lattice configuration. The double coordinate readout is performed by means of copper strip panels. The strip pitch is 3.5 cm for the horizontal
strips and 2.6 cm for the vertical layers. The OPERA RPCs have a rectangular shape, covering an area of about 3.2 m². The sensitive area between the iron slabs (8.75 × 8 m²), is covered by twenty one RPCs arranged on seven rows, each with three RPCs in a line. In total, 1008 RPCs have been installed in the two spectrometers.

3 Physics performances

The OPERA detector will host 155000 bricks for a total target mass of 1350 tons. The signal of the occurrence of νμ → ντ oscillation is the charged current interaction of the ντ’s inside the detector target (ντN → τ−X). The reaction is identified by the detection of the τ lepton in the final state through the decay topology and its decay modes into an electron, a muon, and a single or three charged hadrons:

\[ \tau^- \rightarrow e^- \nu_\tau \nu_e \]
\[ \tau^- \rightarrow \mu^- \nu_\tau \nu_\mu \]
\[ \tau^- \rightarrow (h^- h^+) h^- \nu_\tau (n \pi^0) \]

The branching ratio for the electronic, muonic and hadronic channel are 17.8%, 17.7% and 64.7% respectively. For the typical τ energies expected with the CNGS spectrum the average decay length is \( \sim 450 \mu m \).

Neutrino interactions will occur predominantly inside lead plates. Once the τ lepton is produced, it will decay either within the same plate, or further downstream. In the first case, τ decays are detected by measuring the impact parameter of the daughter track with respect to the tracks originating from the primary vertex, while in the second case the kink angle between the charged decay daughter and the parent direction is evaluated.

The τ search sensitivity, calculated for 5 years of data taking with a total number of \( 45 \times 10^{18} \) integrated p.o.t. per year, is given in table 1.

<table>
<thead>
<tr>
<th>τ decay channels</th>
<th>Signal ( \div \Delta m^2 \text{ (Full mixing)} )</th>
<th>Background</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \tau \rightarrow \mu^- )</td>
<td>2.9 × ( 10^{-3} \text{ (eV}^2 )</td>
<td>0.17</td>
</tr>
<tr>
<td>( \tau \rightarrow e^- )</td>
<td>3.5 × ( 10^{-3} \text{ (eV}^2 )</td>
<td>0.17</td>
</tr>
<tr>
<td>( \tau \rightarrow h^- )</td>
<td>3.1 × ( 10^{-3} \text{ (eV}^2 )</td>
<td>0.24</td>
</tr>
<tr>
<td>( \tau \rightarrow 3h )</td>
<td>0.9 × ( 10^{-3} \text{ (eV}^2 )</td>
<td>0.17</td>
</tr>
<tr>
<td>ALL</td>
<td>10.4 × 15.0</td>
<td>0.76</td>
</tr>
</tbody>
</table>

The main background sources are given by large angle scattering of muons produced in νμCC interactions, secondary hadronic interaction of daughter particles produced at primary νμ interaction vertex and decays of charmed particles produced at primary νμ interaction vertex. Comparing the total number of detected ντ interaction with the estimated background it’s clearly seen that OPERA is quite a background-free experiment. In Figure 2 the ντ observation probability at 3 and 4 σ as a function of \( \Delta m^2 \) is reported.

4 Results from the first runs

The first CNGS run was held in August 2006\textsuperscript{15}. At that time only electronic detectors were installed: the brick filling started indeed at the beginning of 2007. From 18 to 30 August 2006 a total intensity of \( 0.76 \times 10^{18} \) p.o.t. was integrated and 319 neutrino-induced events were collected (interactions in the rock surrounding the detector, in the spectrometers and in the target walls). Thanks to this first technical run the detector geometry was fixed and the full reconstruction of
Figure 2: 3 and 4 $\sigma$ observation probability as a function of $\Delta m^2$.

electronic detectors data tested. It was also possible to fine-tune the synchronization between CERN and Gran Sasso, performed using GPS clocks. Furthermore, the zenith-angle distribution from penetrating muon tracks was reconstructed and the measured mean angle of $3.4 \pm 0.3^\circ$ was well in agreement with the value of $3.3^\circ$ expected for CNGS neutrinos traveling from CERN to the LNGS underground laboratories.

The first OPERA physical run was held in October 2007. At that time about 40% of the target was installed, for a total mass of about 550 tons. In about 4 days of continuous data taking $0.79 \times 10^{18}$ p.o.t. were produced at CERN and 38 neutrino interactions in the OPERA target were triggered by the electronic detectors. The corresponding bricks indicated by the brick finding algorithm were extracted and developed after the cosmic ray exposure and their emulsions sent to the scanning laboratories. In a few hours the first neutrino interactions of the OPERA experiment were successfully located and reconstructed. In Figure 3 the display of two events is shown. The left one is a $\nu_\mu$CC interaction with 5 prongs and a shower reconstructed pointing to the primary interaction vertex ($\gamma$ conversion after a $\pi^0$ decay). In the second a quite energetic shower (about 4.7 GeV) coming from the primary interaction vertex is visible.

This first physical run was quite short but very significative. Indeed it allowed a full testing of the electronic detectors and the data acquisition. Furthermore, the brick finding algorithm was successfully used to locate the bricks were the neutrino interaction occurred. Finally, the target tracker to brick matching was proved to be able to satisfy the expectations and the full scanning strategy validated.

5 Outlook and future plans

The OPERA target will by completed by May 2008. In June a first 150-day period of CNGS beam at nominal intensity is expected to start. About $30 \times 10^{18}$ p.o.t. will be integrated, equivalent to about 3500 neutrino interactions. More then 100 charm decays will be collected, so that the capability to reconstruct $\tau$ decays will be fully exploited. The corresponding number of expected triggered $\nu_\tau$ is 1.3: with some luck the first $\nu_\tau$ candidate event will be observed during the 2008 OPERA run.
Figure 3: Two reconstructed neutrino interaction from the OPERA 2007 run. The event displayed on the left is a $\nu_\mu$CC interaction. The right side shows an event where an energetic shower comes from the interaction vertex.

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NEUTRINO OSCILLATIONS, LOW-ENERGY EXCESS, NUMI NEUTRINOS, AND ANTINEUTRINOS IN MINIBOONE

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The MiniBooNE collaboration published initial results from a search for $\nu_\mu$ to $\nu_e$ oscillations in June 2007. While no evidence of $\nu_e$ appearance was found in the region of neutrino energies expected from LSND under a standard $2\nu$ mixing hypothesis, a significant excess was observed at lower energies. The results from the oscillation analysis, the excess at low energies, an analysis of neutrino events from the NuMI beamline in the MiniBooNE detector, and the current running in $\bar{\nu}$-mode are discussed.

1 Introduction

The primary motivation of the MiniBooNE experiment at Fermi National Laboratory is to search for the oscillation of neutrinos from $\nu_\mu$ into $\nu_e$ flavors over a short baseline consistent with the mixing parameters found by the LSND experiment\(^1\). To achieve the required neutrino flux, MiniBooNE extracts 8.89 GeV/c protons directly from the Booster and impacts them on a Be-target at the center of a focusing horn pulsed at 174 kA. The horn was initially operated in a polarity such that positively charged pions were focused to create a $\nu_\mu$ beam. Ideally the content of the secondary neutrino beam at the detector would be purely $\nu_\mu$, however kaons produced at the target as well as muons from the pion decay chain result in an intrinsic $\nu_e$ contamination that is 0.5% of the $\nu_\mu$ flux, and intrinsic $\bar{\nu}_e$ contamination of only 0.05%. At the event level, intrinsic $\nu_e$ are indistinguishable from oscillation $\nu_e$, but have an energy distribution that extends to higher energies than that expected for $\nu_e$ appearing due to oscillation. The wrong-sign background, i.e. $\bar{\nu}_\mu$ in a $\nu_\mu$ beam or vice versa, forms 6% of the flux at the detector. The $\nu$-mode fluxes at the detector for all neutrino species are plotted as a function of energy in Figure 1(a). In this mode of operation, the experiment has collected a total of $6.6 \times 10^{20}$ protons on target (POT), although results in this document and elsewhere have only been presented for $5.7 \times 10^{20}$ POT.

By switching the polarity of the horn, negatively charged pions were focused to create a predominantly $\bar{\nu}_\mu$ neutrino beam. The $K^- \pi^+$ production at the target in $\bar{\nu}_\mu$ mode is suppressed relative to the corresponding $K^+ \pi^-$ in $\nu$-mode, resulting in a reduced intrinsic $\bar{\nu}_e$ contamination of 0.4%. However, at 0.2% the wrong-sign intrinsic $\nu_e$ contamination in $\bar{\nu}$-mode is higher than the wrong-sign $\bar{\nu}_e$ component was in $\nu$-mode. Overall the total fraction of intrinsic $\nu_e + \bar{\nu}_e$ contamination remains the same. The primary difference in $\bar{\nu}$-mode relative to $\nu$-mode is that leading particle effects cause the $\pi^-$ to have softer momenta and higher production angles, thus making it more difficult for the horn to efficiently focus the negative pions. The overall neutrino
flux at the detector is reduced by approximately a factor of two in \( \nu \)-mode. In addition to this reduction, the cross-section \( \bar{\nu} \) charge-current interactions is smaller than \( \nu \) resulting in total factor of 4 reduction in the overall rate of neutrino interactions per POT. Another consequence of running in \( \bar{\nu} \)-mode is that the harder spectrum of the \( \pi^0 \) make them more difficult to defocus. As shown in Figure 1(b), this results in a much larger wrong-sign component where \( \nu_\mu \) now form 16\% of the beam in \( \bar{\nu} \)-mode. As of this writing, the experiment has collected nearly \( 3.0 \times 10^{20} \) POT with an eventual goal of reaching a total of \( 5.0 \times 10^{20} \) POT over the next year. A complete discussion of the MiniBooNE beamline and the flux prediction has been recently summarized\(^3\) and submitted to PRD.

The average energy of the neutrino beam is 800 MeV, which requires placing the detector at a distance of 0.5 km in order to preserve the \( L/E \) of LSND. The detector consists of 800 tons of pure mineral oil contained in a 12 m diameter sphere. The sphere is divided into an inner and outer region via an optical barrier, with the outer region serving as an active veto lined with 240 8 in photomultiplier tubes (PMTs). The inner region contains 1280 inward-facing PMTs. Although the optical barrier is at a radius of 540 cm, most analyses require events to be contained and the reconstructed vertices to fall within a sphere of radius \( \approx 500 \) cm. A more complete description of the detector has been recently submitted to NIM and can be found on the archive\(^4\).

The fundamental strength of the detector lies in the ability to separate \( \nu_e \)-CCQE signal events with an \( e \) in the final state from \( \nu_\mu \)-CCQE events that produce a \( \mu \). It is also necessary to resolve when two electromagnetic particles are created in order to separate NC-\( \pi^0 \) interactions where the pion immediately decays, \( \pi^0 \rightarrow \gamma \gamma \). Since pure mineral oil is used, the detector is mainly a Cerenkov-type detector, where particle ID is performed through analysis of the multiplicity and topology of rings projected on the outer wall of the detector. However, despite being very pure, high quality Marcol 7 mineral oil, impurities result in some scintillation light. A Michel \( e \) from \( \nu_\mu \) decay will typically have 75\% of its light contained in the prompt Cerenkov ring, while another 25\% is emitted isotropically with a dominant time constant of 34 ns. The upside is that it means the detector can also be used to detect pure NC events where no particles are relativistic enough to emit a Cerenkov cone. The downside is that the optical model is a very complex system where the modeling of the light production and propagation must account for Cerenkov flux, scintillation flux, absorption and readmission of light (fluorescence), Raman and Rayleigh scattering, reflections, along with the PMT response. A satisfactory set of parameters for the optical model, several of which are wavelength dependent, were obtained through a combination of benchtop tests and \textit{in situ} calibrations.
2 The Oscillation Result and the Low Energy Excess

The basics of the oscillation analysis have been described in the original oscillation result paper\(^2\). The steps involve calibrating the raw PMT data, passing both data and Monte Carlo through event reconstruction, developing a robust particle ID for extracting the $\nu_e$-CCQE candidates, applying constraints arising from the well-measured $\nu_\mu$-CCQE and NC-$\pi^0$ samples\(^5,6\), and then performing a $\chi^2$ minimization under a full systematic error covariance matrix.

The analysis was divided into two quasi-independent analyses, referred to as the boosted-decision tree (BDT) and the track-based likelihood (TBL). Both analyses relied on the same underlying Monte Carlo samples to form their background prediction. As such they share identical sources of systematic errors stemming from underlying uncertainties in the flux prediction, understanding of cross-sections in the 1 GeV range, and optical modeling. They diverged starting at the reconstruction stage where the TBL analysis used a more sophisticated (and consequently CPU intensive) set of algorithms. The TBL analysis then constructed a set of maximum likelihoods under the fit hypotheses that the final state consisted of a single $e$, $\nu_\mu$, or $\pi^0$. In addition to the basic pre-cuts shared by both analyses, the TBL analysis formed particle ID cuts using these likelihoods and the reconstructed pion mass. By comparison the BDT analysis used a much faster reconstruction algorithm, constructed a large sample of macroscopic variables, and then input $\approx 170$ of these quantities into a boosted-decision tree to form a single variable as its PID cut. When tested on the Monte Carlo, the BDT analysis had a larger signal-to-background ratio than the TBL. However, when a systematic error analysis was performed by running 1000s of underlying Monte Carlo worlds with changes in the underlying model parameters, the response of the BDT analysis showed a higher sensitivity to systematic errors, which marginally outweighed any gains in signal-to-background. In a predefined procedure, the analysis with the better ultimate sensitivity, TBL, was quoted as the final experimental result, with the BDT serving as a very powerful cross-check. Neither analysis saw an excess of $\nu_e$ events in the region of reconstructed $E_\nu$ above 475 MeV where and LSND-like would have appeared under a 2$\nu$ mixing hypothesis. The resulting limit curves for both analyses are shown in Figure 2(b), with the TBL $\nu_e$ spectrum shown in Figure 2(a). The analyses were comparable enough in sensitivity that after the unblinding, a downward fluctuation in the BDT data set resulted in the limit curve actually being a little better than the TBL analysis at low $\Delta m^2$.

Below the oscillation analysis region, a 3.4 $\sigma$ excess of $96 \pm 28$ events was observed in the 300-475 MeV region. Many consistency checks have been performed on the events in that region.
to verify that they do not exhibit undue pathologies. The $x$, $y$, $z$, and $r$ distributions are consistent with neutrino interactions spatially distributed throughout the the detector, the events are uniformly distributed in time throughout the duration of the run, and perhaps most importantly, a visual inspection of the event displays for events in the low energy region confirms all of the events exhibit the characteristics expected of a single ring from an electromagnetic shower. It should be noted that ring originating from a single $e$ or a single $\gamma$ are indistinguishable. Since the initial publication of the low energy result, the existing analysis was pushed down in threshold to include a bin in the 200-300 MeV range. The excess in this region is similar in significance at $91 \pm 31$ events, although the highly correlated systematic error at low energies precludes a simple quadrature sum of the overall significance.

In addition to performing consistency checks and extending the analysis to lower energies, a full reevaluation of all of the backgrounds and associated systematic errors was undertaken. The results are expected to be presented in a paper to be published later in the year and will include:

- Photonuclear disintegration which was absent from GEANT3, and can create a background if one of the $\gamma$s in a NC-$\pi^0$ event is lost.
- More comprehensive hadronic errors, particularly in the final states after a photonuclear event occurs.
- Better handling of beam $\pi^+$ uncertainties with errors propagated directly from meson production errors.
- Improved measurement of neutrino-induced $\pi^0$ with finer momentum binning.
- Internal MiniBooNE measurement of the coherent/resonant $\pi^0$ fraction.
- Refined cuts that efficiently remove backgrounds coming from $\nu$ interactions in the dirt surrounding the detector.
- Extra $\nu$-mode data acquired since the initial publication, 16% increase in statistics.
- A comprehensive review of how the radiative $\Delta$ decay rate is extracted from the measured $\pi^0$ rate.

![Preliminary Data](image1)

![Preliminary Data](image2)

![Preliminary Data](image3)

Figure 3: Preliminary data to Monte Carlo comparisons for events in the MiniBooNE detector coming from the NuMI beamline. Cuts have been applied to isolate (a) $\nu_\mu$-CCQE, (b) NC-$\pi^0$, and (c) $\nu_e$-CCQE like samples.

## 3 Events in MiniBooNE from the NuMI Beamline

The MiniBooNE detector is positioned 110 mrad off-axis from the NuMI beamline that delivers neutrinos to the MINOS experiment. By comparison, the off-axis angle in T2K is 35 mrad and NOVA is somewhat smaller at 14.5 mrad, however, even at 110 mrad MiniBooNE still sees a
significant flux of neutrinos from NuMI. Since the NuMI and BooNE neutrinos are produced in independent spills, there is no confusion about which beam the neutrino events originate from and a dedicated NuMI trigger can easily be established. At the large off-axis angle, neutrinos in MiniBooNE from the NuMI beamline have a similar average energy as the normal Booster neutrino beam. The distance from the NuMI target to the MB detector is a little larger, so the overall $L/E$ is about 1.4 times larger. The largest difference in the neutrino flux coming from the NuMI beam is in the amount of intrinsic $\nu_e$ contamination, which at 5% is an order of magnitude higher than in the Booster neutrino beam.

A track-based analysis with very similar cuts to the official TBL analysis has been performed on the NuMI events. The preliminary results are shown in Figure 3. The agreement in the $\nu_\mu$-CCQE and NC-$\pi^0$ samples is very good, which validates both the flux prediction from the NuMI beamline and the response of the MiniBooNE detector. The $\nu_e$-CCQE sample shows an excess in the region from 200-900 MeV. However, the large correlated systematic error results in a statistical significance of only 1.3σ.

![Figure 4: (a) World data for $\nu$ cross-sections in charged-current processes, and (b) the projected sensitivity to $\nu_e$ appearance for $5 \times 10^{20}$ POT delivered in $\nu$-mode.](image)

### 4 Running in $\nu$-mode

Currently MiniBooNE is running in $\nu$-mode in order to check for $\nu_\mu$ oscillating into $\nu_e$. However, due to the considerations discussed in Section 1, the lower neutrino rate will reduce the sensitivity relative to the same POT delivered in $\nu$-mode, see Figure 4(b). Preliminary projections for sensitivities with errors ranging from statistics-only to a nearly full systematic error treatment are shown. Since the LSND experiment used a $\nu_\mu$ beam, this will provide a more direct but less sensitive check of the LSND signal.

In addition to $\nu$ oscillations, current knowledge of antineutrino cross-sections at energies below 1 GeV is limited, see Figure 4(a). MiniBooNE will be able to provide measurements of various exclusive NC and CC channels. Absolute measurements will be limited by a flux uncertainty that is currently 15%, however, measuring $\nu/\overline{\nu}$ ratios should have smaller errors and could prove to be very useful for experiments like T2K where the beam will also be operated in both modes with neutrinos of a similar energy.

Finally, valuable information about the low energy excess will be revealed since several of the potential explanations, including anomaly-mediated photons or new gauge bosons presented...
elsewhere in these proceedings, make distinct predictions for how the excess should extrapolate to a $\nu_\mu$ beam.

References

SOME NEW IMPLICATIONS OF THE ANOMALOUS BARYON CURRENT
IN THE STANDARD MODEL

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Phenomenological implications of the anomalous baryon current in the Standard Model are discussed, in particular neutrino-photon interactions at finite baryon density. A pedagogical derivation of the baryon current anomaly is given.

1 Introduction

The baryon current in the Standard Model is not conserved in the presence of electroweak gauge fields. Although classically we have

$$\partial_\mu J^\mu = \partial_\mu \left( \frac{1}{3} \sum_q q \gamma^\mu \bar{q} q \right) = 0,$$

the baryon current divergence acquires quantum corrections when gauge fields are coupled differently to left- and right-handed quarks. For the Standard Model electroweak gauge fields, we have

$$\partial_\mu J^\mu = -\frac{1}{64\pi^2} \varepsilon^{\mu\nu\rho\sigma} \left( g_2^2 F_{\mu\nu}^a F_{\rho\sigma}^a - g_1^2 F_{\mu\nu}^Y F_{Y\rho\sigma}^Y \right) \neq 0,$$

where $F_{\mu\nu}^a = \partial_\mu W_\mu^a - \partial_\nu W_\mu^a + g_2 \varepsilon^{abc} W_\mu^b W_\nu^c$ is the covariant $SU(2)_L$ field strength and $F_{\mu\nu}^Y$ is the weak hypercharge field strength. This curious fact may have profound cosmological implications through the generation of baryon number at the electroweak phase transition.

As discussed in Refs. 3,4 and reviewed in this talk, nonconservation of baryon number is connected to novel effects that can be observed in laboratory experiments, and that may have interesting astrophysical implications. This report begins with a theoretical review by analyzing the baryon number anomaly in analogy to the perhaps more familiar axial anomaly. Turning to
phenomenology, some observable consequences in neutrino scattering experiments are described, and several other directions to explore are mentioned.

# 2 Theoretical excursion

## 2.1 The axial current anomaly and $\pi^0 \to \gamma \gamma$

A famous implication of gauge anomalies is the necessity for a nonzero $\pi^0 \to \gamma \gamma$ amplitude due to the nonconservation of the iso-triplet axial-vector quark current,

$$J_5^\mu = \frac{1}{2} (\bar{u} \gamma^\mu \gamma_5 u - \bar{d} \gamma^\mu \gamma_5 d).$$

In the presence of electromagnetism we have,

$$\partial_\mu J_5^\mu = \frac{e^2}{32\pi^2} \epsilon^{\mu\nu\rho\sigma} F_{\mu\nu} F_{\rho\sigma}. \tag{4}$$

If low-energy QCD is described by a theory of mesons, a nonzero $\pi^0 \to \gamma \gamma$ amplitude is necessary in order to reproduce this result. Let us recall how this works explicitly, by considering the object:

$$\int d^4 x e^{-iq \cdot x} \langle \gamma(p)\gamma(k)| J_5^\mu(x)|0\rangle \equiv \left[ B_5(q, \mu) \right] \times \left[ \epsilon^*_\nu(p) e^*_\rho(k)(2\pi)^4 \delta^4(p + k - q) \right], \tag{5}$$

first at the quark level, and then at the meson level. The field $B_5$ denotes a background field coupled to $J_5^\mu$, and $A$ is the photon. At the quark level, after a proper definition of the relevant triangle diagram that ensures vector current conservation, a standard calculation shows that in lowest order perturbation theory,

$$i q_\mu \left[ B_5 \right] = \frac{e^2}{4\pi^2} \epsilon^{\nu\rho\alpha\beta} p_\alpha k_\beta, \tag{6}$$

consistent with (4). How is this result reproduced in terms of the low-energy effective action where the quarks are replaced by mesons? First, there is no gauge invariant operator connecting $B_5$ and two photons directly, so that

$$B_5 = 0. \tag{7}$$

A nonzero contribution is however obtained from the pion pole (consider the limit of vanishing quark masses),

$$B_5 = -iC_1 q^\mu \times \frac{i}{q^2} \times (-iC_2) \frac{e^2}{4\pi^2} \epsilon^{\nu\rho\alpha\beta} p_\alpha k_\beta. \tag{8}$$
Here $C_1$ denotes the strength of the $\pi$ coupling to the axial current, and $C_2$ is the strength of the pion-photon vertex. From the chiral lagrangian with Wess-Zumino-Witten term\textsuperscript{9,10}, we necessarily have $C_1 = f_\pi, C_2 = 1/f_\pi$. Contracting (8) with $iq_\mu$ reproduces (6) and hence (4). Phrased differently, if low-energy QCD is described by an effective theory of pions, then the process $\pi^0 \to \gamma\gamma$ occurs with a fixed strength.

2.2 The baryon current anomaly

The anomalous baryon current can be treated in close analogy to the anomalous axial-vector current above. We must however pay close attention to which currents are conserved, since in the present case it is no longer true that “vector currents are conserved, axial-vector currents are anomalous,” as the usual intuition suggests. Suppose that we introduce a background field $B_\mu$ coupled to baryon number. Then the baryon current is defined by varying the action with respect to $B_\mu$:

$$J^\mu = \frac{\delta S}{\delta B_\mu},$$

and its divergence is read off from

$$\delta S = - \int d^4x \epsilon(x) \partial_\mu J^\mu,$$

where $\delta B_\mu = \partial_\mu \epsilon$. Thus the problem of calculating the anomalous divergence of the baryon current is reduced to the introduction of $B_\mu$.

However, we must not be too naive in introducing $B_\mu$; otherwise we may start with a gauge invariant theory, but end up with a non-gauge-invariant (i.e., nonsensical) theory. Varying the nonsensical theory would not give the correct symmetry current and its divergence. To be explicit, let us return to the example of the axial-vector current for a single fermion, and suppose that we add the perturbation

$$\bar{\psi}(i\slashed{\partial} + A)\psi \to \bar{\psi}(i\slashed{\partial} + A + B_5 \gamma_5)\psi.$$  

Then the theory naively remains invariant under electromagnetic gauge transformations,

$$\psi \to e^{i\epsilon} \psi, \quad A_\mu \to A_\mu + \partial_\mu \epsilon, \quad B_5 \to B_5 \epsilon.$$  

However, due to the effects of anomalies\textsuperscript{a} the theory is in fact not gauge invariant. For a sensible theory, we must add at the same time as the perturbation (11), a counterterm:

$$\bar{\psi}(i\slashed{\partial} + A)\psi \to \bar{\psi}(i\slashed{\partial} + A + B_5 \gamma_5)\psi + \mathcal{L}_{ct}(A, B_5),$$  

where explicitly,

$$\mathcal{L}_{ct}(A, B_5) = \frac{1}{6\pi^2} \epsilon^{\mu\nu\rho\sigma} B_\mu A_\nu A_\rho A_\sigma.$$  

The results (4) and (6) have an implicit dependence on the choice of counterterm.\textsuperscript{b} In particular, the “Bardeen”\textsuperscript{8} form of the counterterm, of which (14) is an example, is employed to conserve vector currents in the presence of arbitrary background fields.

When nonvectorlike currents are gauged, a different counterterm must be employed. For the general case the explicit counterterm is given in Ref.\textsuperscript{4}. Let us consider the baryon current for a single standard model generation, and for simplicity restrict attention to the neutral gauge bosons $A$ and $Z$. The Bardeen counterterm is then

$$\mathcal{L}_{Bardeen} = \frac{eg_2}{24\pi^2 \cos \theta_W} \epsilon^{\mu\nu\rho\sigma} (B_\mu Z_\nu \partial_\rho A_\sigma + A_\mu Z_\nu \partial_\rho B_\sigma),$$

\textsuperscript{a}That is, due to the effects of the fermion measure, in path integral language.

\textsuperscript{b}The dependence can be made explicit by performing the calculation with Weyl fermions.
whereas the full counterterm is

$$\mathcal{L}_{ct} = \frac{e g_2}{24 \pi^2 c_W} \epsilon^\mu\nu\rho\sigma (-2 B_\mu Z_\nu \partial_\rho A_\sigma + A_\mu Z_\nu \partial_\rho B_\sigma). \quad (16)$$

If we now write

$$S = [S + S_{Bardeen} - S_{ct}] + S_{ct} - S_{Bardeen}, \quad (17)$$

then the variation (10) vanishes for the bracketed combination in (17), and from the remainder we can read off immediately using (15) and (16):

$$\partial_\mu J^\mu = - \frac{e g_2}{8 \pi^2 \cos \theta_W} \epsilon^\mu\nu\rho\sigma \partial_\mu A_\nu \partial_\rho Z_\sigma, \quad (18)$$

yielding the result (2).

In the language of chiral lagrangians, the new counterterm has the novel property that it leaves residual “pseudo-Chern Simons” terms in the action, i.e., terms involving the epsilon tensor, but having no pion fields. Such terms are subtracted if the Bardeen counterterm is used instead, since it can be shown that $\mathcal{L}_{Bardeen} = -\mathcal{L}(\pi = 0)$. Equivalently, the counterterm requires a different boundary condition for “integrating the anomaly” to obtain the anomalous part of the chiral lagrangian$^9$; this is again related to the statement that $\mathcal{L}(\pi = 0) \neq 0$.

### 2.3 Vector mesons

The preceding discussion shows how to incorporate spin-1 background fields into the chiral lagrangian without upsetting gauge invariance in the fundamental gauge fields. In particular, the $SU(2)_L \times U(1)_Y$ gauge anomaly cancellation between quark and lepton sectors is not upset when background fields are coupled to the quark flavor symmetries. With these background field probes in place, it is straightforward to derive properly defined (covariant) currents and the associated anomalous divergences, by an appropriate variation of the action.

The relevance of the background field discussion for vector mesons is twofold. First, since physical spin-1 mesons such as $\rho$ and $\omega$ behave mathematically like these background fields, we have found the “slots” which these fields fit into when constructing our chiral lagrangian. Second, and relatedly, once we know that the vector mesons inhabit these slots, we find new physical effects related to the quark level anomalies, e.g. to the baryon current anomaly. These effects can be observed experimentally. For example, at the level of vector meson dominance, new effects will be described by the interaction

$$\mathcal{L} = \frac{e g_2}{8 \pi^2 \cos \theta_W} \epsilon^\mu\nu\rho\sigma \omega_\mu Z_\nu \partial_\rho A_\sigma. \quad (19)$$

This is in the same spirit as using $\pi^0 \to \gamma \gamma$ as a probe of the axial current anomaly.

The theoretical description can be refined at low energy by integrating out the vector mesons; the vector dominance assumption then translates into a prediction for the coefficients of certain $1/m_\omega^2$ suppressed operators.

### 3 Phenomenology

To see that the vector mesons are indeed described as part of the WZW term structure, we can verify that the same coupling strength is observed in accessible decay modes, such as $\omega \to 3\pi$ and $\omega \to \pi \gamma$. As depicted in the Fig. 1, these are all parts of the baryon current, expressed in terms of the fields, including nuclear sources, in the low-energy chiral lagrangian:

$$J^\mu = \bar{N} \gamma^\mu N + \frac{1}{4 \pi^2} e \epsilon^{\mu\nu\rho\sigma} \left( -\frac{2i}{f_\pi} \partial_\nu \pi^+ + \partial_\rho \pi^- \partial_\sigma \pi^0 - \frac{e}{f_\pi} \partial_\nu \pi^0 \partial_\rho A_\sigma + \frac{e g_2}{2 \cos \theta_W} Z_\nu \partial_\rho A_\sigma + \ldots \right). \quad (20)$$
Figure 1: Different parts of the baryon current. The bottom leg denotes a field such as $\omega$ coupling to this current, and the blob denotes a source of baryon number, such as a nucleus.

For example,

$$\Gamma(\omega \to \pi\gamma) \approx \frac{3\alpha E^3}{64\pi^4 f^2_\pi} \left(\frac{2}{3} g_\omega\right)^2 \approx 0.76 \text{ MeV} \left(\frac{2}{3} g_\omega}{6}\right)^2.$$  \hfill (21)

Similarly, $\frac{2}{3} g_\omega \approx 6$ is obtained for $\omega \to 3\pi$, including the $\omega \to \rho\pi \to 3\pi$ contributions. A consistent, although somewhat uncertain, value of the $\omega$ coupling to the baryon current is also obtained for the first diagram in Fig. 1, using one-meson exchange models of the force between nucleons, and isolating the isoscalar $J^P = 1^-\pi$ channel\textsuperscript{12}. The effective coupling is expected to be somewhat larger in this case, since “$\omega$” is actually representing a tower of resonances.

Figure 2: Analogy to the Primakoff effect: on the left, one of the photons in the $\pi\gamma\gamma$ vertex couples to electric charge; on the right, $\omega$ from the $\omega Z\gamma$ vertex couples to baryon number.

We wish to access the final diagram in Fig. 1, i.e., the pure gauge field part of the baryon current, that is most directly related to the baryon current anomaly. We expect $\omega$ to couple to this part of the current with the same strength as the other parts. Now, if the $Z$ mass were small,\textsuperscript{d} the Standard Model would predict a decay mode,

$$\Gamma(\omega \to Z\gamma) = \frac{3\alpha}{256\pi^4} \frac{E^3}{m_Z^2} \frac{m_Z^2}{\cos^2 \theta_W} \left(\frac{2}{3} g_\omega\right)^2 \left(1 + \frac{m_Z^2}{m_\omega^2}\right).$$ \hfill (22)

Of course, the decay $\omega \to Z\gamma$ is not physically allowed. Nevertheless, processes involving virtual $Z^*$ are allowed, and can lead to interesting effects. Since there will be a weak suppression, we should focus on situations in which the $Z$ is “useful”, e.g. processes involving neutrinos or parity violation. We can also make the $\omega$ “useful”, e.g., by utilizing its strong coupling to baryon number to look for enhanced rates when scattering off nuclei, rather than searching for the tiny branching fraction $\omega \to \gamma\nu\bar{\nu}$. As depicted in Fig 2, this is in analogy to probing the $\pi^0\gamma\gamma$ coupling via the Primakoff effect, where one of the photons couples coherently to the electric charge of the nucleus.

### 3.1 Neutrino scattering

The interaction (19) will induce neutrino-photon interactions in the presence of baryon number. For example, single photons are produced in neutrino-nucleus scattering, as depicted in Fig. 3.

\textsuperscript{a}Note that $g_\omega = \frac{2}{3} g'$ in the conventions of Ref.\textsuperscript{3,4}.

\textsuperscript{d}Consider the limit $m_{u,d} \to 0$, $g_2 \to 0$ with $v$ fixed, so that $m_u^2 \ll m_Z^2 \ll m_\omega^2$. Then the $Z$ will eat mostly Higgs field, and effects of $\pi - Z$ mixing can be ignored.
In the approximation where the nuclear interactions are described by one-meson exchange, there will be competing contributions from virtual $\pi^0$ and $\rho^0$ exchange. However, $\pi^0$ exchange is suppressed by the accidental smallness of $1 - 4\sin^2 \theta_W$ in the Standard Model, and the $\rho^0$ exchange diagram is suppressed in amplitude by $\sim (1 + 1 - 1)^2/(1 + 1 + 1)^2 = 1/9$ relative to $\omega$, due to the fact that $\omega$ is isoscalar, whereas $\omega$ is isotriplet; this can be thought of as a coherence effect at the nucleon level. Further enhancement of the $\omega$ exchange due to coherence over adjacent nucleons can occur in kinematics where small enough momentum is exchanged with the nucleus.

Neglecting effects such as coherence, form factors and recoil, the cross section for the process depicted in Fig. 3 for scattering off an isolated nucleon is

$$\sigma \approx \frac{\alpha g_\omega^4 G_F^2}{480\pi^6 m_\omega^4} E_\nu^6 \approx 2.2 \times 10^{-41} (E_\nu/1\text{GeV})^6 (g_\omega/10)^4 \text{cm}^2. \quad (23)$$

The photon energy distribution in this approximation is

$$\frac{d\sigma}{dE_\gamma} \propto E_\gamma^3 (E_\nu - E_\gamma)^2, \quad (24)$$

and the angular distributions is flat,

$$\frac{d\sigma}{d\cos \theta_\gamma} \propto \text{constant}. \quad (25)$$

Form factors will suppress the cross section at large momentum exchange, pulling the angular distribution forward. As an illustration, the photon energy and angular distribution for a 700 MeV neutrino incident on an isolated nucleon, including nuclear recoil and the form factor.
induced by $\omega(780)$ exchange, is depicted in Fig. 4. A more detailed analysis of single photon events will be presented elsewhere\(^5\).

In the absence of large coherent enhancements, e.g. for scattering on small nuclei, we should ideally use relatively large incident neutrino energies, in order to overcome the mass of $\omega$. Also, if it is not possible to distinguish photon showers from electron showers, a pure $\nu_\mu$ beam should be used in order to avoid a background from charged-current scatters, $\nu_e + n \rightarrow e^- + p$. In fact, these requirements have overlap with experiments looking for $\nu_e$ appearance in a $\nu_\mu$ beam. For example, MiniBooNE\(^1\)\(^3\) and (in the future) T2K\(^1\)\(^4\) have $\nu_\mu$ beams with energy spectra of order several hundreds of MeV, but primarily $\lesssim 1$ GeV, largely within the range of a chiral lagrangian description. Single photons that are mistaken for electrons are a background to $\nu_e$ appearance searches, as depicted in Fig. 5. It is interesting that an excess of events observed by MiniBooNE is in the same order of magnitude as predicted by (23), and has similar characteristics to the distributions in Fig. 4. Experiments with higher energy neutrinos are also of interest, but pass beyond a simple chiral lagrangian description.

3.2 Neutrino pair production

Similar interactions can give rise to photon conversion into neutrino pairs in the presence of baryon number, as depicted in Fig. 6. A nonnegligible contribution to neutron star cooling via this mechanism was computed in Ref.\(^3\). Similar effects will occur in the hot and dense environment of a supernova core.

3.3 Parity violation

Besides neutrino interactions, we can use the $Z$ to mediate parity violation. The interaction (19) will give rise to potentially interesting effects in various parity-violating observables. These will be investigated elsewhere\(^5\).

4 Summary

This report began with a pedagogical derivation of the baryon current anomaly in the Standard Model. The counterterm structure in this derivation is interesting because it requires residual “pseudo-Chern-Simons” terms in the action when background vector fields are coupled to the
quark flavor symmetries. This exercise is significant for phenomenology because the same framework can be used to describe vector meson interactions in vector dominance approximation. The resulting extension of the QCD chiral lagrangian provides a useful guide to new effects, such as “baryon-catalyzed” neutrino-photon interactions and parity violation. Other applications of the formalism that have not been discussed here include a description of “natural parity violating” QCD vector meson decays, such as $f_1 \to \rho \gamma$. It is also interesting to relate this framework to five-dimensional descriptions of QCD, both as a means of constraining “AdS-QCD” models, and potentially using such models to predict undetermined constants appearing in the chiral lagrangian.

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I give a brief overview over various attempts to reconcile the LSND evidence for oscillations with all other global neutrino data, including the results from MiniBooNE. I discuss the status of oscillation schemes with one or more sterile neutrinos and comment on various exotic proposals.

1 Introduction

Reconciling the LSND evidence\(^1\) for $\bar{\nu}_\mu \rightarrow \bar{\nu}_e$ oscillations with the global neutrino data reporting evidence and bounds on oscillations remains a long-standing problem for neutrino phenomenology. Recently the MiniBooNE experiment\(^2,3\) added more information to this question. This experiment searches for $\nu_\mu \rightarrow \nu_e$ appearance with a very similar $L/E_\nu$ range as LSND. No evidence for flavour transitions is found in the energy range where a signal from LSND oscillations is expected ($E > 475$ MeV), whereas an event excess is observed below 475 MeV at a significance of $3\sigma$. Two-flavour oscillations cannot account for such an excess and currently the origin of this excess is under investigation\(^2\), see also\(^4\). MiniBooNE results are inconsistent with a two-neutrino oscillation interpretation of LSND at 98% CL\(^3\), see also\(^5\). The exclusion contour from MiniBooNE is shown in Fig. 1 (left) in comparison to the LSND allowed region and the previous bound from the KARMEN experiment\(^6\), all in the framework of 2-flavour oscillations.

2 Sterile neutrino oscillations

The standard “solution” to the LSND problem is to introduce one or more sterile neutrinos at the eV scale in order to provide the required mass-squared difference to accommodate the LSND signal in addition to “solar” and “atmospheric” oscillations. However, in such schemes there is severe tension between the LSND signal and short-baseline disappearance experiments, most importantly Bugey\(^7\) and CDHS\(^8\), with some contribution also from atmospheric neutrino data\(^9\).
I report here the results from a global analysis including MiniBooNE data within schemes with one, two and three sterile neutrinos.\cite{10}

Four-neutrino oscillations within so-called (3+1) schemes have been only marginally allowed before the recent MiniBooNE results,\cite{11,12,13} and become even more disfavored with the new data. We find that the LSND signal is disfavored by all other null-result short-baseline appearance and disappearance experiments (including MiniBooNE) at the level of 4\sigma.\cite{10} The corresponding upper bound on the effective LSND mixing angle is shown in Fig. 1 (right). Five-neutrino oscillations in (3+2) schemes allow for the possibility of CP violation in short-baseline oscillations. Using the fact that in LSND the signal is in anti-neutrinos, whereas present MiniBooNE data is based on neutrinos, these two experiments become fully compatible in (3+2) schemes.\cite{10} Moreover, in principle there is enough freedom to obtain the low energy excess in MiniBooNE and being consistent at the same time with the null-result in the high energy part as well as with the LSND signal, see Fig. 2 (left, red histogram). However, in the global analysis the tension between appearance and disappearance experiments remains unexplained. This problem is illustrated in Fig. 2 (right) where sections through the allowed regions in the parameter space for appearance and disappearance experiments are shown. An opposite trend is clearly visible: while appearance data require non-zero values for the mixing of $\nu_e$ and $\nu_\mu$ with the eV-scale mass states 4 and 5 in order to explain LSND, disappearance data provide an upper bound on this mixing. The allowed regions touch each other at $\Delta \chi^2 = 9.3$, and a consistency test between these two data samples yields a probability of only 0.18%, i.e., these models can be considered as disfavored at the 3\sigma level.\cite{10} Also, because of the constraint from disappearance experiments the low energy excess in MiniBooNE can not be explained in the global analysis, see Fig. 2 (left, blue histogram). Furthermore, when moving from 4 neutrinos to 5 neutrinos the fit improves only by 6.1 units in $\chi^2$ by introducing 4 more parameters, showing that in (3+2) schemes the tension in the fit remains a sever problem. This is even true in the case of three sterile neutrinos, since adding one more neutrino to (3+2) cannot improve the situation.\cite{10}

3 Exotic proposals

Triggered by these problems many ideas have been presented in order to explain LSND, some of them involving very speculative physics, among them sterile neutrino decay,\cite{15,16}; violation
of the CPT\textsuperscript{17,12,18,19} and/or Lorentz\textsuperscript{20} symmetries, quantum decoherence\textsuperscript{21,22,23} mass-varying neutrinos\textsuperscript{24}, short-cuts of sterile neutrinos in extra dimensions\textsuperscript{25}, a non-standard energy dependence of sterile neutrinos\textsuperscript{26}, or sterile neutrinos interacting with a new gauge boson\textsuperscript{27}. In the following I comment on a personal selection of these exotic proposals, without the ambition of being complete.

\textit{CPT violation.} Triggered by the observation that the LSND signal is in anti-neutrinos, whereas their neutrino data is consistent with no oscillations, it was proposed\textsuperscript{17} that neutrinos and anti-neutrinos have different masses and mixing angles, which violates the CPT symmetry. A first challenge to this idea has been the KamLAND reactor results, which require a $\Delta m^2$ at the solar scale for anti-neutrinos. Subsequently it has been shown that the oscillation signature in SuperK atmospheric neutrino data (which cannot distinguish between $\nu$ and $\bar{\nu}$ events) is strong enough to require a $\Delta m^2 \sim 2.5 \cdot 10^{-3}$ eV$^2$ for neutrinos as well as for anti-neutrinos\textsuperscript{18}, see\textsuperscript{28} for an update. This rules out such an explanation of the LSND signal with three neutrinos at 4.6$\sigma$. However, introducing a sterile neutrino, and allowing for different masses and mixings for neutrinos and anti-neutrinos\textsuperscript{19} is fully consistent with all data, including the MiniBooNE null-result in neutrinos. Such a model should lead to a positive signal in the MiniBooNE anti-neutrino run.

\textit{Sterile neutrino decay.} Pre-MiniBooNE data can be fitted under the hypothesis\textsuperscript{16} of a sterile neutrino, which is produced in pion and muon decays because of a small mixing with muon neutrinos, $|U_{\mu 4}| \approx 0.04$, and then decays into an invisible scalar particle and a light neutrino, predominantly of the electron type. One needs values of $g m_4 \sim$ few eV, $g$ being the neutrino–scalar coupling and $m_4$ the heavy neutrino mass, e.g. $m_4$ in the range from 1 keV to 1 MeV and $g \sim 10^{-6} \ldots 10^{-3}$. This minimal model is in conflict with the null-result of MiniBooNE. It is possible to save this idea by introducing a second sterile neutrino, such that the two heavy neutrinos are very degenerate in mass. If the mass difference is comparable to the decay width, CP violation can be introduced in the decay, and the null-result of MiniBooNE can be reconciled with the LSND signal\textsuperscript{16}.

\textit{Sterile neutrinos with an exotic energy dependence.} Short-baseline data can be divided into low-energy (few MeV) reactor experiments, LSND and KARMEN around 40 MeV, and the high-energy (GeV range) experiments CDHS, MiniBooNE, NOMAD. Based on this observation it turns out that the problems of the fit in (3+1) schemes can be significantly alleviated if one
assumes that the mass or the mixing of the sterile neutrino depend on its energy in an exotic way. For example, assuming that $m_3^2(E_\nu) \propto E_{\nu}^{-7}$ one finds that for $r > 0$ the MiniBooNE exclusion curve is shifted to larger values of $\Delta m^2$, whereas the bound from disappearance experiments is moved towards larger values of the mixing angle, and hence the various data sets become consistent with LSND, compare Fig. 3 (left). At the best fit point with $r \approx 0.3$ the global fit improves by 12.7 units in $\chi^2$ with respect to the standard (3+1) fit. Similar improvement can be obtained if energy dependent mixing of the sterile neutrino is assumed.

Let us note that this is a purely phenomenological observation, and it seems difficult to construct explicit models for such sterile neutrinos. There are models which effectively introduce a non-standard “matter effect” for sterile neutrinos, e.g., via exotic extra dimensions or via postulating a new gauge interaction of the sterile neutrinos. Similar as in the usual MSW case, the sterile neutrino encounters effective mass and mixing which depend on energy. However, in these approaches the matter effect felt by the sterile state has to be some orders of magnitude larger than the standard weak-force matter effect of active neutrinos, in order to be relevant for short-baseline experiments. In such a case, in general very large effects are expected for long-baseline experiments such as MINOS, atmospheric neutrinos, or KamLAND. Unfortunately an explicit demonstration that a successful description of all these data can be maintained in such models is still lacking.

Quantum decoherence. The possibility that the origin of the LSND signal might be quantum decoherence in neutrino oscillations has been considered in

21,22,23. Such effects can be induced by interactions with a stochastic environment: a possible source for this kind of effect might be quantum gravity. The attempts to explain the LSND signal by quantum decoherence in 21,22 seem to be in conflict with present data. Both of these models are ruled out by the bound from NuTeV, $P_{\nu_e \rightarrow \nu_e}, P_{\bar{\nu}_e \rightarrow \bar{\nu}_e} < 5 \times 10^{-3}$ (90% C.L.) 29. Furthermore, the model of 21 (where in addition to decoherence, CPT-violation is also introduced which results in a difference between the oscillation probabilities for neutrinos and anti-neutrinos) cannot account for the spectral distortion in the anti-neutrino signal observed by KamLAND, whereas the scenario of 22 is disfavored by the absence of a signal in KARMEN, NOMAD and MiniBooNE.

Recently we have revisited this idea 23 by introducing a different set of decoherence parameters. We assume that only the neutrino mass state $\nu_3$ is affected by decoherence, whereas the 1-2
sector is completely unaffected, guaranteeing the standard explanation of solar and KamLAND data. Hence, denoting as $\gamma_{ij}$ the parameter which controls the decohering of the mass states $\nu_i$ and $\nu_j$, we have $\gamma_{12} = 0$ and $\gamma_{13} = \gamma_{23} \equiv \gamma$, where we have assumed that decoherence effects are diagonal in the mass basis. Furthermore, we assume that decoherence effects are suppressed for increasing neutrino energies, $\propto E^{-r}_\nu$ with $r \sim 4$. This makes sure that at short-baseline experiments with $E_\nu \gtrsim 1$ GeV such as MiniBooNE, CDHS, NOMAD, and NuTeV no signal is predicted, and at the same time maintains standard oscillations for atmospheric data and MINOS. In this way a satisfactory fit to the global data is obtained. Disappearance and appearance data become fully compatible with a probability of 74%, compared to 0.2% in the case of (3+2) oscillations. The LSND signal is linked to the mixing angle $\theta_{13}$, see Fig. 3(right) and hence, this scenario can be tested at upcoming $\theta_{13}$ searches: while the comparison of near and far detector measurements at reactors should lead to a null-result because of strong damping at low energies, a positive signal for $\theta_{13}$ is expected in long-baseline accelerator experiments.

4 Outlook

Currently MiniBooNE is taking data with anti-neutrinos. This measurement is of crucial importance to test scenarios involving CP (such as (3+2) oscillations) or even CPT violation to reconcile LSND and present MiniBooNE data. Therefore, despite the reduced flux and detection cross section of anti-neutrinos the hope is that enough data will be accumulated in order to achieve good sensitivity in the anti-neutrino mode. Furthermore, it is of high importance to settle the origin of the low energy excess in MiniBooNE. If this effect persists and does not find an “experimental” explanation such as an over-looked background, an explanation in terms of “new physics” seems to be extremely difficult. To the best of my knowledge, so-far no convincing model able to account for the sharp rise with energy while being consistent with global data has been provided yet.

The main goal of upcoming oscillation experiments like Double-Chooz, Daya Bay, T2K, NO$\nu$A is the search for the mixing angle $\theta_{13}$, with typical sensitivities of $30 \sin^2 2\theta_{13} \gtrsim 1\%$. This should be compared to the size of the appearance probability observed in LSND: $P_{\text{LSND}} \approx 0.26\%$. Hence, if $\theta_{13}$ is large enough to be found in those experiments sterile neutrinos may introduce some sub-leading effect, but their presence cannot be confused with a non-zero $\theta_{13}$. Nevertheless, I argue that it could be worth to look for sterile neutrino effects in the next generation of experiments. They would introduce (mostly energy averaged) effects, which could be visible as disappearance signals in the near detectors of these experiments. This has been discussed for the Double-Chooz experiment, but also the near detectors at superbeam experiments should be explored. An interesting effect of (3+2) schemes has been pointed out recently for high energy atmospheric neutrinos in neutrino telescopes. The crucial observation is that for $\Delta m^2 \sim 1$ eV$^2$ the MSW resonance occurs around TeV energies, which leads to large effects for atmospheric neutrinos in this energy range, potentially observable at neutrino telescopes. Another method to test sterile neutrino oscillations would be to put a radioactive source inside a detector with good spatial resolution, which would allow to observe the oscillation pattern within the detector. I stress that in a given exotic scenario such as the examples discussed in sec. 3 signatures in up-coming experiments might be different than for “conventional” sterile neutrino oscillations.

For the subsequent generation of oscillation experiments aiming at sub-percent level precision to test CP violation and the neutrino mass hierarchy, the question of LSND sterile neutrinos is highly relevant. They will lead to a miss-interpretation or (in the best case) to an inconsistency in the results. If eV scale steriles exist with mixing relevant for LSND the optimization in terms of baseline and $E_\nu$ of high precision experiments has to be significantly changed. Therefore, I argue that it is important to settle this question at high significance before decisions on high precision oscillation facilities are taken.
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NEUTRINO CROSS SECTION MEASUREMENTS FOR LONG-BASELINE ACCELERATOR-BASED NEUTRINO OSCILLATION EXPERIMENTS

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Neutrino oscillations are clear evidence for physics beyond the standard model. The goal of next-generation neutrino oscillation experiments is to find a non-zero $\theta_{13}$, the last mixing matrix element for which we only know an upper limit. For this, next-generation long-baseline neutrino oscillation experiments require an order of magnitude better sensitivities. In particular, accelerator-based experiments such as T2K and NOvA experiments need (1) good neutrino energy reconstruction for the precise measurement of $\Delta m^2_{32}$ and $\sin^2 2\theta_{23}$, and (2) good background prediction to measure $\nu_e$ appearance signals. Current and near future high statistics neutrino experiments, such as K2K, MiniBooNE, SciBooNE, MINOS, and MINERvA help both (1) and (2) by precise signal and background channel measurements.

1 next-generation long baseline accelerator-based neutrino oscillation experiments

The goal of next-generation long baseline accelerator-based neutrino oscillation experiments is to measure a non-zero $\theta_{13}$, the last mixing matrix element. The value of $\theta_{13}$ is the important parameter to access beyond the standard model physics. Especially if it were non-zero, then we hope to measure leptonic CP violation which can help to understand leptogenesis, one of the candidate explanations of baryon asymmetry of the universe $^1$.

Currently two experiments are planned, the Tokai-to-Kamioka (T2K) experiment $^2$ ($\sim 800$ MeV, $\sim 300$ km) and the NuMI Off-axis $\nu_e$ Appearance (NOvA) experiment $^3$ ($\sim 2$ GeV, $\sim 800$ km). Both experiments use a $\nu_\mu$ beam and search for $\nu_e$ appearance events to measure $\theta_{13}$ through the equation,

$$P(\nu_\mu \rightarrow \nu_e) = \sin^2 \theta_{23} \sin^2 2\theta_{13} \sin^2 \left( 1.27 \frac{\Delta m^2_{32}(eV^2)L(km)}{E(GeV)} \right).$$ (1)

Since a small $P(\nu_\mu \rightarrow \nu_e)$ is proportional to $\sin^2 \theta_{23}$ and $\sin^2 \left( 1.27 \frac{\Delta m^2_{e\mu}}{E} \right)$, we also need accurate knowledge of these two quantities, and can achieve by the measurements of $\nu_\mu$ disappearance


\[ P(\nu_\mu \rightarrow \nu_\mu) = 1 - \sin^2 2\theta_{23} \sin^2 \left( 1.27 \frac{\Delta m_{32}^2 (\text{eV}^2) L (\text{km})}{E(\text{GeV})} \right), \quad (2) \]

The oscillation parameters are extracted from the shape of \( P(\nu_\mu \rightarrow \nu_\mu) \), a function of reconstructed neutrino energy. Therefore a good extraction of \( \sin^2 2\theta_{23} \) and \( \Delta m_{32}^2 \) rely on good reconstruction of neutrino energy, which is based on better understanding of the signal (\( \nu_\mu \text{CCQE} \)) and background interactions, mainly CC1\( \pi^0 \) interaction (Sec. 2).

The signal of \( \nu_e \) appearance is an electron,

\[ \nu_e + n \rightarrow p + e^-. \quad (3) \]

There are many kind of possible backgrounds for this signal, for example, sometimes \( \nu_\mu \) induced NC\( \pi^0 \) production can mimic a \( \nu_e \) event if one of the decay photons from \( \pi^0 \) decay is undetected. Therefore, it is critical to understand this background channel (Sec. 3).

It is important to perform these cross section measurements prior to oscillation experiments. Although all long baseline accelerator-based neutrino oscillation experiments have near detectors, they exist to constrain neutrino flux uncertainties, and this constraint relies on accurate knowledge of cross section measurements. Fig. 1 shows the world’s data for charged current cross sections. As you can see, existing data are rather sparse and old. Since two experiments, T2K and NOvA, span different energy ranges, we need cross section measurements in both regions because the dominant interaction types will be different, and thus their energy reconstructions and backgrounds are different. Fortunately, we have a lot of new input from current and future neutrino cross section measurements: K2K near detector \(^5 \) (\( \sim 1.2 \text{ GeV} \), completed), MiniBooNE \(^6 \) (\( \sim 800 \text{ MeV} \), ongoing), SciBooNE \(^7 \) (\( \sim 800 \text{ MeV} \), ongoing), MINOS near detector \(^8 \) (\( \sim 2 - 20 \text{ GeV} \), ongoing), and MINERvA \(^9 \) (\( \sim 2 - 20 \text{ GeV} \), approved). We would like to discuss the two main themes of cross section related issues impacting oscillation searches, (1) neutrino energy reconstruction (Sec. 2), and (2) background determination (Sec. 3).

2 Neutrino energy reconstruction

2.1 Neutrino energy reconstruction for T2K

At the T2K energy scale (\( \sim 800 \text{ MeV} \)), the dominant neutrino reactions are \( \nu_\mu \) charged current quasi-elastic (CCQE) interactions,

\[ \nu_\mu + n \rightarrow p + \mu^-. \quad (4) \]

This channel is used to measure \( \nu_\mu \) disappearance, and thus the \( \nu_\mu \) energy reconstruction is critical. Since neutrino oscillation experiments use nuclear targets, understanding of this interaction is not trivial. Recently K2K \(^5,10 \) and MiniBooNE \(^6 \) have reported new measurements of the axial mass, \( M_A \), which are higher than the historical value (Table 1).

In this energy range, the axial vector form factor is the dominant contribution to the cross section and controls the \( Q^2 \) dependence. Inconsistency of their results from the world average, and the consistency between K2K and MiniBooNE is best understood in terms of nuclear effects, because most of the past experiments used deuterium targets whereas K2K and MiniBooNE used oxygen and carbon. Instead of using the world average, both experiments employ their measured \( M_A \) values to better simulate CCQE events in their oscillation analyses. After the
Figure 1: The world data for $\nu_\mu$ charged current cross section divided by neutrino energy. The dominant interaction for T2K and NOvA are quasi-elastic (QE) and deep inelastic scattering (DIS) respectively. The existing data are rather sparse and old, but we have more new input from current and future experiments!

Table 1: The comparison of measured axial mass $M_A$.

<table>
<thead>
<tr>
<th></th>
<th>$M_A$(GeV)</th>
<th>target</th>
</tr>
</thead>
<tbody>
<tr>
<td>K2K (SciFi) $^5$</td>
<td>1.20 ± 0.12</td>
<td>oxygen</td>
</tr>
<tr>
<td>K2K (SciBar) $^10$</td>
<td>1.14 ± 0.11</td>
<td>carbon</td>
</tr>
<tr>
<td>MiniBooNE $^6$</td>
<td>1.23 ± 0.20</td>
<td>carbon</td>
</tr>
<tr>
<td>world average $^{11}$</td>
<td>1.026 ± 0.021</td>
<td>deuteron, etc</td>
</tr>
</tbody>
</table>

Figure 2: (Left) K2K near detector complex. From the left to right, 1 kiloton water Čerenkov detector “1KT”, scintillation-fiber/water target tracker “SciFi”, fully active plastic organic scintillation-bar tracker “SciBar”, and muon range detector “MRD”. (Right) reconstructed $Q^2$ plot for 2-track QE sample from K2K SciFi, data (crosses) and simulation with best-fit $M_A$ (solid) agree well. The shaded region indicates the fraction of signal ($\nu_\mu CCQE$) events.
Figure 3: (Left) Schematic figure of a $\nu_\mu$ CCQE interaction in MiniBooNE. The MiniBooNE detector is a Čerenkov detector filled with mineral oil surrounded by PMTs. The Čerenkov light from the muon (Čerenkov 1) and subsequent Čerenkov light from the decayed electron are used to tag the CCQE event. (Middle) Event display of a muon candidate event in MiniBooNE. Each sphere represents a hit on a PMT, and size and color show charge and time information respectively. Muons create shape-edged Čerenkov ring. The ring center will appear filled-in if the muon is stopping in the tank. (Right) Reconstructed $Q^2$ plot of MiniBooNE, data (dots), simulation before the fit (dashed), and after the fit with $M_A$ and Pauli-blocking (solid). The dotted and dash-dotted lines indicate total background and irreducible background fraction respectively.

$M_A$ adjustment, both experiments see good agreement between data and simulation (Fig 2 and 3).

We can only measure the interaction rate, which is the convolution of flux and cross section \( R = \int \Phi \times \sigma \). So, without knowing flux prediction is perfect, one cannot tune the cross section model from measured interaction rate. MiniBooNE carefully examined this, and showed that their observed data simulation mismatching is not the effect of mismodeling of neutrino flux, but is really a cross section model problem. Fig 4 shows the ratio of data-simulation in the 2-dimensional plane made in muon kinetic energy and angle; left plot is before any cross section model tuning, right plot is after. The key point is that left plot clearly shows that data-simulation disagreements follow equal $Q^2$ lines, not equal $E_\nu$ lines.

\[
R = \int \Phi \times \sigma \rightarrow R[E_\nu, Q^2] = \int \Phi[E_\nu] \times \sigma[Q^2]
\]

This is strong evidence that the MiniBooNE data suggests a problem with the cross section model, and not the beam model, because cross section is the function of $Q^2$, whereas neutrino beam is a function of $E_\nu$.

It is not only important to understand the energy reconstruction of signal events \( i.e., \) CCQE interaction, but also for background channels. For Super-K, the neutrino energy is reconstructed
Figure 5: (Left) (a) CCQE interaction and (b) CC1π interaction. Eq. 6 correctly reconstructs neutrino energy only for (a), (b) can be distinguished from (a) by additional pion, however when pion is lost (by pion absorption for example), (b) becomes indistinguishable from intrinsic backgrounds. When (a) and (b) have the same muon kinematics, the reconstructed neutrino energies are the same, however the true neutrino energy for (b) is higher due to the creation of the pion in the event (neutrino energy mis-reconstruction). (Right) true and reconstructed neutrino energy distribution for Super-K predictions with neutrino oscillations. The shaded region is non-QE (mainly CC1π) channels. As can be seen from the bottom plot, CC1π background events are misreconstructed at lower neutrino energies and hence can fill out the dip created by neutrino oscillations.

Figure 6: (Left) charged current 1π production to CCQE cross section ratio from K2K SciBar analysis. Their result is consistent with past ANL bubble chamber experiment. (Right) charged current inclusive 1π+ production to CCQE cross section ratio from K2K SciBar analysis. Although the errors are large, the cross section obtained is significantly higher than the cross section model used in the K2K experiment.

from the measured muon energy $E\mu$ and angle $\theta\mu$, assuming a CCQE interaction,

$$E^{\nu}_{QE} \sim \frac{M_N E\mu - \frac{1}{2} m^2 \mu}{M_N - E\mu + \sqrt{E^2\mu - m^2 \mu \cos^2 \theta\mu}}.$$  \hspace{1cm} (6)

Here, $M_N$ and $m\mu$ are nucleon and muon masses. Since this formula assumes a 2-body interaction, any interaction involving more than two particles is a source of neutrino energy mis-reconstruction (Fig 5, left). The most notable channel contributing to this is charged current 1 π (CC1π) production. Especially when the detection of the outgoing pion fails for various reasons (pion absorption, detector effect, etc), CC1π events become an irreducible background, and thus they need to understand their relative contribution rather than rejecting them by cuts\(^4\) (Fig. 5, right).

Although neutrino absolute cross sections are notoriously difficult to measure due to uncertainties in the incoming neutrino flux, here they only need to know the kinematic distribution of CC1π events compared with CCQE events. Such measurements were done in K2K (Fig. 6)\(^{12,13}\) and MiniBooNE\(^{14}\).
Figure 7: (Left) SciBooNE detector. It consists of 3 parts, organic plastic scintillation-bar tracker “SciBar”, 11 radiation length lead electromagnetic calorimeter “EC”, and muon range detector “MRD” which can range out muons up to 0.9 GeV. (Middle) SciBooNE event display for $\nu_\mu$ CCQE candidate event. two tracks are seen in “SciBar”, then the longer track (muon) produce hits in both “EC” and “MRD”. (Right) Under the assumption of target nucleon at rest, muon energy and angle completely specify CCQE kinematics, i.e., one can predict the angle of outgoing proton. $\Delta \theta_p$ is defined as an opening angle of this predicted proton track and measured proton track. (a) is the case of CCQE interaction, and $\Delta \theta_p$ is small. However, (b) CC1$\pi$ interaction with invisible pion, $\Delta \theta_p$ is large because predicted track is based on the assumption of 2-body interaction but actual interaction is 3-body.

The SciBooNE experiment$^7$ at FNAL is particularly designed for this purpose (Fig. 7, left and middle). The SciBooNE vertex detector “SciBar”, formerly used at K2K experiment and shipped from Japan to Fermilab, is a high resolution tracker consisting of X-Y plastic organic scintillators with wavelength shifting fibers through the middle of each bar. Since SciBar can reconstruct both proton and muon tracks in a $\nu_\mu$ CCQE interaction (unlike Čerenkov detectors), so the opening angle of the measured proton and the expected outgoing proton (assuming CCQE kinematics) can be used to separate CCQE and CC1$\pi$ events, even in cases where the pion is undetected (right plot of Fig. 7). The goal of the SciBooNE experiment is to measure non-QE to CCQE cross section ratio to 5%, making the non-QE mis-reconstruction uncertainty for T2K negligible$^7$.

2.2 Neutrino energy reconstruction for NOvA and MINOS

The situation is quite different for higher energy scales (~2 GeV). The CCQE assumption is no longer held and calorimetric energy reconstruction provides a much more efficient energy determination:

$$E_\nu \sim E_\mu + E_{\text{showers}}$$  \hspace{1cm} (7)

Here, $E_\mu$ is the energy of muon, usually measured by a muon spectrometer which consists of a dense material to stop muons. $E_{\text{showers}}$ is the energy of both electromagnetic and hadronic showers measured in the calorimeter. This energy reconstruction method is successfully tested by the Main Injector Neutrino Oscillation Search (MINOS) experiment$^8$.

Neutrino energy misreconstruction happens, for example, when hadronic showers are absorbed by nuclei (Fig. 8, left). This is important for precise $\nu_\mu$ disappearance measurements by MINOS, where steel is used as a target but no reliable pion absorption measurements are available. The future Main Injector Experiment for $\nu$-A (MINERvA) has the ability to switch its target and they plan to study nuclear effects (Fig. 8, middle and left) as well as various physics topics from quasi-elastic to DIS$^9$. Their measurements will significantly reduce the uncertainties on $\Delta m_{23}^2$ coming from nuclear cross section modeling in MINOS$^9$. 

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3 Background channel

Since T2K uses water Čerenkov detector “Super-K” as a far detector, the signal of $\theta_{13}$, namely $\nu_e$ appearance is a single electron (Eq. 3) because outgoing protons are below Čerenkov threshold in most cases and therefore invisible.

The notorious background for this signal is the neutral current $\pi^0$ (NC$\pi^0$) interaction,

$$\nu_\mu + N \rightarrow \nu_\mu + N + \pi^0.$$  

(8)

Although $\pi^0$ decays to two photons, there are various reasons to miss one of them, for example, two photons overlap, or one photon is boosted to low energy below threshold. The precise prediction of this channel is critical for any $\nu_e$ appearance experiments. K2K measured the NC$\pi^0$ rate using 1KT detector\(^ {15}\).

Recently, the MiniBooNE experiment made an in-situ measurement of NC$\pi^0$ production on mineral oil which was used to predict background processes more precisely for their $\nu_e$ appearance search\(^ {16}\). Even though the underlying source of the $\pi^0$ may not be known, (\textit{i.e.}, actual resonance model to create the $\pi^0$ is not clear), the difference between the observed and predicted kinematic distribution of $\pi^0$’s can be used to correct the rate of $\pi^0$ events that are misclassified as $\nu_e$ signal events. Since the loss of a photon in the $\pi^0$ decay is mostly a kinematic effect, once correct $\pi^0$ production kinematics are obtained from the data, it is easy to calculate the distribution of $\pi^0$ where one photon is missed. Left plot of Fig. 9 shows data-simulation comparisons for pion mass peak. After the correction, their simulation precisely predicts all observed aspects of NC$\pi^0$ events. The right plot of Fig. 9 shows a kinematic distribution.

This result triggered another interest. This plot clearly shows the existence of NC coherent pion production. However, the K2K experiment saw no evidence for CC coherent pion production at similar energies\(^ {12}\). Since a coherently produced pion has very different kinematics, understanding of this rate is important. Again, further analysis of K2K, MiniBooNE, SciBooNE, MINOS, and MINERvA will shed light on this in the near future.

The fine-grained MINERvA detector will provide critical input for NOvA. Although high statistics data from K2K, MiniBooNE, and SciBooNE will be available, backgrounds of $\nu_e$ appearance search around $\sim 2$ GeV is only effectively accessible by MINERvA experiments. We are expecting negligible cross section error on $\sin^22\theta_{13}$ from NOvA after precise CC and NC measurements from MINERvA\(^ {9}\).
4 Conclusions

The goal of next-generation long baseline accelerator-based neutrino oscillation experiments is to measure a $\nu_e$ appearance signal. The cross section errors arise from (1) misreconstruction of neutrino energy and (2) incorrect background predictions. The inputs from current and future neutrino cross section measurements are critical to the success of future oscillation experiments, such as T2K and NOvA.

References

HARP Collaboration results on the proton-nuclei interactions at few GeV energies

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(on behalf of the HARP Collaboration)

Recent results obtained by the HARP collaboration on the measurements of the double-differential production cross-section of positive and negative pions in proton interactions with nuclear targets from Beryllium to Lead are presented. They cover production at small angles (30-210 mrad) and relatively large momenta up to 8 GeV/c as well as large angles (0.35 - 2.15 rad) and small momenta (0.1 - 0.8 GeV/c). These results are relevant for a detailed understanding of neutrino fluxes in accelerator neutrino experiments, better prediction of atmospheric neutrino fluxes, optimization of a future neutrino factory design and for improvement of hadronic generators widely used by the HEP community in the simulation of hadronic interactions.

1 The HARP experiment

The HARP experiment\(^1,2\) at the CERN PS was designed to make measurements of hadron yields from a large range of nuclear targets and for incident particle momenta from 1.5 GeV/c to 15 GeV/c. The main motivations are the measurement of pion yields for a quantitative design of the proton driver of a future neutrino factory, a substantial improvement in the calculation of the atmospheric neutrino flux and the measurement of particle yields as input for the flux calculation of accelerator neutrino experiments, such as K2K\(^3,4\), MiniBooNE\(^5\) and SciBooNE\(^6\).

The experiment makes use of a large-acceptance spectrometer consisting of a forward and large-angle detection systems. A detailed description of the experimental apparatus can be found in Ref.\(^2\). The forward spectrometer – based on large area drift chambers and a dipole magnet complemented by a set of detectors for particle identification\(^8\); a time-of-flight wall, a large Cherenkov detector and an electromagnetic calorimeter – covers polar angles up to 250 mrad which is well matched to the angular range of interest for the conventional neutrino beams. The large-angle spectrometer – based on a Time Projection Chamber (TPC) and Resistive Plate Chambers (RPCs), located inside a solenoidal magnet – has a large acceptance in the momentum and angular range for the pions relevant to the production of the muons in a neutrino factory. It covers the large majority (~ 70%) of the pions accepted in the focusing system of a typical design.
Figure 1: Overall mechanical layout of the HARP detector. The different sub-detectors are shown. The target is inserted inside the TPC.

Figure 2: Double-differential production cross-section of $\pi^+$ and $\pi^-$ in p-C reactions at 12 GeV/c (points with error bars) and comparison with model predictions.
2 Results obtained with the HARP forward spectrometer

The first HARP physics publication reported measurements of the $\pi^+$ production cross-section from an aluminum target at 12.9 GeV/c proton momentum for the K2K experiment at KEK PS. The results were subsequently applied to the final neutrino oscillation analysis of K2K, allowing a significant reduction of the dominant systematic error associated with the calculation of the so-called far-to-near ratio. Our next result was the measurement of the $\pi^+$ cross-sections from a thin 5% $\lambda_I$ beryllium target at 8.9 GeV/c proton momentum. It contributed to the understanding of the MiniBooNE and SciBooNE neutrino fluxes. They are both produced by the Booster Neutrino Beam at Fermilab which originates from protons accelerated to 8.9 GeV/c by the booster before being collided against a beryllium target.

Further, measurements of the double-differential production cross-section of $\pi^\pm$ in the collision of 12 GeV/c protons and $\pi^\pm$ with thin 5% $\lambda_I$ carbon target and liquid N$_2$ and O$_2$ targets were performed. These measurements are important for a precise calculation of the atmospheric neutrino flux and for a prediction of the development of extended air showers. The results for the pion production on the carbon target, the ratio N$_2$/Carbon, and comparison with models typically used in air shower simulations are shown in Figs. 2 and 3. The conclusion of comparing the predictions of the models to the measured data is that they do predict the ratio of cross-sections and often fail in predicting the absolute rates, especially in certain regions of the phase space.

In practice production targets are not thin and cascade calculations or dedicated measurements with 'replica targets' are needed. HARP has taken, albeit with somewhat lower statistics, and analyzed p+A data at different beam momenta with 100% $\lambda_I$ targets. They can be used for parametrizations or tuning of models. Preliminary spectra are available for p + Be, C, Al, Cu, Sn, Ta, Pb interactions at 3 – 12 GeV/c. The measurements are on the tapes and can be analyzed on demand.
3 Results obtained with the HARP large-angle spectrometer

The HARP TPC is the key detector for the analysis of tracks emerging from the target at large angles with respect to the incoming beam direction. It suffered from a number of shortcomings that were discovered during and after the data taking\(^2\). A description of the measures taken to correct for the effects of them is given in \(^2,15,16\). Wide range of experimental cross-checks has been employed to assess the momentum scale and momentum resolution in the HARP TPC, summarized in our recent paper\(^16\).

A group of people formerly belonging to the HARP collaboration and subsequently detached themselves from it have been criticizing our methods of TPC and RPCs calibration\(^13\). Our arguments against this criticism and for the correctness of our results are presented in \(^14,16\).

A first set of results on the production of pions at large angles have been published by the HARP collaboration in the papers\(^15,18\), based on the analysis of the data in the beginning of each accelerator spill. Track recognition, momentum determination and particle identification were all performed based on the measurements made with the TPC. The reduction of the data set was necessary to avoid problems in the chamber responsible for dynamic distortions to the image of the particle trajectories as the ion charge was building up during each spill. Corrections for such distortions that allow the use of the full statistics have been developed\(^16\) and applied in the analysis. The results exploiting the full spill data have been obtained recently\(^17\). They are fully compatible with the previous ones and cover pion production by proton beams in a momentum range from 3 GeV/c to 12 GeV/c hitting Be, C, Al, Cu, Sn, Ta and Pb targets with a thickness of 5% \(\lambda\) in the angular and momentum phase space \(100\ \text{MeV/c} \leq p < 800\ \text{MeV/c}\) and 0.35 rad \(\leq \theta < 2.15\) rad in the laboratory frame.

As an example we show in Fig. 4 the results for the double-differential cross-sections \(d^2\sigma/dpd\theta\) at 5 GeV/c incident proton beam momentum and Ta target compared to the respective predictions of several different generator models used in GEANT4 and MARS simulation packages. The comparison between data and models is reasonable, but some discrepancies are evident for

![Figure 4: Double-differential π⁺ production cross sections for p-Ta at 5 GeV/c and comparison with GEANT4 and MARS MC predictions, using several generator models.](image)
Figure 5: The dependence on the beam momentum and on the atomic number $A$ of the $\pi^-$ (right) and $\pi^+$ (left) production yields in $p$–Be, $p$–C, $p$–Al, $p$–Cu, $p$–Sn, $p$–Ta, $p$–Pb interactions averaged over the forward angular region ($0.350 \text{ rad} \leq \theta < 1.550 \text{ rad}$) and momentum region $100 \text{ MeV}/c \leq p < 700 \text{ MeV}/c$. The results are given in arbitrary units, with a consistent scale for all panels.

Figure 6: Predictions of the $\pi^+$ (closed symbols) and $\pi^-$ (open symbols) yields for different designs of the neutrino factory focusing stage. Integrated yields and the integrated yields normalized to the kinetic energy of the proton for $p$–Ta and $p$–Pb interactions. The circles indicate the integral over the full HARP acceptance ($100 \text{ MeV}/c < p < 700 \text{ MeV}/c$ and $0.35 \text{ rad} < \theta < 0.95 \text{ rad}$), the squares are integrated over $0.35 \text{ rad} < \theta < 0.95 \text{ rad}$, while the diamonds are calculated for the smaller angular range and $250 \text{ MeV}/c < p < 500 \text{ MeV}/c$. Although the units are indicated as “arbitrary”, for the largest region the yield is expressed as $d^2\sigma/dp d\Omega$ in mb/(GeV/c sr). For the other regions the same normalization is chosen, but now scaled with the relative bin size to show visually the correct ratio of number of pions produced in these kinematic regions.
some models. Discrepancies up to a factor of three are seen. For details see the full paper\textsuperscript{17}. The dependence on the beam momentum and on the atomic number $A$ of integrated yields are presented in Fig. 5. Predictions of the $\pi^+$ and $\pi^-$ integrated yields relevant for the design of the neutrino factory focusing stage are given in Fig. 6.

4 Conclusions

The full set of HARP data is in process of publishing now. It covers the pion production by protons and pions on nuclear targets spanning the full periodic table of elements and large solid angle and momentum range in the difficult energy region between 3 and 15 GeV/c of incident momentum. HARP results fill in an essential gap in the available experimental information for soft hadron production and help in the understanding of neutrino fluxes in accelerator neutrino experiments, prediction of atmospheric neutrino fluxes, optimization of a future neutrino factory design and may be used for improvements of the event generators for simulation of hadronic interactions.

Acknowledgments

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References

LOW ENERGY EFFECTS OF SEESAW NEUTRINO MASS MODELS

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Seesaw models all lead to the same universal dimension 5 low-energy effective operator. As a result effects of this lepton number violating operator, including neutrino masses, cannot distinguish these models. However, effects of the dimension 6 operators, which conserve lepton number and differ for each model, could. For all seesaw models we perform a general analysis of the structure and phenomenology of dimension 6 operators. We discuss how the associated effects could be within experimental reach if a decoupling between dimension 5 and 6 operator coefficients occurs.

1 Introduction

The experimental observation of non-zero neutrino masses and mixings constitutes evidence for physics beyond the Standard Model (SM) and points to the existence of a new, yet unknown, physics scale. If this scale, $M$, is larger than the electroweak scale, the low energy effects of this new physics can be phrased in terms of a generic effective theory, order by order in $1/M$. At lowest order in this expansion, $\mathcal{O}(1/M)$, the symmetries and particle content of the standard model allow only one operator

$$\delta \mathcal{L}^{d=5} = \frac{1}{2} c_{\alpha \beta} \bar{\ell}_L^\alpha \delta \phi^\dagger \ell_L^\beta + \text{h.c.}$$

where $\ell_L$ stands for the lepton weak doublets, greek letters denote flavour indices and $\delta \phi$ is related to the standard Brout-Englert-Higgs (Higgs for short) doublet $\phi \equiv (\phi^+, \phi^0)$ by $\delta \phi = i \tau_2 \phi^*$. This dimension 5 operator has the particularity to induce Majorana neutrino masses once the Higgs acquires a vacuum expectation value: $m_\nu = -\frac{v^2}{2} c^{d=5}$, with $\langle \phi^0 \rangle = v/\sqrt{2} = 174 \text{ GeV}$. This means that the neutrino masses, which constitute the first ever observed laboratory evidence for physics beyond the Standard Model, is nothing but the first evidence for physics beyond the Standard Model we could have expected in the $1/M$ expansion. Moreover in such a way
the smallness of the neutrino masses can find a nice explanation: the seesaw mechanism: the neutrino masses are suppressed by $1/M$, so are naturally small if $M$ is large. These facts, as well as the absence of exotic experimental signals other than neutrino masses strongly point for such a heavy scale explanation of the neutrino masses. This leads to the question of how to generate the dimension 5 operator from explicit new physics models. From the tree level exchange of heavy particles there are only 3 basic ways to generate it: through right-handed neutrino exchange (type-I seesaw), scalar Higgs triplet exchange (type-II seesaw) or fermionic triplet exchange (type-III seesaw).

However since the dimension 5 operator is unique, and since one can show that any of these 3 models could generate any flavour structure for the $d_{\alpha\beta}^5$ matrix, these models cannot be distinguished from it. To distinguish them one must therefore rely on possible effects at next order in the $1/M$ expansion, that is to say from dimension 6 operators, which turn out to be different in each model.

### 2 Determination of the dimension 6 operators

If the SM is augmented by right-handed neutrinos $N_R$, in full generality there are 3 new terms which can be written

$$\mathcal{L} \ni i \bar{N}_R \tilde{\theta} N_R - \bar{\ell}_L \tilde{\phi} Y^i_N N_R - \frac{1}{2} \bar{N}_R M_N N_R^c + \text{h.c.}$$

with $Y_N$ the Yukawa interactions and $M_N$ the Majorana mass term, i.e. the new physics scale(s). Flavour indices are implicit in these expressions and we will work in a basis in which $M_N$ is a diagonal complex matrix.

Alternatively if the SM is augmented by a scalar triplet with hypercharge $2$, $\Delta = (\Delta_1, \Delta_2, \Delta_3)$ related to the charge eigenstates by $\Delta^{++} \equiv \frac{1}{\sqrt{2}}(\Delta^1 - i \Delta^2)$, $\Delta^+ \equiv \Delta^3$, $\Delta^0 \equiv \frac{1}{\sqrt{2}}(\Delta^1 + i \Delta^2)$, there are 8 new terms which can be written

$$\mathcal{L}_{\Delta} = \left( D_{\mu} \Delta \right) \left( \bar{D}^{\mu} \Delta \right) + \left( \bar{\ell}_L Y_\Delta (\bar{\tau} \cdot \Delta) \ell_L + \mu_\Delta \tilde{\phi} (\bar{\tau} \cdot \Delta)^i \phi + \text{h.c.} \right)$$

$$- \frac{1}{2} \lambda_2 \left( \Delta^i \Delta^i \right)^2 + \lambda_3 \left( \tilde{\phi}^i \phi \right) \left( \Delta^i \Delta^i \right) + \frac{\lambda_4}{2} \left( \Delta^i T^i \Delta \right)^2 + \lambda_5 \left( \Delta^i T^i \Delta \right) \phi^i \tau^i \phi$$

where summation over the $SU(2)$ indices $i$ and flavour indices is assumed, with $T^i$ the 3 by 3 generators of $SU(2)_L$. Particularly important for the neutrino masses are the $Y_\Delta$ Yukawa interactions, and the $\mu_\Delta$ trilinear scalar interaction, which together induce the dimension 5 operator from the exchange of the scalar triplet.

Alternatively if the SM is augmented by hypercharge 0 fermion triplets, $\Sigma = (\Sigma^1, \Sigma^2, \Sigma^3)$, related to the charge components by $\Sigma^\pm \equiv \frac{\Sigma^1 \pm i \Sigma^2}{\sqrt{2}}$, $\Sigma^0 \equiv \Sigma^3$, there are 3 new terms which can be written:

$$\mathcal{L}_{\Sigma} = i \bar{\Sigma}_R D^\mu \Sigma_R - \left[ \frac{1}{2} \bar{\Sigma}_R M_{12} \Sigma_R + \Sigma_R Y_{\Sigma} (\tilde{\phi}^i \tau^i \ell_L) + \text{h.c.} \right].$$

Integrating the heavy fields, using a standard procedure whose details can be found in $^2$, these three models generate the dimension 5 operator, as well as various dimension 6 operators. The structure of these operators as well the form of their coefficients as a function of the various couplings are given in Table 1.

The main features of the dimension 6 operators are the following. In the fermionic seesaw model, type-I and type-III, the (unique) dimension 6 operator is derivative and, once $\phi$ is replaced by its vev, induces a non flavour diagonal contribution to the kinetic terms of the light leptons. This requires to redefine the light lepton fields by a non-unitary transformation to
Table 1: Coefficients of the $d = 5$ operator, $c^{d=5}$, and $d = 6$ operators and their coefficients, $c^{d=6}$, in the three basic seesaw theories.

<table>
<thead>
<tr>
<th>Model</th>
<th>Effective Lagrangian $\mathcal{L}_{\text{eff}} = c_i \mathcal{O}_i$</th>
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<tbody>
<tr>
<td></td>
<td>$c^{d=5}$</td>
</tr>
<tr>
<td>Fermionic Singlet</td>
<td>$Y^T \frac{1}{M_N} Y_N$</td>
</tr>
<tr>
<td>Scalar Triplet</td>
<td>$4Y \Delta \frac{\mu \lambda}{M_N^2}$</td>
</tr>
<tr>
<td></td>
<td></td>
</tr>
<tr>
<td>Fermionic Triplet</td>
<td>$Y^T \frac{1}{M_N} Y_N$</td>
</tr>
</tbody>
</table>

get properly normalized and flavour diagonal kinetic terms. This redefinition applied to the gauge interaction terms leads to non flavour diagonal gauge interactions. In the type-I model, only the kinetic term of neutrinos is affected, which can be understood from the fact that only the neutrinos mix with the right-handed neutrinos. Therefore the gauge interactions affected by this redefinition are the ones involving a light neutrino, $W^- l_i \nu_j$ interactions and $Z^- \nu_i \nu_j$ interactions. For the type-III model both neutrino and charged lepton fields must be redefined, both kinetic terms are affected, because both neutrinos and charged leptons mix with the neutral and charged components of the fermion triplets respectively. Consequently non flavour diagonal $Z^- l_i l_j$ are induced too. This interaction is obtained, in other words, from a $l_i \Sigma^- l_j$ transition with the $Z$ attached to any of these three particles. This leads to a very rich rare lepton process phenomenology, as discussed in section 4 below. For the type-II model since the heavy state is a scalar there is no fermion mixing induced, i.e. no non-unitarity effects, but alternatively a four-lepton interaction is induced through simple exchange of the scalar triplet. This also leads to a very rich phenomenology.

3 What about the size of dimension 6 operator effects?

In general, given the size of the neutrino masses, the magnitude of the dimension 6 operator effects is expected far below the sensitivity of present experiments. For example in the type-I model both dim-5 and dim-6 coefficients in Table 1 involve 2 Yukawa couplings so that one could in general expect that $c_{d=6} \simeq c_{d=5}/M_N \sim m_\nu/M_N v^2$ which even for $M_N$ as low as 1 TeV is very suppressed ($m_\nu/M_N \sim 10^{-13}$ in this case). With such coefficient the observation of any rare lepton process is hopeless. However this naive estimate turns out not to be necessarily valid at all. While the dimension 5 operator breaks lepton number, the dimension 6 ones do not. This reflects itself in the fact that both coefficients do not involve the same combinations of the Yukawa couplings even if they both involve 2 Yukawa couplings, see Table 1. As a result there is no symmetry reason why, if the neutrino mass matrix entries, so $c_{d=5}$ coefficients, are suppressed, the $c_{d=6}$ coefficients should also be.

Any scenario which would proceed in the following 2 steps would result in unsuppressed $c_{d=6}$ coefficients even if the $c_{d=5}$ ones do are suppressed: first assume a scenario where lepton number
is conserved with not too large heavy state mass $M$ and with large Yukawa couplings $Y$. In this case one gets $c_{d=6} \sim Y^2/M^2$ large and $m_\nu = 0$. Then introduce a small source of lepton number violation from a small perturbation parameter $\mu$: this small perturbation will not affect much the dim-6 coefficients but will lead to dimension 5 coefficients naturally suppressed by an extra $\mu/M$ factor, one gets $c_{d=5} \sim f(Y)\mu/M^2$, i.e. $m_\nu \sim f(Y)\mu^2/M$, with $f(Y)$ a function of the Yukawa couplings.

This mechanism which we call Direct Lepton number Violation (because the suppression of neutrino masses comes proportionally to a small entry in the numerator rather than from a large entry in the denominator), turns out to be possible in each of the 3 seesaw models. In the type-II model it is straightforward. Since the neutrino mass and the dim-6 four lepton operator coefficients have already the DLV form from the beginning, see Table 1, the decoupling of both coefficients is automatic. Large dim-6 coefficients require large $Y_\Delta$ couplings with not too large scalar triplet mass, while the size of neutrino masses involve another parameter, $\mu_\Delta$, which can be taken small to have sufficiently suppressed neutrino masses. Lepton number is conserved if $\mu_\Delta$ vanishes. In type-I and type-III scenarios the DLV scenario is much less automatic but turns out to be feasible too. For example for the simple case with one light neutrino $\nu$ and 2 right-handed neutrinos $N_{1,2}$, assume that the lepton number of these 3 particles is 1, −1, 1 respectively. In this case lepton number conservation allows two heavy states in the $N_1$-$N_2$ sector, from the $MN_1N_2$ bilinear term, while it also allows a Yukawa interaction of the type $Y\nu_i\phi_1$. These 2 terms lead to a non-vanishing dim-6 coefficient which can be large, $c_{d=6} \sim Y^2/M^2$, but they do not lead to any neutrino mass since lepton number is conserved, no matter the size of $Y$ and $M$. Introduce now a small L violating term, for example of the form $\mu N_2 N_2$: this leads to neutrino masses of the form $Y^2\mu^2/M^2$ just as in the DLV framework above. This mechanism for the type-I model is known as the "inverse seesaw" mechanism. It can be generalized to the 3 light neutrino plus 3 heavy neutrino case and works also for the type-III model just in the same way as in the type-I model. Of course it requires to give up one of the virtue of the canonical GUT seesaw mechanism which is that neutrino masses can be obtained with just enough suppression from assuming heavy states with mass not far from the GUT scale. But adopting a phenomenological point of view that any possible observable effects of the physics at the origin of the neutrino masses should be searched for, it is interesting to study this possibility in details.

4 Phenomenology of dimension 6 operators effects

There is a long list of rare processes which can be induced by dimension 6 operators: rare lepton decays: $\mu \to e\gamma$, $\tau \to \mu\gamma$, $\tau \to e\gamma$, $\mu \to eee$, $\tau \to 3\ell$,...; deviations to universality of gauge interactions in $W \to l\bar{\nu}$, $\tau \to l\nu\bar{\nu}$, $Z \to \ell\bar{\ell}$, $\pi \to l\nu\bar{\nu}$,...; flavour changing $Z$ decays: $Z \to l_i\bar{l}_j$; $Z$ invisible width from $Z \to \nu\bar{\nu}$ decays; correction to the $\rho$ parameter and $W$ mass in type II model;... Since none of these effects has yet been observed this leads to upper bounds on the dim-6 coefficients, that is to say on combinations of Yukawa couplings and heavy state masses. For the type-I model, combining all constraints, we get the following bounds on the coefficients of the unique dim-6 operator in Table 1:

$$\frac{v^2}{2} |c_{d=6}|_{a\beta} = \frac{v^2}{2} \left| \frac{1}{|M_N|^2} Y_N \right|_{a\beta} \begin{pmatrix} 10^{-2} & 7.0 \cdot 10^{-5} & 1.6 \cdot 10^{-2} \\ 7.0 \cdot 10^{-5} & 10^{-2} & 1.0 \cdot 10^{-2} \\ 1.6 \cdot 10^{-2} & 1.0 \cdot 10^{-2} & 10^{-2} \end{pmatrix}.$$ (5)

Similarly for type-III one obtains:

$$\frac{v^2}{2} \left| c_{d=6} \right|_{a\beta} = \frac{v^2}{2} \left| \frac{1}{|M_\Sigma|^2} Y_\Sigma \right|_{a\beta} \begin{pmatrix} 3 \cdot 10^{-3} & 1.1 \cdot 10^{-6} & 1.2 \cdot 10^{-3} \\ 1.1 \cdot 10^{-6} & 4 \cdot 10^{-3} & 1.2 \cdot 10^{-3} \\ 1.2 \cdot 10^{-3} & 1.2 \cdot 10^{-3} & 4 \cdot 10^{-3} \end{pmatrix}.$$ (6)
The bounds in the type-III model are stronger than in the type-I model, especially for the off-diagonal entries. This is due to the fact that, as said above, in the type-III model, unlike in the type-I model, flavour changing processes with charged fermions are generated already at tree level through the $Z$-$l_l$-$l_j$ couplings. In particular the most stringent bound, which is in the $\mu$-$e$ channel, comes from $\mu \rightarrow eee$. It is induced at tree level in the type-III model from a $\mu$-$e$ transition on a same fermionic line with emission of a $Z$ which decays in an electron pair. In the type-I model this transition can be done only at the one-loop level and leads to a weaker bound than from $\mu \rightarrow e\gamma$ (also induced at one-loop). Similarly the two other off-diagonal constraints in Eq. (6) come from $\tau \rightarrow \mu^+\mu^-$ and $\tau \rightarrow e^+e^-e^-$ in the same way, while for the type-I model in Eq. (5) they come from $\tau \rightarrow l\gamma$. As for the diagonal entries they come from universality tests. The constraints from $\mu \rightarrow e\gamma$ and $\tau \rightarrow l\gamma$ in the type-III model have been discussed in this conference by Florian Bonnet. A very interesting feature of the type-III model is that it predicts fixed ratios between the $\mu \rightarrow eee$ and $\mu \rightarrow e\gamma$ rates and similarly for $\tau$ decays. This offers the opportunity to test the model in a particularly clean way. An interesting bound also arises from $\mu$ to $e$ conversion in atomic nuclei. As for the effect of the type-III model to the $g-2$ of the muon, it turns out to be too small to be able to explain the discrepancy between theoretical expectations and experimental data.

Similarly in the type-II model one gets a series of constraints whose most important ones are given in Table. 2. As in the type-III model, $\mu \rightarrow eee$ and $\tau \rightarrow 3f$ are induced at tree level (from the simple exchange of a scalar triplet) and give the best constraints on the off-diagonal elements.

The general trend for the bounds above is that the Yukawa couplings must be smaller than $\sim 10^{-1} \cdot (M/1 \text{ TeV})$ or smaller in some cases. Note that these bounds already basically exclude the possibility to produce the heavy states at LHC through Yukawa driven processes even for mass below the LHC energies, see also the discussion by F. del Aguila in this conference. Drell-Yan pair production of type-II and type-III heavy states is nevertheless perfectly possible.
if they are light enough. If all couplings were of order unity the bounds above translate in a lower bound for the masses of the heavy states: $M_N > 21$ TeV, $M_\Delta > 294$ TeV and $M_\Sigma > 170$ TeV. More details can be found in $^{2,8}$.

## 5 Summary

Rare processes associated to dimension 6 operator effects present a unique opportunity for distinguishing seesaw neutrino mass models, especially if the heavy states are too heavy to be produced at colliders. Although in general these effects are expected to be far from being within experimental reach, it is not excluded that they might be large enough to be reachable soon in special cases. Such cases can be justified within the (common to all seesaw models) Direct Lepton number Violation theoretical framework, based on the assumption of approximate lepton number conservation. The associated phenomenology is very rich and leads, through a long series of processes, to the Yukawa couplings upper bounds given in section 4. This provides a strong motivation for the new generation of experiments aiming to measure rare leptonic processes.$^{12}$

## Acknowledgments

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NEUTRINOLESS DOUBLE BETA DECAY SEARCH WITH
CUORICINO AND CUORE EXPERIMENTS

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CUORICINO is a bolometric experiment on Neutrinoless Double Decay (0ν-DBD). With its 40.7 kg mass of TeO$_2$ it is the most massive 0ν-DBD presently running and it has proven the feasibility of the CUORE experiment, whose aim is to be sensitive to the effective neutrino mass down to few tens of meV. We report here latest CUORICINO results and prospects for the future CUORE experiment.

1 Introduction

The positive results obtained in the last few years in neutrino oscillation experiments have given convincing and model independent evidences that neutrinos are massive and mixed particles. The obtained data are compatible with two possible mass patterns, or hierarchies, the normal: $m_1 < m_2 < m_3$, and the inverted hierarchy: $m_3 < m_1 < m_2$. Unfortunately oscillation experiments are only sensitive to neutrino mass squared differences and cannot give any information about neutrino nature (Dirac or Majorana particle) and absolute mass scale. Beta decay experiments are sensitive to the absolute mass scale but cannot determine neutrinos nature. Experiments looking for the Neutrinoless Double Beta Decay (0ν-DBD) of even-even nuclei have the highest sensitivity to the effective neutrino mass ($m_{\beta\beta}$) and to the neutrino nature. A positive signal would imply that neutrino are Majorana particles and can even lead to the measure of the absolute mass scale.

The use of the bolometric technique offers the unique possibility to investigate different 0ν-DBD candidates with a considerable high energy resolution, needed to separate the 2ν contribution from the 0ν peak. The CUORE experiment$^1$, to search 0ν-DBD of $^{130}$Te, will start its assembling phase in 2008 and it aims to reach a sensitivity on $m_{\beta\beta}$ better than 50 meV.
CUORICINO\textsuperscript{2} represent not only the first stage of CUORE, but also the most massive $0\nu$-DBD experiment presently running.

\section{Bolometric Technique}

Bolometers are sensitive calorimeters that measure the energy deposited by particle or photon interactions by measuring the corresponding rise in temperature. The CUORICINO bolometers are single crystals of TeO\textsubscript{2} that are dielectric and diamagnetic. In these materials the heat capacity is due almost exclusively to the lattice degrees of freedom and its value at 10 mK ($\sim 2 \times 10^{-9}$J·K\textsuperscript{-1}) has been extrapolated from the measured value at 232 K. With these values of the parameters, an energy deposition of a few keV will result in a measurable temperature increase of the crystal ($\sim 0.2$ mK/MeV). In CUORICINO the crystal temperature variations are measured by thermistors glued to each crystal that, in order to obtain usable signals for such small temperature changes, have to be very sensitive.

The thermistors are heavily doped high-resistance germanium semiconductors with an impurity concentration slightly below the metal-insulator transition. High quality thermistors require a very homogeneous doping concentration: CUORICINO uses Neutron Transmutation Doped (NTD) germanium thermistors, this is achieved by means of uniform thermal neutron irradiation throughout the entire volume in a nuclear reactor. The electrical conductivity of these devices, which is due to variable range hopping (VHR) of the electrons, depends very sensitively on the temperature. The resistance varies with temperature according to $R = R_0 \exp\left(\frac{T_a}{T}\right)^\gamma$, where $R_0$ and $T_a$ depend on the doping concentration and $\gamma = 1/2$.

\section{Cuoricino Setup}

CUORICINO is an array of 62 crystals of TeO\textsubscript{2} with a total active mass of 40.7 kg, that corresponds to a mass of $^{130}$Te of $\sim$ 11 kg. The tower is located inside the cryostat situated in the Hall A of Laboratori Nazionali del Gran Sasso (LNGS) of INFN. CUORICINO’s 62 crystals are arranged in a tower made by 13 planes (Figure 1), 11 of them are filled with 4 cubes of 5 cm side while the other two with 9 crystals $3 \times 3 \times 6$ cm\textsuperscript{3} each. Four $3 \times 3 \times 6$ cm\textsuperscript{3} crystals are enriched, two of which in $^{128}$Te, 82.3 \% isotopic abundance, and the other two in $^{130}$Te, isotopic abundance of 75 \%.

All the materials composing the detector were selected to be low contaminated with radioactive isotopes. To avoid external vibrations to reach the detectors the tower is mechanically decoupled from the cryostat through a steel spring. In order to shield against the radioactive contaminants from the materials of the refrigerator, a 1.2 cm shield of Roman lead with $^{210}$Pb activity less than mBq/kg is framed around the array to reduce the activity of the thermal shields. The cryostat is externally shielded by means of two layers of lead of 10 cm minimal thickness each. The background due to environmental neutrons is reduced by a layer of Borated Polyethylene of 10 cm minimum thickness. The refrigerator operates inside a Plexiglas anti-radon box flushed with clean N\textsubscript{2} and inside a Faraday cage to reduce electromagnetic interferences. CUORICINO is operated at a temperature of $\sim$ 8 mK with a spread of $\sim$ 1 mK. The energy calibration is performed before and after each subset of runs, which lasts about a month, by exposing the array to two thoriated tungsten wires inserted in immediate vicinity of the refrigerator.
4 Cuoricino Results

CUORICINO first measurement started in March 2003 and ended in October 2003. After a substantial operation of maintenance in April 2004 the second run of CUORICINO started. The average resolution FWHM is $7.5 \pm 2.9$ keV for the bigger size and of $9.6 \pm 3.5$ keV for the small size crystals. The duty cycle of the experiment, since August 2004 is $\sim 73\%$. Discarding the time needed for energy calibration measurement (3 days every 3–4 weeks) the total live time is $63\%$. The background spectra collected up to Aug 2007, corresponding to a total statistic of $15.53$ kg ($^{130}$Te)·year, is presented in figure 2. Apart the $^{60}$Co sum line, no other unexpected peak is found near the $2530$ keV $0$\nu DBD region of $^{130}$Te. The background level is $0.18 \pm 0.01$ c/keV/Kg/y and the corresponding lower limit on the $0$\nu DBD of $^{130}$Te is $3.1 \times 10^{24}$ y (90\% C.L.). This limit leads to a constraint on the electron neutrino effective Majorana mass ranging from $0.20$ to $0.68$ eV, depending on the nuclear matrix elements considered in the computation.

5 The Cuore Experiment

The CUORE detector will consist of an array of 988 TeO$_2$ bolometers arranged in a cylindrical configuration of 19 towers containing 52 crystals each (Figure 3), for a total mass of $\sim$741 kg. Each of these towers is a CUORICINO-like detector consisting of 13 modules, 4 detectors each. Assuming a background of $B=0.01$ c/keV/kg/y, achievable with a slight improvement of the current available material selection and cleaning techniques, and an energy resolution $\Gamma(2.5$ MeV)$=5$ keV, we get a sensitivity $S_{0\nu}$ on the half life (90 \% C.L.) of $5.8 \cdot 10^{25} \sqrt{t}$ years ($4.1 \cdot 10^{25} \sqrt{t}$ years for $\Gamma=10$ keV), which in 5 years of statistics would provide $m_{\beta\beta}$ bounds in the range $0.024$–$0.13$ eV. However, the R&D to be carried out in CUORE, if successful, would provide a value of $B\sim 0.001$ c/keV/kg/y, i.e. a detection sensitivity of $S_{0\nu} \sim 1.86 \cdot 10^{26} \sqrt{t}$ years.
Figure 2: CUORICINO spectrum in the $0\nu\beta\beta$ region

(1.2 $10^{26}\sqrt{t}$ years for $\Gamma=10$ keV), or $m_{33}$ bounds in the range $\sim 0.016-0.085$ eV in 5 years. TeO$_2$ crystals made with $^{130}$Te enriched material have been already operated in MiDBD and CUORICINO, making an enriched CUORE a feasible option. Assuming a 95% enrichment in $^{130}$Te and a background level of $B=0.001$ c/keV/kg/\text{y}, the sensitivity would become $S_{90}\sim 8.3\times10^{20}/\sqrt{t}$ years. For an exposure of 5 years, the corresponding $m_{33}$ bounds would range from 8 meV to 45 meV depending on the nuclear matrix element calculations.

References

Figure 3: CUORE detector: the bolometers array is made of 19 CUORICINO like towers
SEARCH FOR NEUTRINOLESS DOUBLE BETA DECAY IN $^{153}$Nd WITH THE NEMO3 EXPERIMENT

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The double beta decay experiment NEMO3 has been taking data since February 2003. The aim of the experiment is to search for neutrinoless double beta decay with 16 kg of enriched isotopes. Using 939 days of data, a preliminary result was obtained for $^{153}$Nd: $T_{1/2}^{0\nu} = (9.20_{-0.37}^{+0.36}\text{(stat.)} \pm 0.62\text{(syst.)}) \times 10^{19}$ y. No neutrinoless double beta decay signal was observed and a limit on the half-life of the process was set to $T_{1/2}^{0\nu} > 1.8 \times 10^{22}$ y at 90% confidence level. An overview of the results previously obtained for other isotopes is also given.

1 Introduction

Neutrinoless double beta decay ($0\nu\beta\beta$) for a nucleus of atomic number $A$ and charge $Z$ is the process: $(A, Z) \rightarrow (A, Z + 2) + 2e^-$. The process belongs to physics beyond the Standard Model as it violates the conservation of the total lepton number. Its observation would prove that the neutrino is a Majorana particle. Neutrinoless double beta decay may occur through several mechanisms, among which the decay with light neutrino exchange that would grant access to the mass scale. The half-life of the $0\nu\beta\beta$ process is then given by:

$$ (T_{1/2}^{0\nu})^{-1} = |M_{0\nu}(A, Z)|^2 G^{\nu}(Q, Z) \langle m_\nu \rangle^2 $$  \hspace{1cm} (1)

with $M_{0\nu}(A, Z)$ a nuclear matrix element (NME) obtained from theoretical calculations, $G^{\nu}(Q, Z)$ a phase space factor depending on the Q value of the process, and $\langle m_\nu \rangle$ the effective mass of the neutrino. This effective mass is given by $\langle m_\nu \rangle = \sum_{i=1}^{4} U_{ei}^2 m_i$ where $U_{ei}$ stands for the squared elements of the PMNS matrix and $m_i$ the mass associated to the mass eigenstate $i$.

2 The NEMO3 experiment

2.1 The NEMO3 detector

The NEMO3 experiment has been taking data since February 2003. The detector is located in the Modane Underground Laboratory (LSM) under a rock cover of 4800 m water equivalent. It accommodates 10 kg of double beta emitter foils. With its cylindrical geometry, the NEMO3 detector is divided into 20 sectors in the middle of which the foils were installed. The main isotopes used for the search of the neutrinoless double beta decay are $^{100}$Mo ($\sim 7$ kg) and $^{82}$Se ($\sim 1$ kg). Smaller amount of other isotopes are also used to study the two-neutrino double beta
Table 1: Isotopes in the NEMO3 detector.

<table>
<thead>
<tr>
<th>Mass (g)</th>
<th>69.14</th>
<th>93.2</th>
<th>405</th>
<th>36.6</th>
<th>9.4</th>
<th>7.0</th>
<th>454</th>
<th>491</th>
<th>621</th>
</tr>
</thead>
<tbody>
<tr>
<td>Q_{BS} (keV)</td>
<td>3034</td>
<td>2995</td>
<td>2805</td>
<td>3367</td>
<td>3350</td>
<td>4772</td>
<td>2529</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

decay ($2v\beta\beta$) process. The sectors containing the natural tellurium and the copper foils are dedicated to the background measurement. The mass of each isotope within the detector is summarized in Table 1.

The principle of the NEMO3 detector is the identification of the electrons in a double beta decay process and the measurement of their individual energy. Therefore, a tracking chamber is associated to a calorimeter. The tracking volume consists of 6180 drift cells operated in Geiger mode. It allows the reconstruction of the trajectory of charged particles and the determination of the position of the vertex with a resolution of 5 mm transversely to the source foil plane and 8 mm longitudinally. A 25 Gauss magnetic field is generated by a coil surrounding the detector and allows the identification of the charge of the particles. The drift gas inside the tracking volume is helium with added ethyl alcohol (4%), argon (1%) and water (0.1%). The calorimeter consists of 1940 photomultiplier tubes associated to plastic scintillator blocks. The energy resolution ranges from 14.1% to 17.7% for 1 MeV electrons. The time resolution is 250 ps. The tracker - calorimeter association allows to identify electrons, positrons, gammas and alpha particles as well as to measure their time of flight. The detector is protected from gammas by an iron shielding, and from neutrons present inside the laboratory by wood and a borated water shielding. A tent coupled to a radon-free air factory surrounds the detector. A detailed description of the detector was published in 1.

2.2 Background model

The background sources can be classified into three groups: the external background from incoming $\gamma$, the radon present inside the tracker volume and the internal background from radioactive contamination of the sources. The activities of these background sources can be obtained through the use of a set of control channels corresponding to different event topologies on the NEMO3 data. These measured activities were compared to measurements performed with HPGe and radon detectors.

At the beginning of the experiment, the radon inside the tracking chamber, and more precisely the decay product $^{214}$Bi present in its radioactive chain, was found to be the predominant background source. Radon is present in the air of the laboratory as it originates from the rock surrounding it and can penetrate inside the detector through small leaks. A tent coupled to radon-free air factory was then installed around the detector in October 2004 in order to decrease the presence of radon inside the tracker. The data taken before that date are referred to as Phase I data. After October 2004, the data are Phase II data.

3 NEMO3 results

3.1 $2v\beta\beta$ decay of $^{150}$Nd

Measurements of the $2v\beta\beta$ half-life can be performed for 7 isotopes in the NEMO3 experiment. This process is the following: $(A,Z) \rightarrow (A,Z+2) + 2e^- + 2\nu_e$. It constitutes the ultimate background in the $0v\beta\beta$ search because of the energy resolution of the detector.

The $\beta\beta$ type events are selected by requiring two tracks with a curvature compatible with a negative charge. Each track is associated to a separate energy deposit in the calorimeter greater
than 200 keV. Both track should originate from a common vertex located inside the source foil. The time of flight measurement for both electrons must be consistent with the hypothesis that they were emitted from the foil.

A preliminary measurement of the half-life of $^{155}$Nd was obtained for a 36.6 g sample from data collected between February 2003 and December 2006, corresponding to 939 days of data taken during the Phases I and II of the experiment. A total of 2828 $\beta\beta$ type events were observed with a signal-over-background ratio of 2.7. The distribution of the energy sum of the electrons in $\beta\beta$ type events and the angle between them are shown in Fig. 1. The background subtracted data and the $2\nu\beta\beta$ signal expectation obtained from Monte-Carlo are in good agreement. The $2\nu\beta\beta$ selection efficiency is 7.2%. The measured half-life is $T_{1/2}^{2\nu} = (9.20^{+0.37}_{-0.22}(\text{stat.}) \pm 0.62(\text{syst.})) \times 10^{18}$ y. This value is in between two previous results obtained from experiments led using time projection chambers $^2$, $^3$.

![Figure 1: Distribution of the energy sum of the two electrons $E_{\text{sum}}$ (left graph) and the angle between the two electrons $\cos\theta$ (right graph), background subtracted. The data are shown as points and the subtracted background as a red histogram.](image)

### 3.2 $2\nu\beta\beta$ decay of other isotopes

The half-life of the $2\nu\beta\beta$ process for the other isotopes inside NEMO3 was measured using the data from Phase I or combined Phases I and II and the preliminary results are given in Table 2. These results are important for nuclear theory as they help constrain the nuclear models and thus improve the NME calculations. The NME are a source of uncertainty when translating the half-lives of the $0\nu\beta\beta$ into effective neutrino masses.

### 3.3 Search for neutrinoless double beta decay of $^{155}$Nd

A $0\nu\beta\beta$ decay signal would correspond to an excess of $\beta\beta$ type events around the energy of the transition $Q_{\beta\beta}$ in the distribution of the energy sum of the electrons. Indeed, the theoretical peak at $Q_{\beta\beta}$ would be smeared out by the energy resolution of the calorimeter. For $^{155}$Nd, the $Q_{\beta\beta}$ value is 3.367 MeV. Fig. 2 shows that no excess of events in the distribution of the energy sum of the electrons for the $\beta\beta$ type events originating from the $^{155}$Nd sample is observed for 939 days of data collection. A limit on the half-life of the $0\nu\beta\beta$ process was subsequently set using the $CL_s$ method $^5$ for $E_{\text{sum}} > 2.5$ MeV. The corresponding $0\nu\beta\beta$ selection efficiency is 19%. The limit on the half-life obtained is $T_{1/2}^{0\nu} > 1.8 \times 10^{22}$ years at 90% confidence level (CL), which translates into an upper limit on the effective Majorana mass of the neutrino in the range...
Table 2: Half-lives of $2\nu\beta\beta$ processes measured using Phase I data (360 days). The $^{130}$Te results use Phases I and II data (534 days).

<table>
<thead>
<tr>
<th>Isotope</th>
<th>Signal/Background</th>
<th>$T_{1/2}$ [10$^{16}$ years]</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{100}$Mo</td>
<td>40</td>
<td>$0.711 \pm 0.002$(stat.) $\pm 0.054$(syst.)$^4$</td>
</tr>
<tr>
<td>$^{82}$Se</td>
<td>4</td>
<td>$9.6 \pm 0.3$(stat.) $\pm 1.0$(syst.)$^4$</td>
</tr>
<tr>
<td>$^{110}$Cd</td>
<td>7.5</td>
<td>$2.8 \pm 0.1$(stat.) $\pm 0.3$(syst.)</td>
</tr>
<tr>
<td>$^{95}$Zr</td>
<td>1</td>
<td>$2.0 \pm 0.3$(stat.) $\pm 0.2$(syst.)</td>
</tr>
<tr>
<td>$^{48}$Ca</td>
<td>~ 10</td>
<td>$3.9 \pm 0.7$(stat.) $\pm 0.6$(syst.)</td>
</tr>
<tr>
<td>$^{130}$Te</td>
<td>0.25</td>
<td>$76 \pm 15$(stat.) $\pm 8$(syst.)</td>
</tr>
</tbody>
</table>

$\langle m_\nu \rangle < 1.9 - 2.7$ eV according to the NME calculations in$^7$ and $\langle m_\nu \rangle < 5.4 - 8.5$ eV according to$^8$. The limit on the half-life was improved by one order of magnitude compared to the previous result, $T_{1/2}^{0\nu} > 1.7 \times 10^{21}$ years at 90% CL$^6$.

Along with the light neutrino exchange, other mechanisms can mediate $0\nu\beta\beta$ decay. Among these can be found the models with Majoron emission characterized by a spectral index $n$ and the model with a right-handed ($V+A$) contribution in the Lagrangian. In the assumption of a $0\nu\beta\beta$ process involving right currents ($V+A$), the limit on the half-life was found to be $T_{1/2}^{0\nu} > 1.27 \times 10^{22}$ years at 90% confidence level. For a $0\nu\beta\beta$ process with the emission of a Majoron (spectral index $n = 1$) the limit obtained is $T_{1/2}^{0\nu} > 1.55 \times 10^{21}$ years at 95% CL.

![NEMO 3 Preliminary](image)

Figure 2: Distribution of the energy sum of the two electrons $E_{sum}$ for $E_{sum} > 2.5$ MeV for data (points) compared to the total background (in red) consisting of the radioactive background (in grey) plus the $2\nu\beta\beta$. A MC simulation of a $0\nu\beta\beta$ signal with a half-life of $1.45 \times 10^{22}$ years is shown in blue.

### 3.4 Search for $0\nu\beta\beta$ of $^{100}$Mo and $^{82}$Se

$^{100}$Mo and $^{82}$Se are the main isotopes in the NEMO3 detector used for the search for neutrinoless double beta decay. For both isotopes, 639 days of data spread over the two phases of the experiment were analyzed. For $^{100}$Mo the $Q_{\beta\beta}$ value is 3.034 MeV. No excess of events was found in the $2.8 < E_{sum} < 3.2$ MeV energy window, as 12.1 events are expected from Monte Carlo and 11 were observed. The $0\nu\beta\beta$ selection efficiency is 8.2%, which leads to a limit on the
Table 3: Constraints on $T_{1/2}$ in years (90% CL) for (V+A) and Majoron emission processes from NEMO3 data. 
\( \lambda \) is a (V+A) Lagrangian parameter and \( g \) is a Majoron to neutrino coupling strength.

<table>
<thead>
<tr>
<th>0νββ process</th>
<th>$^{100}$Mo</th>
<th>$^{82}$Se</th>
</tr>
</thead>
<tbody>
<tr>
<td>(V+A) current</td>
<td>( &gt; 3.2 \times 10^{24} ) ( \lambda &lt; 1.8 \times 10^{-6} )</td>
<td>( &gt; 1.2 \times 10^{24} ) ( \lambda &lt; 2.8 \times 10^{-6} )</td>
</tr>
<tr>
<td>( n = 1 )</td>
<td>( &gt; 2.7 \times 10^{22} ) ( g &lt; (0.4 - 1.8) \times 10^{-4} )</td>
<td>( &gt; 1.5 \times 10^{22} ) ( g &lt; (0.7 - 1.9) \times 10^{-4} )</td>
</tr>
<tr>
<td>( n = 3 )</td>
<td>( &gt; 1.7 \times 10^{22} ) ( g &gt; 6.0 \times 10^{-4} )</td>
<td>( &gt; 6.0 \times 10^{-4} )</td>
</tr>
<tr>
<td>( n = 5 )</td>
<td>( &gt; 1.0 \times 10^{22} ) ( g &gt; 3.1 \times 10^{-4} )</td>
<td>( &gt; 3.1 \times 10^{-4} )</td>
</tr>
<tr>
<td>( n = 7 )</td>
<td>( &gt; 7.0 \times 10^{19} ) ( g &gt; 5.0 \times 10^{-4} )</td>
<td>( &gt; 5.0 \times 10^{-4} )</td>
</tr>
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</table>

half-life of the 0νββ process $T_{1/2}^{0\nu} > 5.8 \times 10^{23}$ years at 90% CL. This translates into an upper limit on the neutrino mass of \( \left< m_\nu \right> < 0.8 - 1.3 \text{ eV} \) using the values of the NME in 9. The $Q_{\beta\beta}$ value for $^{82}$Se is 2.995 MeV. As for $^{100}$Mo, no excess of events was found around the $Q_{\beta\beta}$ value in the distribution of the energy sum of the electrons. In the 2.65 < $E_{\text{sum}}$ < 3.20 MeV energy window, 7 events were observed for 6.4 estimated by Monte Carlo simulations. With a 14.4% 0νββ selection efficiency, the obtained limit on the half-life is $T_{1/2}^{0\nu} > 2.1 \times 10^{23}$ years at 90% CL. The corresponding upper limit on the mass of the neutrino using 10 is then \( \left< m_\nu \right> < 1.4 - 2.2 \text{ eV} \). The results obtained for the half-life of the 0νββ process of $^{100}$Mo and $^{82}$Se through the (V+A) and the Majoron mechanisms 11 are summarized in Table 3.

4 Summary

The NEMO3 experiment allows the measurement of 2νββ decays with very high statistics. A preliminary result was obtained for the 2νββ half-life of $^{150}$Nd for a 36.6 g sample: $T_{1/2}^{2\nu} = (9.20^{+0.25}_{-0.22}(\text{stat.}) \pm 0.62(\text{syst.})) \times 10^{18}$ years. A limit on the the 0νββ half-life of $T_{1/2}^{0\nu} > 1.8 \times 10^{22}$ years was estimated at 90% confidence level. The $Q_{\beta\beta}$ value of $^{150}$Nd is one of the highest among ββ emitters and lies above the typical energies of many background sources. Also, that isotope has a large phase space factor. These characteristics make $^{150}$Nd a promising candidate for the SuperNEMO experiment aiming at a sensitivity of 50 meV on the Majorana neutrino mass with 100 kg of enriched isotopes. The enrichment process is currently under study by the Nd-150 collaboration for the SuperNEMO and SNO++ projects.

References


LHC TEST OF THE SEE-SAW

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We discuss the prospects for detecting right-handed neutrinos which are introduced in the see-saw mechanism at future colliders. This requires a very accurate cancellation between contributions from different right-handed neutrinos to the light neutrino mass matrix. We search for possible symmetries behind this cancellation and find that they have to include lepton number conservation. Light neutrino masses can be generated as a result of small symmetry-breaking perturbations. The impact of these perturbations on LHC physics is negligible, so that the mechanism of neutrino mass generation and LHC physics are decoupled in general. In constrained cases, accelerator observables and neutrino masses and mixings can be correlated.

1 Introduction

The (type-I) see-saw mechanism generates small neutrino masses in a natural way, introducing right-handed (RH) neutrinos that are singlets under the Standard Model (SM) gauge group and can therefore have large Majorana masses. The light neutrino mass matrix is approximately given by

\[ m_\nu = -m_D m_R^{-1} m_D^T, \]

where \( m_D \) is the Dirac mass matrix and \( m_R \) is the Majorana mass matrix of the heavy singlets. A direct test of the see-saw mechanism requires the detection of these heavy neutrinos and the measurement of their Yukawa couplings. Using Eq. (1) in the case of only one generation and \( m_\nu \sim 0.1 \text{ eV} \), we obtain the estimate \( m_R \sim 10^{14} \text{ GeV} \), if the Dirac neutrino masses are close to the electroweak scale. The singlets may have masses as small as 100 GeV, within the energy reach of the LHC and other future colliders, if the Dirac masses are a bit smaller than the electron mass, which does not appear completely unreasonable either. However, the RH

\[ ^{a}\text{Talk presented at the XLIIIrd Rencontres de Moriond, March 1–8, 2008, La Thuile, Italy. Based on work done in collaboration with Alexei Smirnov.} \]
neutrinos interact with the SM particles only via Yukawa couplings,\(^b\) which are tiny in this case. Thus, we expect the RH neutrinos to be either way too heavy or way too weakly coupled to be observable at colliders.

However, this conclusion can be avoided provided that there are two or more RH neutrinos \(7,8,9,10,11,12,13,14,15,16,17,18,19,1,20\). Their contributions to the light neutrino masses can cancel, opening up the possibility of rather light singlets with large Yukawa couplings but exactly vanishing light neutrino masses. Non-vanishing masses are generated by small perturbations of the cancellation structure. In this setup, the RH neutrinos may be observable in future collider experiments. This possibility has attracted renewed interest recently, see e.g. \(21,22,23,24,25,26,20\).

In the following, we will discuss the prospects for discovering RH neutrinos at colliders from the point of view of theory. We will consider the cancellation of contributions to the light neutrino mass matrix and possible underlying symmetries in the next section. After briefly discussing small perturbations of the leading-order mass matrices that yield viable masses for the light neutrinos, we will turn to consequences for signatures at colliders. Within the setups relying on a symmetry, lepton number violation is unobservable, while lepton-flavour-violating processes can have sizable amplitudes. Finally, we will comment on the implications a detection of RH neutrinos would have for our understanding of the mechanism of neutrino mass generation.

2 Cancellations and Symmetries

2.1 Vanishing Light Masses

For three generations of left- and right-handed neutrinos, the contributions of the RH neutrinos to the light mass matrix cancel exactly, if and only if \(^{10,13,14,1,27}\) the Dirac mass matrix has rank 1,

\[
m_D = m \begin{pmatrix} y_1 & y_2 & y_3 \\ \alpha y_1 & \alpha y_2 & \alpha y_3 \\ \beta y_1 & \beta y_2 & \beta y_3 \end{pmatrix},
\]

and if

\[
\frac{y_1^2}{M_1} + \frac{y_2^2}{M_2} + \frac{y_3^2}{M_3} = 0 ,
\]

where \(M_i\) are the singlet masses. The mass parameters are defined in the basis where the singlet mass matrix is diagonal. The case of two RH neutrinos is analogous\(^{9,11,12}\), while for four or more RH neutrinos there are additional possibilities. The cancellation is valid to all orders in \(m_Dm_R^{-1}\). The overall scale of the Yukawa couplings is not restricted by the cancellation condition (3) and hence allowed to be large enough to make the detection of RH neutrinos possible. The only relevant constraint is the experimental bound on the mixing

\[
V = m_Dm_R^{-1}
\]

between active and singlet neutrinos,\(^{28}\)

\[
\sum_i |V_{\alpha i}|^2 \lesssim 10^{-2} \quad (\alpha = e, \mu, \tau)
\]

2.2 Underlying Symmetries

Without a symmetry motivation, the cancellation condition (3) amounts to severe fine-tuning and is unstable against radiative corrections. Let us therefore discuss symmetries leading to the

\(^b\)This is the case in the minimal extension of the SM we consider here. Of course, the situation is very different if the RH neutrinos have additional interactions, for example with TeV-scale SU(2)\(_R\) gauge bosons.
cancellation. We will restrict ourselves to the case of three singlets. A well-known possibility is imposing lepton number conservation\textsuperscript{7,8,15,17,18}. The assignment $L(\nu_L) = L(\nu_1^R) = -L(\nu_2^R) = 1$, $L(\nu_3^R) = 0$ implies

$$m_R = \begin{pmatrix} 0 & M & 0 \\ M & 0 & 0 \\ 0 & 0 & M_3 \end{pmatrix}, \quad m_D = m \begin{pmatrix} a & 0 & 0 \\ b & 0 & 0 \\ c & 0 & 0 \end{pmatrix}. \quad (6)$$

Two singlets form a Dirac neutrino with mass $M$, while the third one decouples.

An important question is whether lepton number conservation is also a necessary condition for the cancellation of light neutrino masses, i.e. whether the cancellation can result from a symmetry that does not contain $L$ conservation. One can show that there is always a conserved lepton number, if the cancellation occurs and if all three singlets have equal masses\textsuperscript{1}. Let us therefore consider the case where the singlets involved in the cancellation, say $\nu_1^R$ and $\nu_2^R$, have different masses and where the condition (3) is imposed by a symmetry at the energy scale $M_2$. Below this scale, the symmetry is broken. The neutrino masses change due to the renormalisation group running. The contributions from the two singlets to $m_\nu$ run differently between $M_1$ and $M_2$ in the SM\textsuperscript{29}, so that the cancellation is destroyed. A rough estimate yields

$$m_\nu(M_1) \sim 10^{-4} \text{ GeV} \ln \frac{M_2}{M_1} \quad (7)$$

at $M_1$, which is unacceptable unless $\nu_1^R$ and $\nu_2^R$ are degenerate. Of course, this problem persists if also the third singlet contributes to the cancellation.

Thus, the cancellation of light neutrino masses can only be realised without fine-tuning, if the RH neutrinos involved in the cancellation have equal masses, which implies lepton number conservation. Therefore, any symmetry leading to vanishing neutrino masses via this cancellation has to contain the corresponding U(1)$_L$ as a subgroup or accidental symmetry.

2.3 Small Perturbations

Non-zero masses for the light neutrinos are obtained by introducing small lepton-number-violating entries in the mass matrices (6). In the most general case,

$$m_R = \begin{pmatrix} \epsilon_1 M & M & \epsilon_{13} M \\ M & \epsilon_2 M & \epsilon_{23} M \\ \epsilon_{13} M & \epsilon_{23} M & M_3 \end{pmatrix}, \quad m_D = m \begin{pmatrix} a & \delta_a & \epsilon_a \\ b & \delta_b & \epsilon_b \\ c & \delta_c & \epsilon_c \end{pmatrix}. \quad (8)$$

The smallness of the observed neutrino masses leads to the restriction

$$\epsilon_2, \delta_{a,b,c} \lesssim 10^{-10} \quad (9)$$

for $\max(a, b, c) \sim 1$, $m/M \sim 0.1$, $M \sim 100$ GeV (as required by observability of RH neutrinos at LHC\textsuperscript{21,22,24,26}), provided that there are no special relations between the small parameters causing additional cancellations. The perturbations $\epsilon_{23}$ and $\epsilon_{a,b,c}$ appear quadratically in $m_\nu$ and are correspondingly less severely constrained. Finally, $\epsilon_1$ and $\epsilon_{13}$ do not lead to neutrino masses at the tree level at all but do contribute via loop diagrams\textsuperscript{12}, so that they are only slightly less constrained than the other parameters.

The most general mass matrices of Eq. (8) contain many free parameters, so that there is no clear imprint of the considered setup in the light neutrino mass matrix. A more interesting phenomenology is possible in constrained cases, some of which have been considered earlier in the context of leptogenesis\textsuperscript{30,18}. For example, if all small parameters are of the same order of magnitude,

$$m_\nu \approx \frac{m^2}{M} \left[ \epsilon_2 vv^T - (v v_3^T + v_3 v^T) \right], \quad (10)$$
where we have abbreviated the first and second column of $m_D$ by $v$ and $v_\delta$, respectively. The light neutrino masses are strongly hierarchical, since $m_\nu$ has rank 2 and hence one vanishing eigenvalue. The large Yukawa couplings $a, b, c$ are determined by the light neutrino masses and mixing parameters, which leads to predictions for correlations between the branching ratios of different lepton-flavour-violating decays in supersymmetric see-saw models. Likewise, the amplitudes of LFV processes at colliders are correlated, as we will discuss shortly.

### 3 Signals at Colliders

A striking signature of RH neutrinos at colliders would be lepton-number-violating (LNV) processes with like-sign charged leptons in the final state. However, we have argued that all symmetries guaranteeing the required suppression of the light neutrino masses lead to the conservation of lepton number, so that the amplitudes of such processes vanish. Any $L$ violation is severely restricted by the smallness of neutrino masses and can therefore not lead to sizable amplitudes. Consequently, in the absence of fine-tuning, LNV signals are expected to be unobservable.

Another option are events with different leptons such as $\mu^-\tau^+$ in the final state, since these have a relatively small SM background as well. Such signals are unlikely to be observable at LHC, however. In the considered scenarios, the mechanism leading to the cancellation of neutrino masses causes the terms in the corresponding amplitudes to add up constructively, leading to

$$A_{\alpha\beta} \propto \frac{m_\nu^2}{M^2} (a, b, c)_\alpha (a^*, b^*, c^*)_\beta$$

for the mass matrices of Eq. (8), where $\alpha \neq \beta$ denote the flavours of the charged leptons. If the cross sections are large enough for a detection at colliders, flavour-violating decays of charged leptons mediated by the RH neutrinos should be observable in upcoming experiments as well, since their amplitudes depend on the same combination of parameters. In the constrained case that yields Eq. (10), $a, b, c$ can be determined from the light neutrino mass parameters, as mentioned above, so that the ratios $A_{e\mu}/A_{e\tau}$ and $A_{e\mu}/A_{\mu\tau}$ are predicted.

At the ILC, the resonant production of RH neutrinos is possible for $|V_{ei}| \gtrsim 0.01$. By observing the branching ratios for the subsequent decays into charged leptons, one could then determine the mixings of the heavy neutrinos with the different left-handed doublets directly.

### 4 Summary and Discussion

We have discussed the prospects for testing the see-saw mechanism of neutrino mass generation in collider experiments. We have assumed the existence of right-handed neutrinos with masses close to the electroweak scale (but no other new particles or interactions). The couplings of these neutrinos to the SM particles can only be large enough to make their observation at colliders possible, if different contributions to the light neutrino masses nearly cancel. This cancellation is then the main reason for the smallness of the observed neutrino masses, while the see-saw mechanism plays only a minor role. Therefore, we have to conclude that a direct test of the see-saw mechanism at the LHC or the ILC is not possible.

If one defines the leading-order mass matrices in such a way that they correspond to exactly vanishing light neutrino masses, non-zero masses appear as a result of small perturbations of this structure. One may then ask whether these perturbations could have consequences for signals at colliders and thus allow for a test of the mechanism of neutrino mass generation. Unfortunately, the smallness of the light neutrino masses immediately tells us that all perturbations are tiny and therefore irrelevant for collider signatures. Thus, the answer to this second question is
negative, too. Collider experiments are only sensitive to the leading-order mass matrices which do not lead to neutrino masses.

As a consequence, a connection between collider physics and neutrino masses can only be established, if the perturbations are introduced in such a way that the leading-order parameters are related to the light neutrino masses and mixings. In the most general case, this is not possible because there are too many free parameters. Then collider physics decouples completely from the light neutrino masses and their generation. However, the situation is better in constrained setups where only some of the perturbations are present or dominant. In the cases we discussed, a strong mass hierarchy is expected. To the extent that the leading-order Yukawa couplings are fixed by the measured neutrino masses and mixings, correlations between the branching ratios of lepton-flavour-violating processes can be obtained. This applies both to reactions at colliders and to LFV decays of charged leptons. Finally, $e^+e^-$ colliders may be able to determine the mixings of RH neutrinos with the different flavours directly. Pursuing all these experimental options provides a chance to test constrained setups of the kind we have described. Of course, even in this optimistic case it is impossible to exclude the existence of additional, very heavy RH neutrinos contributing to neutrino masses via the standard see-saw mechanism.

Without an underlying symmetry, the described cancellation of the light neutrino masses amounts to severe fine-tuning. We have therefore discussed symmetry motivations. We have argued that every symmetry realising the cancellation has to include lepton number conservation. Otherwise, the cancellation is unstable against radiative corrections, so that fine-tuning is still required. Thus, both lepton number violation and light neutrino masses arise due to small perturbations of the leading-order mass matrices, and their magnitudes are related. Therefore, we expect lepton-number-violating signals at colliders to be unobservable in untuned scenarios. The cross sections for lepton-flavour-violating processes are not suppressed, so that LHC experiments might be able to observe such reactions. If this is the case, lepton flavour violation should also be observable in decays of charged leptons in the near future.

For completeness, let us briefly consider different see-saw scenarios as well, where the particles responsible for generating neutrino masses are not gauge singlets. In such a case, they can be produced by gauge interactions. Consequently, large Yukawa couplings and thus the discussed cancellation of light neutrino masses are no longer required. Neither is it necessary to impose lepton number conservation in order to motivate this cancellation by a symmetry. Therefore, lepton-number-violating processes can be detectable via their signature of like-sign charged leptons. One example for such a scenario is left-right symmetry close to the TeV scale. Here the right-handed neutrinos can be produced via interactions with the new gauge bosons $W_R$ and $Z'$. In the type-II see-saw setup, where neutrino masses arise from the vacuum expectation value of a scalar triplet $\Delta$, the new particles can be produced in reactions like $q\bar{q} \rightarrow \gamma, Z \rightarrow \Delta^{++}\Delta^{--}$. Precise measurements of the decay rates $\Gamma(\Delta^{++} \rightarrow l^+_\alpha l^-_\beta)$ may even allow to probe the Majorana phases in the lepton mixing matrix. In the type-III see-saw mechanism, fermionic triplets $T$ are responsible for neutrino masses. Again, they may be detected by observing like-sign charged leptons, for instance in the process $q\bar{q} \rightarrow W^+ \rightarrow T^+T^0 \rightarrow l^+_\alpha l^-_\beta + \text{jets}$, where the couplings relevant for the triplet decays are related to the light neutrino masses.

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**References**

Neutrinos: astrophysics and astroparticle aspects
Neutrinos produced in gamma-ray bursters (GRBers) may provide a unique probe for the physics of these extreme astrophysical systems. Here we discuss neutrino production in inelastic neutron-proton collisions within the relativistic outflows associated with GRBers. We consider both the widely used fireball model and a recently proposed magneto-hydrodynamic (MHD) model for the GRB outflow.

1 Introduction

Gamma-ray bursts (GRBs) are short and energetic flashes of gamma rays (\(\sim 100\) keV), reaching Earth from apparently random direction at a rate of a few per day (see Ref.\(^1\) for a review). The luminosity of these bursts may be very large, sufficient to temporarily outshine all other gamma-ray sources combined. Following their accidental discovery in 1967, the origin of these remarkable events has puzzled astronomers for three decades. In particular, the question whether gamma-ray bursts were produced by sources within our galaxy or at cosmological distances has remained under debate until the 1990s.

In the last decade there has been tremendous progress in our understanding of GRBs. This is largely due to observations of GRB afterglows – periods of prolonged broad-band electromagnetic emission following the actual burst – that were first discovered in 1997. The cosmological distance scale has been established by redshift measurements of the afterglow, and in some cases afterglow observations have allowed for identification of the host galaxy, providing further clues as to the nature of gamma-ray bursters (GRBers). There is presently compelling evidence that long-duration GRBs (the subclass of GRBers lasting more than 2 sec) are ultimately caused by core-collapse of massive stars\(^2\), although the situation for short-duration GRBs (lasting less than 2 sec) is less clear.
One can easily estimate that the energy released in a stellar core-collapse matches that required to power a GRB, but the mechanism responsible for the energy transfer is far from obvious. The widely accepted framework to describe this is divided into four phases. In the initial phase, core-collapse of the massive star results in a black hole-accretion disk system. This launches a jet, a collimated outflow of plasma that contains a small baryonic component. In the accelerating phase, this plasma accelerates to a velocity close to that of light (Lorentz factor $\sim 300$). In this acceleration process, the initial energy of the plasma is transferred to bulk kinetic energy of the baryons that are contained in the plasma. In the coasting phase, the outflow moves with a fixed velocity through the pre-burst stellar environment. Here dissipation of the kinetic energy in the flow, most likely as synchrotron emission of shock-accelerated electrons, gives rise to the actual GRB. Finally, the afterglow is attributed to the interaction of the outflow with the external medium during the afterglow phase.

Although the framework described above successfully explains the general features of the observations, many questions remain. Arguably one of the most important issues is the nature of the relativistic outflow. Within the widely used fireball model, it is understood that the plasma is initially dominated by thermal energy. Alternatively, the energy may be predominantly in electromagnetic form. Such outflows are expected naturally when a magnetized accretion disk is surrounding the central black hole. Further questions concern for example the initial collimation of the flow, where magnetic fields may also play an important dynamical role, and the details of the energy dissipation process, which is likely to involve some particle acceleration mechanism such as shock acceleration.

Besides the intrinsic motivation to better understand the physics of GRBers, further motivation is provided by the connection to other fields of physics. Since GRBers are believed to be efficient astrophysical particle accelerators, they are candidate sources of high-energy neutrinos and cosmic rays and provide a laboratory to study the acceleration mechanism. Furthermore, it has been proposed to use GRBs as standard candles to constrain the evolution of the universe. Finally, there are more speculative proposals, e.g. to use the arrival times of low- and high-energy emission to constrain Lorentz violating interactions.

Neutrinos are promising probes of the environment of GRB sources. Neutrino emission is complementary to the electromagnetic emission in two respects. First, neutrinos mostly trace the hadronic component of GRB outflows whereas electromagnetic radiation mostly traces the leptonic component. Second, neutrinos can leave the GRB source when it is still optically thick. Substantial neutrino production may be expected in various phases of a developing GRB. In the initial phase, neutrino emission can constrain the formation of GRB fireballs. Within the fireball model, the dominant neutrino production process in this phase is electron-positron annihilation (providing a counterexample to the mostly hadronic production mechanisms). Under favorable circumstances, this may give rise to copious neutrino production. However, this mechanism is not sufficiently efficient to carry away the bulk of the fireball energy or to qualitatively modify the dynamical behavior of the fireball. In both the coasting and afterglow phases of a developing GRB, it is believed that kinetic energy is dissipated through shock acceleration of electrons. These shocks will likely also accelerate any protons contained in the fireball. Interactions of these high-energy protons with target nucleons or photons give rise to a flux of high-energy neutrinos that offers good detection prospects with the upcoming km$^3$ neutrino detectors such as IceCube. These neutrinos provide information on the nature of the flow, in particular the strength of the hadronic component, and on the energy dissipation process.

Here we report on a different mechanism to create neutrinos in GRBers, namely inelastic neutron-proton ($np$) collisions that occur during the accelerating phase. We compare the typical neutrino emission through this mechanism for two competing models: the fireball model and the recently proposed ‘AC’ model, which assumes that the energy in the outflow is predominantly electromagnetic. The motivation of this work is to estimate the detection prospects of this
neutrino emission and to investigate whether it can be used to differentiate between the fireball model and the AC model. The $np$ mechanism has been considered within the fireball model before.\textsuperscript{12,13} Our estimates are more pessimistic than existing ones, which can be traced to the more accurate modeling of the inelastic $np$ cross section adopted in our work. For the AC model, the mechanism was first considered in Ref.\textsuperscript{14}, which forms the basis of the present discussion.

In the following section we discuss the dynamics of GRB outflows containing neutrons and protons, both within the fireball model and within the AC model. We then discuss neutrino production through $np$ interactions, and finally we present our conclusions.

2 Dynamics

2.1 Acceleration in the fireball model and the AC model

A striking feature of GRB models is the bulk relativistic motion. This ingredient is motivated by an observational paradox: the short timescales and large energies suggest a huge energy density and thus an optically thick source. This then implies that the photon spectrum should be quasi-thermal, while observations show that it is not. Relativistic motion solves this problem by increasing the physical timescale compared to that inferred from observations, and by decreasing the photon energy in the source compared to the observed energy. The mechanism to accelerate the flow to relativistic velocities differs between models. In the fireball model, acceleration results from the pressure that photons exert on the optically thick fireball. In this case the dynamics of the flow may be approximated with\textsuperscript{15}

\begin{equation}
\Gamma \propto r,
\end{equation}

where $\Gamma$ is the Lorentz factor of the flow, and $r$ the radius of the flow (i.e., the distance from the central black hole). In the AC model, the energy to accelerate the outflow is provided by magnetic reconnection, a mechanism that converts electromagnetic energy into heat and bulk motion. When the magnetic field lines predominantly change polarity in the flow direction, as we will assume, the dynamics of the flow may be approximated with\textsuperscript{16}

\begin{equation}
\Gamma \propto r^{1/3}.
\end{equation}

Comparison with eq. (1) shows that the acceleration of the flow is much more gradual in the AC model than in the fireball model. As we will see, this directly affects the neutrino flux from $np$ collisions.

In both the fireball model and the AC model, acceleration of the flow stops when there is no more energy available to further accelerate the baryons. In the fireball model, the acceleration of the flow can also be terminated when the flow, whose energy density decreases with increasing radius, becomes optically thin.

2.2 Neutron-richness

Since the baryons that are contained in the flow are to be accelerated to high Lorentz factors, the initial baryon density cannot be too large. This requirement is generally stated in terms of a dimensionless baryon loading parameter

\begin{equation}
\eta \equiv L/\dot{M}c^2 \sim 10^3,
\end{equation}

where $L$ denotes the total luminosity of the flow and $\dot{M}$ the mass flux. Near the central black hole, the typical energy density is larger than nuclear binding energies so that the baryonic component will consist predominantly of free protons and neutrons. The ratio of neutrons to
protons at the base of the outflow is determined by the competition of electron capture on protons and positron capture on neutrons. Recent studies\textsuperscript{17} favor a neutron-rich environment, so that we expect that the outflow associated with a developing GRB is initially also neutron-rich. The neutron-to-proton ratio is parameterized with
\[ \xi \equiv \frac{\dot{M}_n}{\dot{M}_p} \sim 1, \]
where \( \dot{M}_n(p) \) denotes the neutron (proton) mass flux. At larger radii, where the energy densities are smaller, nucleosynthesizing reactions reduce the number of free neutrons. However, a significant amount of neutrons is expected in the flow up to the radius where neutron decay becomes important. This radius is much larger than the radii relevant to \( np \) collisions and thus neutron decay is not important for the mechanism considered in this work.

2.3 Neutron decoupling and pion production

Eqs. (1) and (2) are idealized approximations that are only valid when the baryons contained in the plasma play no dynamical role. Detailed numerical studies\textsuperscript{14,18} indicate that a reasonably strong baryonic component affects the dynamics. However, eqs. (1) and (2) provide a reasonable approximation to the full dynamical behavior that captures the properties which are essential to the particle production problem discussed here. We will thus neglect the dynamical importance of nucleons in this section.

Regardless of the mechanism that accelerates the flow, protons are strongly coupled to the other plasma components by electromagnetic interactions and follow the dynamics of the flow. The neutrons, on the other hand, are only coupled to the plasma through inelastic \( np \) collisions. The nucleon number densities are initially very large so that the \( np \) interaction timescale is much shorter than the dynamical timescale. In this regime, the neutrons and protons essentially behave as a single fluid. As the outflow expands, the number densities decrease and the scattering timescale increases. When the \( np \) scattering timescale becomes smaller than the dynamical timescale, the neutrons effectively decouple from the plasma and coast with a certain terminal velocity.

When the flow is still in the accelerating phase at \( np \) decoupling, the protons are accelerated further and consequently a bulk velocity difference develops between protons and neutrons. If this velocity becomes sufficiently large, pions can be created in inelastic \( np \) collisions. The threshold condition to produce pions may be expressed as \( \chi \equiv \Gamma_p/\Gamma_n > \chi_\pi \equiv 2.15 \), where \( \Gamma_p(n) \) denotes the proton (neutron) Lorentz factor. Approximating the dynamics of the outflow with \( \Gamma \propto r^p \) (where \( p = 1 \) corresponds to the fireball model and \( p = 1/3 \) to the AC model), we observe that the radius where pion production occurs \( r_\pi \) and the decoupling radius \( r_{np} \) are related through \( r_\pi \sim r_{np}\chi_\pi^{1/p} \). Hence, in the fireball model the pion production radius is roughly twice the decoupling radius, while in the AC model it is an order of magnitude larger.

If the outflow contains many baryons, the available amount of energy per baryon is relatively small. In this case the acceleration of the flow may saturate before \( np \) decoupling, thus preventing inelastic collisions. Hence a sufficiently ‘pure’ flow (\( \eta \gtrsim 500 \) for the fireball model, or \( \eta \gtrsim 200 \) for the AC model) is required for particle production in inelastic \( np \) collisions.

3 Particle production in neutron-proton collisions

3.1 Interaction probability

The probability \( d\tau \) for a neutron moving with dimensionless velocity \( \beta_n \) to interact with a proton population moving with dimensionless velocity \( \beta_p \), within an infinitesimal radius \( r \ldots r + dr \) is\textsuperscript{14}
\[ d\tau = \sigma \Gamma_p n'_p \left( \frac{\beta_p - \beta_n}{\beta_n} \right) dr \simeq \frac{\sigma n'_p}{2\Gamma_n} \left( \chi - 1 \right) dr, \]
where \( n'_p \) denotes the comoving proton density, \( \sigma \) is the inelastic \( np \) cross section\(^a\), and we have assumed that \( \beta_n \simeq 1 \) and \( \beta_p \simeq 1 \) in the second equality. For outflows that follow an acceleration profile \( \Gamma \propto r^p \) up to infinity, integrating eq. (5) gives the probability \( \tau \) for an inelastic \( np \) collision to occur somewhere between the pion production radius and infinity. The result is independent of any model parameters except the index \( p \). Performing this integral, we find that \( \tau \simeq 0.2 \) for the fireball model \( (p = 1) \) and \( \tau \simeq 0.008 \) for the AC model \( (p = 1/3) \). A comparison of these estimates with numerical results\(^{14}\) shows that the estimate on \( \tau \) is fairly accurate for the AC model over a large range of parameters. For the fireball model, however, this procedure tends to overestimate the optical depth. The reason for this is that, for a large range of model parameters, the flow becomes optically thin shortly after pion production becomes possible. This prevents further acceleration of the flow. Hence the acceleration profile \( \Gamma \propto r \) does not hold up to large radii and the above estimate is not very accurate. Numerical results indicate that a typical value for the fireball model is \( \tau_{FB} \simeq 0.05 \), while for the AC model \( \tau_{AC} \simeq 0.01 \).

Qualitatively, this difference can directly be understood from the dynamics: in the AC model, pion production is only possible at radii an order of magnitude larger than the \( np \) decoupling radius. This implies that the number density of target protons has decreased significantly since decoupling, leading to a small interaction probability. For the fireball model, pion production occurs closer to the decoupling radius, where the dilution of target protons is not so strong.

### 3.2 Neutrino emission

The neutrino fluence from a single GRB source at proper distance \( D_p \) can be expressed as

\[
\Phi_\nu \simeq 1.5 N_n \tau / 4 \pi D_p^2
\]

where \( N_n \) denotes the isotropic-equivalent number of neutrons in the flow, \( \tau \) is the \( np \) interaction probability, and we have taken the average number of neutrinos (adding flavors and antiparticles) per \( np \) scattering equal to 1.5\(^{14}\). Using \( N_n \simeq \xi_0 / (1 + \xi_0) \times E / (\eta m_n c^2) \), where \( \xi_0 \) is the initial neutron-to-proton flux ratio (cf. eq. (4)), \( E \) is the total isotropic-equivalent burst energy, \( \eta \) is the baryon loading parameter (cf. eq. (3)), and \( m_n \) the neutron mass, we find the following neutrino fluences for a burst at redshift \( z = 1 \) for the fireball model and the AC model, respectively:

\[
\Phi^{FB}_\nu \simeq 10^{-4} \text{ cm}^{-2} \left( \frac{\tau}{0.05} \right) \left( \frac{2 \xi_0}{1 + \xi_0} \right) \left( \frac{E}{10^{53} \text{ erg}} \right) \left( \frac{\eta}{10^3} \right)^{-1};
\]

\[
\Phi^{AC}_\nu \simeq 2 \times 10^{-5} \text{ cm}^{-2} \left( \frac{\tau}{0.01} \right) \left( \frac{2 \xi_0}{1 + \xi_0} \right) \left( \frac{E}{10^{53} \text{ erg}} \right) \left( \frac{\eta}{10^3} \right)^{-1}.
\]

Using the fact that pions are created near threshold, and assuming a roughly isotropic distribution in the center-of-mass frame, one finds that the typical observed neutrino energy is \( \sim 50 \) GeV for the fireball model and \( \sim 70 \) GeV for the AC model\(^{14}\). The typical energy for the AC model is slightly higher because charged pions will be accelerated by the plasma before decay.

For the fireball model, the flux estimate (6) is roughly an order of magnitude below previous estimates\(^{13}\). This difference can be attributed to a more accurate treatment of the \( np \) interaction (in Ref.\(^{13}\) it is assumed that \( \tau \simeq 1 \)). For the AC model, the interaction probability is smaller by another factor \( \sim 5 \). This difference results from the more gradual acceleration of the flow and is thus directly linked to its nature. Unfortunately, the detection prospects with the upcoming \( \text{km}^3 \) neutrino detectors such as IceCube are very poor due to the relatively low neutrino energy: for reference values of the parameters we expect less than 1 event per year for a combined, diffuse flux of 1000 GRBers per year for either model. This GRB rate is rather optimistic if one takes into account that \( np \) decoupling only occurs for sufficiently pure (high-\( \eta \)) GRBers. We thus conclude that realistic detection prospects for the neutrino flux studied here requires a detector with larger effective area at sub-100 GeV energies than the upcoming \( \text{km}^3 \) detectors.

\(^{a}\)We refer the reader to Ref.\(^{14}\) for the adopted approximation for \( \sigma \).
4 Discussion

Neutrino emission offers a promising way to further our understanding of gamma-ray bursters. Neutrinos carry information that is complementary to electromagnetic emission because they can escape from optically thick regions and because they predominantly trace the hadronic component of GRB sources. This offers a unique way to constrain the nature of the relativistic outflow associated with GRBs. However, due to their feeble interactions in detectors at Earth, it remains a challenging task to identify concrete realizations of this potential.

Here we have discussed neutrino production in inelastic neutron-proton collisions that occur when neutrons have decoupled from the outflow associated with GRBs. We have estimated the characteristic neutrino flux within the widely used fireball model and the more recently introduced AC model. The characteristic neutrino fluxes and energies are distinctively different for the two models, directly reflecting the dynamics and hence the nature of the flow. Unfortunately, the relatively low neutrino energy precludes any realistic detection prospects with the upcoming km$^3$ detectors such as IceCube.

Apart from neutrino production through charged pion decay, one also expects the production of gamma rays through the decay of neutral pions produced in $np$ interactions. The plasma is optically thick to these gamma rays, and hence they cannot directly leave the plasma. In fact, the energy that is injected in the flow through this mechanism is reprocessed (through synchrotron radiation, pair production, and Inverse Compton scattering) and emitted in a different energy band. The typical energy of this reprocessed emission is $\sim$10 GeV for the fireball model and $\sim$100 keV for the AC model, and the expected fluence is detectable up to large redshifts with the GLAST satellite.\textsuperscript{14} Detection of this emission would favor the fireball model, and constrain the baryon loading of the flow.

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HIGH ENERGY NEUTRINOS FROM THE COLD: STATUS AND PROSPECTS OF THE ICECUBE EXPERIMENT

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The primary motivation for building neutrino telescopes is to open the road for neutrino astronomy, and to offer another observational window for the study of cosmic ray origins. Other physics topics, such as the search for WIMPs, can also be developed with neutrino telescope. As of March 2008, the IceCube detector, with half of its strings deployed, is the world largest neutrino telescope taking data to date and it will reach its completion in 2011. Data taken with the growing detector are being analyzed. The results of some of these works are summarized here. AMANDA has been successfully integrated into IceCube data acquisition system and continues to accumulate data. Results obtained using only AMANDA data taken between the years 2000 and 2006 are also presented. The future of IceCube and the extensions in both low and high energy regions will finally be discussed in the last section.

1 Motivations for neutrino astronomy with IceCube

We expect the acceleration of cosmic rays in astrophysical objects to be accompanied by the production of high energy neutrinos via $pp$ or $p\gamma$ interactions at the acceleration site\(^1\). The detection of these neutrinos could provide us with fundamental information about these sources of cosmic rays, the most violent objects in the universe. The preferred candidates for astro-accelerators are expected to have large-scale strong shocks and/or strong magnetic field, such as active galactic nuclei (AGN), supernovae remnants, microquasars or gamma ray bursts. The study of such objects is on-going in gamma ray astronomy which has produced an impressive harvest of results in the last few years. Moreover, the recent announcement of the Auger collaboration of a possible correlation between the arrival direction of cosmic rays with energies in excess of $6\times10^{19}$ eV and AGN\(^2\) reinforces the interest in the studies on bottom-up processes for cosmic rays production and multi-messenger studies. For instance, if an AGN origin of the highest energy cosmic rays were to be confirmed, the neutrino flux from those objects could be within reach of a kilometer scale detector such as IceCube within a few years\(^3\).

Neutrinos are a completely unique tool for the study of the cosmic ray sources. Charged protons below $10^{18.5}$ eV are bent by the inter-galactic magnetic fields and no longer point back to their sources thus making proton astronomy impossible at low energies. As for neutral messengers, the gamma rays are strongly attenuated above 50 TeV because of their interaction with the infrared background, reaching us only from galactic sources at high energies. The neutrinos, with their weak interaction cross section, are the only particles that allow us to explore the non thermic universe at cosmological distances and at all energies. Nevertheless, if this very small cross section is an advantage for the propagation from the source, the detection of neutrinos from astronomical events requires a very large detection volume. A kilometer cube scale is
needed to allow the detection of a measurable number of events in one year from the expected diffuse neutrino flux at the Waxman-Bahcall bound. Several large scale neutrino telescopes are currently taking data or under development using either water (Baikal, ANTARES, NEMO, NESTOR and the future kilometer scale KM3Net) or ice as a detection medium (IceCube and its sub-detector AMANDA). The sensitivity of IceCube to astrophysical neutrino sources is given in a previous article. We will give here a brief overview of the results of AMANDA for neutrino astrophysics in section 4, and will present the first point source search with a partial configuration of IceCube in section 5. In addition to these analyses, IceCube will also look for high energy neutrinos that are expected to be created by interaction of the cosmic ray protons with the background radiation. More specific studies done with IceCube for supernovae detection and GRB searches can also be found in the literature.

2 Other physics potential of the experiment

In addition to high energy neutrino astronomy, IceCube’s scientific reach extends to particle physics by looking for neutrino from annihilation of weakly interacting massive particles (WIMP), like the neutralino. These dark matter candidates are expected to accumulate in gravitational potential wells such as the Earth or Sun. If the annihilation products of two trapped neutralinos include neutrinos, a neutrino flux excess may be observed in the direction of the center of the Earth or Sun. As a consequence, neutrino telescopes like IceCube, IceCube DeepCore (see section 6.2) and AMANDA aim at the indirect detection of dark matter with the Sun or the Earth as effective neutrino sources. The results obtained with AMANDA are presented in section 4.

Other topics of particle physics can also be addressed but will not be discussed here, like the search for magnetic monopoles, strange quark matter or SUSY Q-balls.

3 The IceCube detector

IceCube is located at the geographic South Pole. Neutrinos are detected through their charged current interactions in the ice of the detector volume (or in the ice surrounding the detector for muon and tau neutrinos). The Cherenkov light produced by the charged lepton resulting from this interaction travels through the transparent ice and is collected by the digital optical modules (DOMs) of IceCube.

The current IceCube design consists of 80 strings, each bearing 60 DOMs. They are deployed at depths between 1450 and 2450 m below the surface of the ice, forming the in-ice part of the detector. The DOMs have a spacing of 17 m on each string and the strings form a triangular grid pattern with an inter-string spacing of 125 m, providing a 1 km$^3$ instrumented volume. The buried detector is topped on the surface by an array of 80 stations called IceTop for the study of extensive air showers (see fig.1). Each IceTop station, located above an IceCube string, consists of two tanks filled with ice. Each of those tanks contains two DOMs of same design as the one used for the in-ice part of the detector. The surface array can be operated looking for anti-coincidence with the in-ice events to reject downward-going muons or in coincidence, to provide a useful tool for cosmic ray composition studies.

Each DOM used by IceCube comprises a 10$^7$ Hamamatsu R7081-02 photomultiplier tube (PMT) housed in a glass pressure vessel and in situ data acquisition electronics. This electronics is the heart of the IceCube data acquisition system: it reads out, digitizes, processes and buffers the signals from the PMT. When the individual trigger conditions are met at the DOM, it reports fully digitized waveforms to a software-based trigger and event builder on the surface. The electronics acquires in parallel on Analog Transient Waveform Digitizers (ATWDs) at 300 megasamples per second (MSPS) sampling over a 425 ns window. In addition the electronics
also records the signal with a coarser 40 MSPS sampling over a 6.4 \( \mu \text{s} \) window to catch the late part of the signals. Two parallel sets of ATWDs on each DOM operate in alternation so that one is active and ready to acquire while the other is read out. This design greatly reduces the dead-time of an individual DOM. The time calibration yields a timing resolution with a RMS narrower than 2ns for the signal sent by the DOM to the surface\(^7\). The noise rate due to random hits observed in-ice DOMs is of the order of 300 Hz. This very low value gives us the possibility to monitor the DOM hit rates and to use it to have a sensitivity to low energy (MeV) neutrinos from supernova core collapse throughout the Milky Way and out to the Large Magellanic Cloud\(^8\).

IceCube has also integrated its predecessor, the AMANDA detector, as it is now surrounded by IceCube (see fig. 1). AMANDA consists in 677 analog optical modules distributed on 19 strings with a much denser configuration than IceCube (string spacing of approximately 40m), giving it a lower energy threshold. The AMANDA optical modules are less sophisticated than the IceCube DOMs. The pulse processing electronics and data acquisition system is on the surface and the signal from AMANDA OMs has to be transmitted over roughly 1 km before being treated. Roughly half of the 677 AMANDA OMs transmit their signals to the surface over optical fibers, which allows for a timing accuracy of 2 to 3 ns, comparable to the one of the DOMs, although with greatly reduced dynamic range. The other half of the OMs are connected to the surface only by electrical cables, which stretch the pulses substantially thus separation of successive pulses is prevented. For relatively low energy events, the dense configuration of AMANDA gives it a considerable advantage over IceCube. Moreover, IceCube strings surrounding AMANDA can be used as an active veto against cosmic ray muons, making the combined IceCube + AMANDA detector considerably more effective for low energy studies than AMANDA alone. The DeepCore upgrade, whose construction will start next austral summer, will provide IceCube with a dense subdetector using DOM technology. This upgrade will open many possibilities in the low energy region and WIMPs studies as discussed in section 6.2.

The data taking with the partially finished IceCube detector is running smoothly and the detector is operating as expected. The detector began taking data in 2006 with a nine strings configuration (IC-9) and with a 22 strings configuration in 2007 (IC-22). The data acquisition with AMANDA also continues, enabling analyses done with more than 7 years of accumulated
data. The analysis of the IC-9 configuration of IceCube has already lead to first results with atmospheric neutrinos which are detailed in section 5. The analysis of IC-22 data is on-going and will be finished during the summer 2008.

4 Summary of current AMANDA results

4.1 Search for astrophysical sources

Between 2000 and 2004, AMANDA-II, the final configuration of the AMANDA detector as an independent entity, has been taking data. Results on the 5 years of the dataset have been reported. This subset yields 4282 up-going neutrino candidates with an estimated background contamination of approximately 5%. The analysis for point sources in the Northern hemisphere sky\(^ {18}\) for this dataset yielded no statistically significant point source of neutrinos as can be seen in Fig. 2. The highest positive deviation corresponds to about \(3.7\sigma\). The probability of such a deviation or higher due to background, estimated with 100 equivalent sky surveys of events with randomized right ascension, is 69%. Based on these studies, an upper limit has been placed on a reference \(E^{-2}\) point source flux of muon neutrinos averaged over declination in the Northern hemisphere sky at 90% confidence level: \(E^2d\phi/dE < 5.5 \times 10^{-8}\) GeV cm\(^{-2}\) s\(^{-1}\) in the energy range of 1.6 TeV to 2.5 TeV.

Over the same period of time, a search for neutrino emission from 32 specific candidate sources chosen based on observations at various wavelengths in the electromagnetic spectrum has been performed\(^ {18}\). No statistically significant evidence for neutrino emission was found from any of the candidate sources. The highest observed significance, with 8 observed events compared to 4.7 expected background events (1.2\(\sigma\), is at the location of the GeV blazar 3C273. The second highest excess (1.1\(\sigma\)) is from the direction of the Crab Nebula, with 10 observed events compared to 6.7 expected background events.

In addition to searches for individual sources of neutrinos, AMANDA data taken between 2000 and 2003 have been used to set a limit on possible diffuse fluxes of neutrinos. Populations of distant sources could lead to such a diffuse flux that would clearly prove the acceleration of hadrons in astrophysical sources even if the sources cannot be resolved. This diffuse flux can be distinguished from the background of atmospheric neutrinos due to its harder spectra, expected from most astrophysical sources. This study relies on the the number of triggered OMs which serve as an energy estimator for AMANDA. A limit of \(E^2d\phi/dE < 7.4 \times 10^{-8}\) GeV cm\(^{-2}\) s\(^{-1}\) sr\(^{-1}\) is placed on the diffuse muon neutrino flux in the energy range from 16 TeV to 2.5 PeV at 90% confidence level\(^ {19}\). Additionally, AMANDA has searched for an all-flavour diffuse flux from the Southern sky, a work on these three years of data places a limit of \(E^2d\phi/dE < 2.7 \times 10^{-7}\) GeV cm\(^{-2}\) s\(^{-1}\) sr\(^{-1}\), in the energy range of \(2 \times 10^5\) to \(10^9\) GeV\(^ {23}\).
4.2 Searches for neutralino dark matter

AMANDA can be used to search for neutralino dark matter by looking for a neutrino flux excess from the center of the Earth \(^{20}\) or from the Sun \(^{21}\). The respective limits obtained with the 2001-2003 dataset for the Earth and the 2001 dataset for the Sun are given in Fig. 3. The figures show the muon flux limit from neutralino annihilations, along with the results from other indirect searches and predictions from theoretical models. Disfavoured models by recent direct searches with the XENON 10 experiment \(^{22}\) are shown as green dots.

5 First results from the IceCube 9 strings configuration

The IC-9 dataset has a total livetime of 137.4 days taken between June and November 2006. 234 neutrino candidates were identified on this data sample with \(211 \pm 76 \) (syst.) \(\pm 14\) (stat.) events expected from atmospheric neutrinos and less than 10\% pollution by the background of down-going muons \(^{24}\).

The zenith and azimuth angle distributions of these neutrino candidates are shown on Fig. 4. The agreement with simulation is good except for a discrepancy near the horizon due to a residual contamination of down going muons. This discrepancy would disappear with tighter event selection. One can notice the 2 strong peaks in the azimuth angle distribution (entry on the right in Fig. 4 on the right), due to the very asymmetric configuration of the detector and corresponding to the long axis of IC-9.

These data have been used to search for a possible accumulation of events in the sky \(^{25}\). The resulting sky-average point-source sensitivity for a source with an \(E^{-2}\) spectrum is \(E^2d\phi/dE = 12 \times 10^{-8}\ \text{GeVcm}^{-2}\text{s}^{-1}\) which is already comparable with what was obtained with the 5 years of AMANDA II data presented in section 4. The events were treated with a likelihood based analysis that makes use of the angular distribution of the background with a source hypothesis compared to a background only hypothesis obtained by scrambling the data in right ascension. The first significance map obtained with IceCube for the Northern hemisphere sky is shown on fig. 5. This map doesn’t show any significant deviation from uniformity. The most significant excess, with a 3.3 \(\sigma\) significance is at r.a. = 276.6\(^{\circ}\), dec. = 20.4\(^{\circ}\). This is comparable with a
random fluctuation of a uniform background as 60% of the datasets scrambled in right ascension show an excess of 3.3 σ or higher somewhere in the sky. A search for neutrinos coming from 26 galactic and extragalactic preselected objects has also been performed on this dataset. In addition, the most significant excess over the expected background on these sources was found at the Crab nebula with 1.77σ, which again is consistent with random fluctuations. Like in AMANDA, IceCube data can be used to probe the diffuse flux of neutrino from an unresolved population of astrophysical sources\textsuperscript{26}. The sensitivity of IC-9 is $1.4 \times 10^{-7}$ GeVcm$^{-2}$s$^{-1}$sr$^{-1}$ which is only a factor of 2 above the AMANDA-II sensitivity despite the much shorter integrated exposure time. Large improvements can be expected from both longer operation of IceCube with even more strings and refinement of analysis techniques.

6 Conclusion: The future of IceCube

6.1 The next years of IceCube

The accumulated exposure of the IceCube 9 strings configuration does not allow us yet to reach the integrated exposure level required to probe astrophysical neutrino signals. Nevertheless,
various analyses are developing and are very promising\cite{ref24}. These results confirm the stability of data taking, the good quality of the data recorded and experiment simulation. During the coming years, IceCube will continue to grow and will in 2009 reach an integrated exposure of 1 km$^3$\cdot yr. This will be an important milestone as it represents roughly what is needed to reach the level of detection for an astrophysical neutrino flux\cite{ref5}. When completed, the acceptance of the detector will naturally be larger, but it will also have an improved performance for reconstruction due to its larger size. In the case of the search for point sources for instance, the longer lever arm for the reconstruction of the muons tracks will lead to a better angular resolution of the detector which will become better than a degree.

6.2 One step further: extension of the IceCube detector at low energies with DeepCore

The capabilities of IceCube will be extended at both lower and higher energies in the near future. Starting next austral summer, a compact core of 6 strings using IceCube’s DOM technology, called the DeepCore detector, will start to be deployed near the center of the main in-ice detector. The interstring spacing will be of the order of 72 m, allowing for the exploration of energies as low as 10-20 GeV. The surrounding IceCube strings will be used as an active veto to reduce the atmospheric muon background. The energy range that is explored is very important for dark matter searches that were initiated with AMANDA. Moreover, the ability to select contained events opens the search for downgoing astrophysical neutrino signals at low energies. This will allow one to look above the current horizon of IceCube, even opening the possibility to look at the galactic center or sources like RX J1713.7-3946\cite{ref27}.

6.3 The second step: extensions at higher energies

At EeV energies, on the other end of the energy range, an extension of IceCube is also studied. The radio or the acoustic signal generated by neutrino interacting in the ice can be detected with a high energy extension of IceCube. With a much increased detection volume, we will aim at detecting the GZK neutrino flux. With attenuation lengths of the order of the kilometer for acoustic (kHz frequency range) and for radio signals (MHz frequency range), a sparse instrumentation will suffice for this extension. Two projects are currently explored for this extension: AURA (Askarian Underice Radio Array) for the radio signal\cite{ref28} and SPATS (South Pole Acoustic Test Setup) for the acoustic signal\cite{ref29}. They are currently studying the polar ice and developing the hardware necessary for the building of a hybrid detector enclosing IceCube in another array of strings with a much larger spacing that will allow to study these very scarce and energetic events.

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The Antares Neutrino Telescope: first results

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The Antares Collaboration is completing the deployment of a 12 lines underwater detector, 2500m deep in the Mediterranean Sea, dedicated to high energy neutrino astronomy. Starting with the first line in 2006, 10 lines were continuously recording data by the end of 2007, which allow us to reconstruct downward-going cosmic muons, and search for the first upward-going \( \nu \)-induced muons. Calibration topics will be described and preliminary results presented.

1 The neutrino as a new high-energy messenger

The advantage of using neutrinos as new messengers lies firstly on their weak interaction cross-section; unlike protons (\( E_{\text{cut-off}} \sim 5 \times 10^{19} \text{eV}, l_{\text{free path}} \sim 50 \text{ Mpc} \)) or \( \gamma \) (\( E_{\text{cut-off}} \sim 10^{14} \text{eV}, l_{\text{free path}} \sim 10 \text{ Mpc} \)), they provide a cosmological-range unaltered information from the very heart of their sources. Secondly, charged particles are deflected by magnetic fields, with a mean deflection \( \Delta \theta \sim L(\text{kpc}) \frac{ZB(\mu G)}{E(\text{EeV})} \), yielding for Galactic Sources \( \Delta \theta \sim 12^\circ \) at \( 10^{19} \text{eV} \). Neutrinos on the other hand point directly to their sources and exact production site.

The neutrinos Antares is aiming at are typically TeV neutrinos from AGNs (supermassive black holes believed to be hosted in the center of each galaxy), typically 30 orders of magnitude lower in flux\(^1\) than solar neutrinos. The detection of those specific neutrinos requires under water/ice instruments, or alternatively acoustic/radio techniques in the PeV-EeV range and air showers arrays above 1 EeV. In spite of efforts in those various energy ranges, since the detection of the MeV neutrino burst from SN 1987A by Kamiokande/Baksan/IMB/Mont-Blanc\(^2\) no astrophysical source for neutrinos above a few GeV has ever been identified.

\(^{1}\)http://antares.in2p3.fr
2 TeV cosmic neutrinos: production and detection

Sources for TeV \( \nu \) are typically compact objects (neutron stars/black holes), from which often emerge relativistic plasma jets with a still unclear composition - leptonic or hadronic?

2.1 Sources of TeV cosmic neutrinos?

Most of these sources have already been extensively studied from radio wavelengths up to \( \gamma \)-rays. These photons can be produced by \( e^- \) via inverse compton effect (on ambient photon field)/synchrotron radiation, or by protons/nuclei via photoproduction of \( \pi^0/\pi^\pm \):

\[
p/A + p/\gamma \rightarrow \pi^0 \pi^\pm, \quad \text{with} \quad \pi^0 \rightarrow \gamma \gamma, \quad \pi^\pm \rightarrow \nu_\mu \mu, \quad \mu \rightarrow \nu_\mu \nu_\tau e
\]

(1)

In the former scenario, no neutrinos are produced, whereas in the latter, the neutrino flux is directly related to the gamma flux: a TeV neutrino detection from gamma sources would then yield a unique way to probe the inner processes of the most powerful events in the universe. Several hints exist which indicates that hadrons could be accelerated up to very high energies. Firstly, the combined radio, X-rays and \( \gamma \)-rays observations of the shell-type supernova remnant RX J1713.7-3946\(^3\) favour the production of photons via \( \pi^0 \) decay (figure 1, left). Secondly, the correlations between X and \( \gamma \) for the Blazar 1ES1959+650\(^4\) prove the existence of \( \gamma \) flares not visible in X (figure 1, right), which is difficult to account for in purely leptonic models. Finally, it should be reminded that the so-called GZK cut-off (interaction of ultra-high energy cosmic rays with the CMB) is a guaranteed source of sub-EeV neutrinos\(^5\).

![Figure 1: Left: Multiwavelength observations of the SNR RXJ 1713.7-39; the solid curve at energies above \( 10^7 \) eV corresponds to \( \pi^0 \)-decay \( \gamma \)-ray emission, whereas the dashed and dash-dotted curves indicate the inverse Compton (IC) and Nonthermal Bremsstrahlung (NB) emissions, respectively. Right: Whipple vs RXTE flux, for the Blazar 1ES1959+650, which shows the existence of orphan \( \gamma \) flares (in red).](image)

2.2 Practical issues for their detection

ANTARES can be seen as a fixed target experiment: a cosmic muon neutrino interacts in the Earth and produces a muon that propagates in sea water. The Čerenkov light emitted by the muon is detected by an array of photomultipliers arranged in strings, able to reconstruct the energy and direction of the incident muon/neutrino\(^6\).

The main physical backgrounds are twofold. Atmospheric muons (\( \sim 1/s \) at the reconstruction level in ANTARES), produced in the upper atmosphere by the interaction of cosmic rays, can be strongly suppressed because of their downward direction. Upward-going atmospheric neutrinos (\( \sim 10/\)days in ANTARES) on the other hand are more delicate to identify: they have exactly the same signature as the expected cosmic signal ANTARES awaits for.
For a given neutrino flux $\Phi_\nu$, the number of events expected for a telescope of effective area $A_\mu$ (i.e. the size of a detector 100% efficient for muons) can be estimated as follows:

$$N_\mu \propto \Phi_\nu \times P_{\text{absorption}}(\theta, E_\nu) \times \sigma_\nu \times R_\mu \times A_\mu,$$

(2)

where typically $\sigma_\nu \approx 2 \times 10^{-34}$ cm$^2$, $R_\mu \approx 10$ km, and $A_\mu \approx 0.06$ km$^2$ (roughly the geometrical surface for ANTARES for reconstructed events with an angular resolution below 1°) at 100 TeV.

The neutrino luminosity $L_\nu = 4\pi d^2 \Phi_\nu$ needed to detect $N_\nu$ events can then be written:

$$L_\nu \approx 10^{46} N_\nu \left( \frac{d}{4 \text{ Gpc}} \right)^2 \left( \frac{E_\nu}{100 \text{ TeV}} \right)^{1-\alpha} \left( \frac{A_\mu T_{\text{obs}}}{\text{km}^2 \text{yr}} \right)^{-1} \text{erg/s},$$

(3)

for a source observed over a time $T_{\text{obs}}$ : $\alpha \sim 1/0.5$ below/above 100 TeV. Typically for blazars ($d \sim \text{Gpc}$ and $L \sim 10^{47}$ erg/s), the required effective area is $A_\mu \sim 1$ km$^2$, far beyond the reach of ANTARES. For galactic sources and $L \sim 10^{35}$ erg/s, the necessary effective area goes down to $A_\mu \sim 0.1$ km$^2$, typically the size of ANTARES.

3 ANTARES : description, performances & milestones

Two main kinds of signals can be detected with ANTARES: $\mu$ tracks initiated by the charged current interaction of a $\nu_\mu$ in the Earth, and showers produced by the interaction of a neutrino (mainly $\nu_e$ and $\nu_\tau$ by charged or neutral current channels) in water. Those signals are faint signals, and because of light scattering and absorption in water, their detection require single-photon-electron-sensitive devices. The measurements of the time of the hits (time resolution of the order of ns) and the amplitude of the hits (with a resolution of about 30%), together with the position of the hits (by measuring the position of each PMT, to reach a resolution of about 10 cm) are needed to achieve the reconstruction of those signals with the desired resolution.

Muon tracks are detected via their directional Čerenkov light (angle in water $\approx 42^\circ$) and can be reconstructed with an angular resolution below 0.3° above 10 TeV (the resolution below this energy is dominated by the kinematics of the interaction). The energy resolution is quite poor, a factor 2-3 on average, restricted by the granularity/density of the light sensors and the fact that the muon traverses the detector. Showers produced by $\nu_e$ on the other hand emit quasi-isotropic light, and can be reconstructed with a better energy resolution (roughly 30 %) but with a poorer angular resolution $\sim 3-5 \degree$.

3.1 Detector description

The ANTARES neutrino telescope, deployed at 2500 m below sea surface, 40 km off the coast of Toulon (Southern France) is composed of 12 strings, with 25 storeys each containing a triplet of
10° photomultipliers oriented at 45 degrees downward to be optimally sensitive to upward going muons. As of March 2008, 10 lines were connected and continuously taking data since end of 2007: the first line was operating as soon as March 2006, the second line in September 2006, and in January 2007 5 lines in whole were operational. A schematic description of the detector, together with the layout of the lines, can be found in figure 3 (left plot). The full completion of the telescope should be performed by summer 2008.

An instrumented line is also present on site, to perform environmental measurements: seawater temperature, salinity, sound velocity probes, as well as speed of the sea current and direction, all parameters required for an optimum track reconstruction, and for studies of biological backgrounds. The quality of sea water in particular, and its knowledge, is a fundamental parameter for Čerenkov photon detection. The absorption length at the Antares site at 470 nm is roughly 60 m, with an effective scattering length of 300 m. It is a combination of this water quality (scattering and chromatic dispersion, accounting for 1.5 ns at a distance of 40 m) and of the timing performances described below which finally takes down the angular resolution at the 0.2° level at high energy (where the neutrino and the produced muon are essentially colinear).

![Image of the Antares telescope and its components.](image)

**Figure 3:** Left: Description, position and layout of the 12 lines of the Antares telescope, 10 of which are currently taking data. Right: Instruments on board of one of the Antares storeys.

The right panel of figure 3 displays the content of one of the Antares storeys. PMTs are enclosed in pressure-resistant spheres, and a Titanium cylinder contains the front-end electronics. The intrinsic photoelectron transit time spread between the photocathode and its first dynode is roughly 1.3 ns, and the last dynode signal being digitised by a devoted chip, the ARS, gives a resolution better than 0.5 ns. The tilt/compass cards, and hydrophones on some of the storeys, allow us to measure continuously the position of the optical modules (see section 4). Finally, time calibration (see section 4) can be performed using a laser and LED beacons.

### 3.2 Physics performances

The energy resolution is a crucial element for the study of diffuse ν flux. The link between extragalactic sources of both cosmic rays, γ-rays and ν leads to severe limits on the ν diffuse flux expressed in the Waxman-Bahcall (WB) upper bound: $E^2 \Phi < 4.5 \times 10^{-8} \text{GeV.cm}^{-2} \text{s}^{-1} \text{sr}^{-1}$. After 3 years, Antares is expected to set an upper limit of $E^2 \Phi < 3.9 \times 10^{-8} \text{GeV.cm}^{-2} \text{s}^{-1} \text{sr}^{-1}$, just below the WB estimate.

The Antares sensitivity to point-like sources can be estimated as a function of the declination of a potential source: figure 4 (left) shows that Antares will be able to observe the Galactic Centre and other interesting γ sources, for most of the time complementary to IceCube. The 90% upper limit for $\nu_\mu + \bar{\nu}_\mu$ flux in case of null signal after 1 year is $E^2 \frac{dN}{dE_\nu} = 4 \times 10^{-8} \text{GeV.cm}^{-2} \text{s}^{-1}$ at $\delta = -90^\circ$, and rises to $1.5 \times 10^{-7} \text{GeV.cm}^{-2} \text{s}^{-1}$ at $\delta = +40^\circ$. 


Those limits improve those of MACRO for the Southern Sky, as can be seen in figure 4 (right), and are comparable to those obtained by AMANDA II for the Northern Sky.\(^\text{10}\)

Figure 4: Left: Sky as visible by ANTARES and AMANDA/IceCube, in Galactic coordinates, the circle indicating the Galactic Centre. Right: Sensitivity for Point-like sources.

### 3.3 ANTARES Milestones

Conceiving and building a Neutrino Telescope in the Mediterranean require much preparatory work: the proposal of the experiment dates back to 1999.\(^\text{11}\) Nine years were thus needed to realise the 10 lines (soon 12!) that are currently taking data. Here is a short historical overview, before describing the calibration and results of ANTARES:

- **1996-2000**: Validation of the Project - Water properties were first studied in order to choose the best site, and marine technologies were developed and improved. This period ended with the deployment of a demonstrator line and the reconstruction of the first atmospheric muons in the ANTARES Collaboration.

- **2001-2004**: Final R&D, first deployments - The Electro-Optical cable between the shore and the site was deployed in 2001, the Junction Box (distribution of power to lines) was operational in 2002. Finally, a Prototype Sector Line (similar to a final ANTARES line, but with only one sector consisting of 5 storeys) successfully took data between end of 2002 until its recovery in July 2003; a Mini-Instrumentation Line (MIL, environmental probes mainly) was also operated for a few months between Feb. and May 2003.\(^\text{13}\)

- **2005-2007**: Construction, deployment and operation - The MIL was recovered to be upgraded with two storeys of Optical Modules (MILOM), and took data for 2 years (April 2005 - March 2007), before the deployment, connection and operation of the first complete ANTARES line in March 2006.\(^\text{14}\)

### 4 ANTARES in operation: calibration of a neutrino telescope

To be able to extract physical results from raw data, a neutrino telescope like any other detector has to be understood and calibrated: some aspects of this calibration will be reviewed here.

#### 4.1 Acoustic positioning

The particularity of an underwater neutrino telescope, as compared with a $\nu$ Telescope in ice (ICECUBE), is that the lines, maintained as vertical as possible with a buoy, are moving under the
influence of water currents. Hence, a reconstruction of the line shape is needed for the positions of each individual PMTs to be known with an accuracy of 10-cm, which is required to achieve the $0.2^\circ$ angular resolution at high energy. This is performed by an acoustic positioning system\textsuperscript{15}, as shown in figure 5 (left panel). Transponders on the sea ground emit signals, detected by hydrophones equipping some of the storeys. Together with data from tilt/compass cards, this allows for the determination of shape of the line; the actual position of the top storey can differ from the original straight line position by up to 15 m for strong sea currents! If not accounted for, this would imply an error in the absolute positioning of a source of several degrees.

![Figure 5: Left: Principles of acoustic positioning. Right: Principles of Time Calibration with LED Beacons.](image)

### 4.2 Time calibration

An error of 0.3 ns in the photon arrival time on one of the PMTs is equivalent to a 10-cm error on its position: timing performances are thus as primordial as the line positioning. They can be studied using the light emitted by LED Beacons\textsuperscript{16}. This blue light is detected by PMTs on adjacent lines, and the timing resolution as shown in figure 5 (right panel) can be estimated to be *(horizontal case)* as low as 0.7 ns, which is then dominated by the electronics.

### 4.3 $^{40}$K calibration

Sea water contains $^{40}$K which is a $\beta$ emitter, the $e^-$ in turn emitting Čerenkov radiation. Adjacent PMTs can thus coincidently detect this light, and this $^{40}$K calibration is a powerful way to estimate the acceptance of each optical module\textsuperscript{17}.

### 5 Antares in operation: first signals and selected results

The trigger rate of Antares is roughly 1/s, mostly corresponding to atmospheric muons, 70% of which are multiple quasi-parallel muons, arriving at the same time in the detector\textsuperscript{18}. A nice muon bundle seen with 10 lines is displayed in figure 6. Showers developing along a $\mu$ track can also be observed (fig. 6, right).

#### 5.1 Line 1 data: first estimate for atmospheric muons flux

The angular distribution of reconstructed events can be transformed into an intensity *versus* depth (using acceptance corrections from simulation) in the region of uniform acceptance: each value of the zenith angle corresponds to a certain slant depth through the water mass above the detector. To compute the muon vertical intensity, the distribution of muons at sea level has to
be taken into account\textsuperscript{19}. The results obtained using Line 1 data, with low sea currents from May to September 2006 (equivalent live time 10 days), are shown in figure 7. The errors of the order of 50% are dominated by the PMT acceptance. The agreement between data and other published values is good, showing that physics results can be extracted even with only 1 line!

### 5.2 Data with 5 lines: neutrino candidates

Figure 7 (right) shows a zenith angle distribution obtained with 5 Lines data (February-May 2007, equivalent live time 54 days). These data contain roughly $5 \times 10^6$ events, reconstructed with a 90% efficiency. After quality cuts, 20000 events remain, for which $\cos \theta$ is shown. The events reconstructed as upgoing ($\cos \theta > 0.1$) are 55 neutrino candidates, the events reconstructed as downgoing corresponding to atmospheric muons. The peak at -1 are vertically downward-going atmospheric $\mu$, or muon bundles, very nicely reconstructed. Finally the slight excess near $\cos \theta \sim 1$ is an acceptance effect: the telescope is more sensitive to purely vertical tracks.

![Figure 6: A $\mu$ bundle seen in 10 lines (altitude of hits vs time of hits, left) and a shower in 1 line (right), from Line 1 data (top) and Monte-Carlo (bottom).](image)

![Figure 7: Left: Line 1 data, vertical intensity of atmospheric muons versus depth (water equivalent). Right: 5 lines data, distribution in zenith angle of (selected) events.](image)
6 Conclusions and outlook

The 2 remaining ANTARES lines should be taking data by summer 2008, giving birth to the biggest Neutrino Telescope in the Northern Hemisphere. The current 10-lines telescope is already operating, and its acoustic positioning is fully functional. Despite its smaller size with respect to ICECUBE\textsuperscript{20}, ANTARES observes the Galactic Centre and other potential sources of TeV $\nu$ not accessible from the South Pole, leaving some margins for unexpected discoveries.

Furthermore, ANTARES is a part of the gcn, Gamma-ray bursts Coordinates Network\textsuperscript{21}, dedicated to $\gamma$-ray bursts, thought to be potential sources of high energy neutrinos: satellites/telescopes broadcast real-time alerts, which in turn trigger the recording of all ANTARES data within a 2 minutes time window\textsuperscript{22}. Over a period of 15 months, 172 gcn alerts were distributed to ANTARES, and the telescope took data for 152 of them, corresponding to a $\sim$ 90\% live time!

ANTARES must be seen as the first stage towards a km$^3$-scale telescope, for which European institutes involved in current $\nu$ astronomy projects (ANTARES, NEMO, NESTOR) are already collaborating. This network, KM3NeT\textsuperscript{23}, will give birth to a telescope with which neutrinos will be as common messengers as gamma-rays are now.

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TESTING DARK MATTER WITH NEUTRINO DETECTORS

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Neutrinos are the least detectable Standard Model particle. By making use of this fact, we consider dark matter annihilations and decays in the galactic halo and show how present and future neutrino detectors could be used to set general limits on the dark matter annihilation cross section and on the dark matter lifetime.

1 Introduction

With the next generation of neutrino experiments we will enter the era of precision measurements in neutrino physics. As a consequence, a lot of efforts are being dedicated to decide which are the best experimental set-ups. However and in addition to the detailed study of neutrino parameters, present and future neutrino detectors, thanks to their great capabilities, might also be used for other purposes. Among the possible synergies of these detectors, they could be used to test some of the properties of the dark matter (DM) of the Universe. For instance, it has been pointed out that by using the spectral information of neutrinos coming from annihilations of DM particles in the center of the Sun, some of the DM properties could be reconstructed. In this talk however, we consider neutrinos coming from DM annihilations or decays in our galactic halo and show how they can be used to test some other DM properties.

We will use the fact that among the Standard Model (SM) particles, neutrinos are the least detectable ones. Therefore, if we assume that the only SM products from the DM annihilations (decays) are neutrinos, a limit on their flux, conservatively and in a model-independent way, sets an upper (lower) bound on the DM annihilation cross section (lifetime). This is the most conservative assumption from the detection point of view, that is, the worst possible case. Any other channel (into at least one SM particle) would produce photons and hence would give rise to a much more stringent limit. Let us stress that this is not an assumption about a particular and realistic case. On the other hand, for the reasons just stated, it is valid for any generic model, in which DM annihilates (decays) at least into one SM particle. Hence, the bounds so
obtained are bounds on the total annihilation cross section (lifetime) of the DM particle and not only on its partial annihilation cross section (lifetime) due to the annihilation (decay) channel into neutrinos.

In this talk, and following and reviewing the approach of Refs.\textsuperscript{2,3,4,5}, we consider this case and evaluate the potential neutrino flux from DM annihilation (decay) in the whole Milky Way, which we compare with the relevant backgrounds for detection. In such a way, we obtain general constrains on the DM annihilation cross section and on the DM lifetime, which are more stringent than previous ones\textsuperscript{6,7,8,9,10}.

2 Neutrino Fluxes from the Milky Way

Detailed structure formation simulations show that cold DM clusters hierarchically in halos which allows the formation of large scale structure in the Universe to be successfully reproduced. In the case of spherically symmetric matter density with isotropic velocity dispersion, the simulated DM profile in the galaxies can be parametrized via

$$\rho(r) = \rho_{\text{sc}} \left( \frac{R_{\text{sc}}}{r} \right)^{\gamma} \left[ 1 + \left( \frac{R_{\text{sc}}}{r_s} \right)^{\alpha} \right]^{(\beta - \gamma)/\alpha},$$

where $R_{\text{sc}} = 8.5$ kpc is the solar radius circle, $\rho_{\text{sc}}$ is the DM density at $R_{\text{sc}}$, $r_s$ is the scale radius, $\gamma$ is the inner cusp index, $\beta$ is the slope as $r \to \infty$ and $\alpha$ determines the exact shape of the profile in regions around $r_s$. Commonly used profiles\textsuperscript{11,12,13} (see also Ref.\textsuperscript{14}) tend to agree at large scales, although they differ considerably in the inner part of the galaxy.

The differential neutrino plus antineutrino flux per flavor from DM annihilation or decay in a cone of half-angle $\psi$ around the galactic center, covering a field of view $\Delta \Omega = 2 \pi (1 - \cos \psi)$, is given by

$$\frac{d\Phi}{dE_{\nu}} = \frac{\Delta \Omega}{4 \pi} \mathcal{P}_k(E_{\nu}, m_\chi) R_{\text{sc}} \rho_0^k \mathcal{J}_{\Delta \Omega, k},$$

where $m_\chi$ is the DM mass, $\rho_0 = 0.3$ GeV cm\textsuperscript{-3} is a normalizing DM density, which is equal to the commonly quoted DM density at $R_{\text{sc}}$, and $\mathcal{J}_{\Delta \Omega, k}$ is the average in the field of view (around the galactic center) of the line of sight integration of the DM density (for decays, $k = 1$) or of its square (for annihilations, $k = 2$), which is given by

$$\mathcal{J}_{\Delta \Omega, k} = \frac{2 \pi}{\Delta \Omega} \frac{1}{R_{\text{sc}} \rho_0^k} \int_{\cos \psi}^{1} \int_{0}^{l_{\text{max}}} \rho(r)^k \ d l (\cos \psi'),$$

where $r = \sqrt{R_{\text{sc}}^2 - 2 R_{\text{sc}} \cos \psi' + l'^2}$ and $l_{\text{max}} = \sqrt{(R_{\text{halo}}^2 - \sin^2 \psi R_{\text{sc}}^2) + R_{\text{sc}} \cos \psi}$. The contribution at large scales is negligible and thus, this integral barely depends on the size of the halo for $R_{\text{halo}} \gtrsim$ few tens of kpc.

The factor $\mathcal{P}_k$ embeds all the dependences on the particle physics model and it reads

$$\mathcal{P}_1 = \frac{1}{3} \frac{dN_1}{dE_{\nu}} \frac{1}{m_\chi \tau_\chi} \quad \text{for decays and} \quad \mathcal{P}_2 = \frac{1}{3} \frac{dN_2}{dE_{\nu}} \frac{\langle \sigma_A v \rangle}{2 m_\chi^2} \quad \text{for annihilations},$$

where the neutrino plus antineutrino spectrum per flavor is given by

$$\frac{dN_1}{dE_{\nu}} = 2 \delta(E_{\nu} - \frac{m_\chi}{2}) \quad \text{for decays and} \quad \frac{dN_2}{dE_{\nu}} = 2 \delta(E_{\nu} - m_\chi) \quad \text{for annihilations},$$

and the factor of $1/3$ comes from the assumption that the annihilation or decay branching ratio is the same for the three neutrino flavors. Let us note that this is not a very restrictive assumption, for even even when only one flavor is predominantly produced, there is a guaranteed
flux of neutrinos in all flavors thanks to the averaged neutrino oscillations between the source and the detector. Hence, although different initial flavor ratios would give rise to different flavor ratios at detection, the small differences affect little our results and for simplicity herein we consider flavor democracy.

2.1 Annihilations versus Decays: DM Halo Uncertainties

As mentioned above, while DM profiles tend to agree at large scales, uncertainties are still present for the inner region of the galaxy. In the two cases considered (annihilations and decay), the overall normalization of the flux is affected by the value of $J_{\Delta \Omega, k}$. However, in the case of DM annihilations, it scales as $\rho^2$, whereas for DM decays, it scales as $\rho$. Our lack of knowledge of the halo profile is hence much more important for the neutrino flux from DM annihilations. For the three profiles considered here, astrophysical uncertainties can induce errors of up to a factor of 6 for the case of DM decays, but they can be as large as a factor of ~100 for DM annihilations. In addition, if the DM mass is not known, DM annihilation and DM decay in the halo might have the same signatures. However, due to the fact that the dependence on the DM halo density is different for each case, in case of a positive signal, directional information would be crucial to distinguish between these two possibilities.

For concreteness, in what follows we present results using the Navarro, Frenk and White (NFW) simulation as our canonical profile.

3 Neutrino Bounds

In order to obtain the constraints on the DM annihilation cross section and DM lifetime we assume that DM annihilates or decays only into neutrinos. If DM annihilates or decays into SM particles, neutrinos (and antineutrinos) are the least detectable ones. Any other possible annihilation or decay mode would produce gamma rays, which are much easier to detect, and would allow to set a much stronger (and model-dependent) bound. Thus, the most conservative approach is to assume that only neutrinos are produced in DM annihilations or decays. Even in this conservative case, it has been shown that stringent limits can be obtained by comparing the expected time-integrated annihilation signal of all galactic halos and the signal from annihilations or decays in the Milky Way Halo with the background at these energies.

3.1 The Atmospheric Neutrino Background

For $E_\nu \gtrsim 100$ MeV, the main source of background for a possible neutrino signal from DM annihilations or decays is the flux of atmospheric neutrinos, which is well known up to energies of ~100 TeV. Thus, in order to obtain a bound on the DM annihilation cross section and lifetime we need to compare these two fluxes, and in particular we consider the $\nu_\mu + \bar{\nu}_\mu$ spectra calculated with FLUKA.

In this energy range, we will follow the approach of Ref. By assuming that the only resultant products of DM annihilation (decay) are neutrino-antineutrino pairs, we first obtain a general bound by comparing the $\langle \nu_\mu + \bar{\nu}_\mu \rangle$ neutrino flux from DM annihilation (decays) in the halo with the corresponding atmospheric neutrino flux for $E_\nu \sim 100$ MeV–100 TeV in an energy bin of width $\Delta \log_{10} E_\nu = 0.3$ around $E_\nu = m_\chi$ ($E_\nu = m_\chi/2$). For each value of $m_\chi$, the limit on $\langle \sigma A \rangle (\tau_\chi)$ is obtained by setting its value so that the neutrino flux from DM annihilations (decays) in the Milky Way equals the atmospheric neutrino spectrum integrated in the chosen energy bin. The reason for choosing this energy bin is mainly that the neutrino signal is sharply peaked around a neutrino energy equal to the DM mass (half of the DM mass) and this choice is within the experimental limits of neutrino detectors.
The most conservative bound is obtained by using the full-sky signal, and this is shown in both panels of Fig. 1 where the dark areas represent the excluded regions. However, a better limit can be obtained by using angular information. This is mainly limited by the kinematics of the interaction. In general, neutrino detectors are only able to detect the produced lepton and its relative direction with respect to the incoming neutrino depends on the neutrino energy as $\Delta \theta \sim 30^\circ \times \sqrt{\text{GeV}/E_\nu}$. As in Ref. 3 and being conservative, we consider a field of view with a half-angle cone of $30^\circ$ ($30^\circ \times \sqrt{10\text{GeV}/E_\nu}$) for neutrinos with energies above (below) 10 GeV. This limit is shown in both panels of Fig. 1 by the dashed lines (light areas), which improves upon the previous case by a factor of a few for $E_\nu > 5$ GeV.

### 3.2 MeV Dark Matter

As we have just described, it is expected that a more detailed analysis, making a more careful use of the directional as well as energy information for a given detector, will improve these results. Note for instance that for energies $\sim 1$-100 GeV neutrino oscillations would give rise to a zenith-dependent background, whereas we expect a nearly flat background for other energies for which oscillations do not take place. We now show how a more careful treatment of the energy resolution and backgrounds can substantially improve these limits.

Here we describe the analysis followed in Refs. 4,5 to set neutrino constraints on the DM total annihilation cross section and DM lifetime in the energy range $15\text{ MeV} \lesssim E_\nu \lesssim 130\text{ MeV}$. In this energy range the best data comes from the search for the diffuse supernova background by the Super-Kamiokande (SK) detector which has looked at positrons (via the inverse beta-decay reaction, $\nu_e + p \rightarrow e^+ + n$) in the energy interval 18 MeV–82 MeV 16. As for these energies there is no direction information, we consider the full-sky $\nu_e$ signal. In this search, the two main sources of background are the atmospheric $\nu_e$ and $\nu_e$ flux and the Michel electrons and positrons from the decays of sub-threshold muons. Below 18 MeV, muon-induced spallation products are the dominant background, and below $\sim 10$ MeV, the signal would be buried below the reactor antineutrino background.

Although for $E_\nu \lesssim 80\text{ MeV}$ the dominant interaction is the inverse beta-decay reaction (with free protons), the interactions of neutrinos (and antineutrinos) with the oxygen nuclei contribute significantly and must be considered. For our analysis we have included both the interactions of
\( \bar{\nu}_e \) with free protons and the interactions of \( \nu_e \) and \( \bar{\nu}_e \) with bound nucleons, by considering, in the latter case, a relativistic Fermi gas model \(^{12}\) with a Fermi surface momentum of 225 MeV and a binding energy of 27 MeV. We then compare the shape of the background spectrum to that of the signal by performing a \( \chi^2 \) analysis, analogous to that of the SK collaboration \(^{16}\). In this way, we can extract the limits on the DM annihilation cross section and DM lifetime \(^{4,5}\). Hence, we consider the sixteen 4-MeV bins in which the data were divided and define the following \( \chi^2 \) function \(^{16}\)

\[
\chi^2 = \sum_{l=1}^{16} \frac{[(\alpha \cdot A_l) + (\beta \cdot B_l) + (\gamma \cdot C_l) - N_l]^2}{\sigma^2_{\text{stat}} + \sigma^2_{\text{sys}}},
\]

where the sum \( l \) is over all energy bins, \( N_l \) is the number of events in the \( l \)th bin, and \( A_l \), \( B_l \) and \( C_l \) are the fractions of the DM annihilation or decay signal, Michel electron (positron) and atmospheric \( \nu_e \) and \( \bar{\nu}_e \) spectra that are in the \( l \)th bin, respectively. The fractions \( A_l \) are calculated taking into account the energy resolution of SK, interactions with free and bound protons and the correct differential cross sections \(^4\). The fractions \( B_l \) are calculated taking into consideration that in water 18.4\% of the \( \mu^- \) produced below Čerenkov threshold (\( p_\mu < 120 \text{ MeV} \)) get trapped and enter a K-shell orbit around the oxygen nucleus and thus, the electron spectrum from the decay is slightly distorted with respect to the well-known Michel spectrum \(^{18}\). In the calculation of the fractions \( B_l \) and \( C_l \) we have used the low energy atmospheric neutrino flux calculation with FLUKA \(^{19}\). Note that, in a two-neutrino approximation and for energies below \( \sim 300 \) MeV (where most of the background comes from), half of the \( \nu_e \) have oscillated to \( \nu_\tau \), whereas \( \nu_\tau \) remain unoscillated. Although this approximation is not appropriate, in principle, to calculate the low energy atmospheric neutrino background, however, for practical purposes, it introduces very small corrections \(^{20}\). Thus, in order to calculate \( B_l \) and \( C_l \) we use the two-neutrino approximation. The fitting parameters in the \( \chi^2 \)-function are \( \alpha \), \( \beta \) and \( \gamma \), which represent the total number of each type of event. For the systematic error we take \( \sigma_{\text{sys}} = 6\% \) for all energy bins \(^{16}\).

In absence of a DM signal, a 90\% confidence level (C.L.) limit can be set on \( \alpha \) for each value of the DM mass. The limiting \( \alpha_{90} \) is defined as

\[
\int_0^{\alpha_{90}} P(\alpha) \, d\alpha = 0.9,
\]

where \( P(\alpha) = K \cdot e^{-\chi^2_{\alpha}/2} \) is the relative probability and \( \chi^2_{\alpha} \) is the minimum \( \chi^2 \) for each \( \alpha \). The normalizing constant \( K \) is such that \( \int_0^{\infty} P(\alpha) \, d\alpha = 1 \). It is straightforward to translate the limit on \( \alpha \) into limits of the total DM annihilation cross section and DM lifetime and these 90\% CL bounds are shown in both panels of Fig. 1 by the hatched areas and they clearly improve (and extend to lower masses) by about an order of magnitude upon the general and very conservative bound obtained with the simple analysis described above for higher energies.

4 Conclusions

In this talk we have shown how neutrino detectors can also be used to test some of the DM properties and have obtained general bounds on the DM annihilation cross section and DM lifetime, which greatly improve over previous limits \(^{6,7,8,9,10}\). In order to do so, we have assumed that the only SM products from DM annihilations or decays are neutrinos, which are the least detectable particles of the SM. By making this assumption we have obtained conservative but model-independent bounds. In a simple way and for energies between \( \sim 100 \text{ MeV} \) and \( \sim 100 \text{ TeV} \), we have considered the potential signal from DM annihilations or decays in the Milky Way and

\(^{*}\)Note that there is an error in Eq.(8) of Ref.\(^{4}\). Nevertheless, this implies very small corrections to the results presented. I thank O. L. G. Peres for pointing this out.
have compared it to the atmospheric neutrino background. The general bounds are obtained by considering this potential signal and imposing that it has to be at most equal to the background in a given energy interval. We have also shown how this crude, but already very stringent limit, can be substantially improved by more detailed analysis which make careful use of the angular and energy resolution of the detectors, as well as of backgrounds. In this way, we have obtained the 90% CL bounds on the DM annihilation cross section and DM lifetime for $m_x \lesssim 200$ MeV, which is about an order of magnitude more stringent.

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References

Oscillations of neutrino emerging from a supernova core are studied. In this extremely high density region neutrino self interactions induce collective flavor transitions. When collective transitions are decoupled from matter oscillations, as for our chosen matter profile, an analytical interpretation of the collective effects is possible, by means of a mechanical analogy with a spherical pendulum. For inverted neutrino hierarchy the neutrino propagation can be divided in three regimes: synchronization, bipolar oscillations, and spectral split. Our simulation shows that averaging over neutrino trajectories does not alter the nature of these three regimes.

1 Introduction

Supernova neutrino oscillations are a very important tool to study astrophysical processes and to better understand neutrino properties\textsuperscript{1}. When neutrinos leave the surface of the neutrinosphere, they undergo vacuum and matter oscillations. Beside this, in the first few hundred kilometers neutrino-neutrino interactions induce collective flavor transitions, whose effect can be very important, depending on the neutrino mass hierarchy. Self-interaction effects are expected to be non negligible when $\mu(r) \sim \omega$, where $\mu(r)$ is the neutrino potential associated to the neutrino background ($\mu = \sqrt{2} G_F (N_\nu(r) + \overline{N}_\nu(r))$, analogously to the MSW potential $\lambda = \sqrt{2} G_F N_{\nu_{e-}}(r)$) and $\omega$ is the vacuum oscillation frequency. We neglect the solar mass square difference $\delta m^2 = m_2^2 - m_1^2 \ll \Delta m^2 = |m_3^2 - m_{1,2}^2|$, and consider a two-neutrino mixing scenario where the oscillations are governed by the mixing angle $\theta_{13}$. Since in the supernova context $\nu_\mu$ and $\nu_\tau$ cannot be distinguished we generically speak of $\nu_e \leftrightarrow \nu_x$ oscillations. In our work we assume $\Delta m^2 = 10^{-3} \text{ eV}^2$ and $\sin^2 \theta_{13} = 10^{-4}$. Figure 1 shows the radial profiles of the matter

\textsuperscript{a}Speaker
potential $\lambda(r)$ and of the neutrino potential $\mu(r)$, and the approximate ranges where collective flavor transitions of different type occur: synchronization, bipolar oscillation and spectral split. The nonlinearity of the self interactions induce neutrino oscillations very different from the ordinary MSW effect. When undergoing collective flavor transition neutrinos and antineutrinos of any energy behave similarly, as we will see in the following. This kind of transitions occurs for small $r$, well before the ordinary MSW resonance, allowing for a clear interpretation of the numerical simulations. For matter profiles different from our own, the MSW resonance condition can occur in the same region of the collective transitions: shallow electron density profiles can trigger MSW effects around $O(100)$ km. In that case it is much more difficult to disentangle collective from MSW effects in the results of the simulations.

2 Reference model and pendulum analogy

In our work, we use normalized thermal spectra with $\langle E_e \rangle = 10$ MeV, $\langle E_\nu \rangle = 15$ MeV, and $\langle E_x \rangle = \langle E_x \rangle = 24$ MeV for $\nu_e$, $\nu_x$, $\nu_x$ and $\nu_x$, respectively. The geometry of the model, the so called “bulb model” $^2$, has a spherical symmetry, since we assume that neutrinos are half-isotropically emitted from the neutrinosphere. Along any radial trajectory there is, therefore, a cylindrical symmetry. By virtue of that, we need only two independent variables to describe the neutrino propagation and interaction: the distance form the supernova center $r$, and the angle $\vartheta$ between two interacting neutrinos. If the dependence on $\vartheta$ is integrated out, we speak of “single-angle” approximation, while the general situation of variable $\vartheta$ is dubbed “multi-angle” case. The numerical simulation in the multi-angle case is extremely challenging, since it requires the solution of a large system (size of order $10^5$) of coupled non-linear equations. The propagation
of neutrinos of given energy $E$ is studied through the Liouville equation for the density matrix. By expanding the density matrix on the Pauli matrices and on the identity, the equations of motion can be expressed in terms of two polarization vectors, $\mathbf{P}(E)$ and $\overline{\mathbf{P}}(E)$, for neutrinos and antineutrinos, respectively. By introducing a vector $\mathbf{B}$ that depends on the mixing angle $\theta_{13}$, and a vector $\mathbf{D} = \mathbf{J} - \mathbf{J}$ that is the difference between the integral over the energy of $\mathbf{P}$ and $\overline{\mathbf{P}}$, the equations of motion can be written as

$$\dot{\mathbf{P}} = (+\omega \mathbf{B} + \lambda \mathbf{z} + \mu \mathbf{D}) \times \mathbf{P},$$

$$\dot{\overline{\mathbf{P}}} = (-\omega \mathbf{B} + \lambda \mathbf{z} + \mu \mathbf{D}) \times \overline{\mathbf{P}}.$$  \hspace{1cm} (1)

In the general case, the polarization vectors depend also on the neutrino emission angle $\theta_0$ (the neutrino incidence angle $\theta$ can be expressed in terms of $r$ and of the emission angle at the neutrinosphere $\theta_0$). The electron neutrinos survival probability $P_{ee}$ is a function of the polarization vector, $P_{ee} = 1/2(1 + P_i^f/P_i^i)$, where the $i$ and $f$ refer to the initial and final state respectively (analogously for antineutrinos). The equations of motion for $\mathbf{P}(E)$ and $\overline{\mathbf{P}}(E)$ can be reduced (under reasonable approximations) to the equations of motion of a gyrosopic pendulum, a spherical pendulum of unit length in a constant gravity field, characterized by a point-like massive bob spinning around the pendulum axis with constant angular momentum. The pendulum inertia is inversely proportional to $\mu(r)$, while its total angular momentum depends on the difference of the integrated polarization vectors $\mathbf{J}$ and $\overline{\mathbf{J}}$. The motion of a spherical pendulum is, in general, a combination of a precession and a nutation. In the case of normal hierarchy of the neutrino mass spectrum the pendulum starts close to the stable, downward position and stays close to it, as $\mu$ slowly decreases and no collective effect is present. In the inverted hierarchy case, the pendulum starts close to the “unstable,” upward position. At the beginning, for small $r$, when $\mu$ is large ($m$ is small), the bob spin dominates and the pendulum remains precessing in the upward position conserving angular momentum, a situation named synchronization. Nevertheless, since $\mu$ decreases with $r$, at a certain point any $\theta_{13} \neq 0$ triggers the fall of the pendulum and its subsequent nutations, the so called bipolar oscillations. The increase of the
pendulum inertia with \( r \) reduces the amplitude of the nutations, and bipolar oscillations are expected to vanish when self-interaction and vacuum effects are of the same size. At this point, at the end of the bipolar regime, self-interaction effects do not completely vanish and the spectral split builds up: a “stepwise swap” between the \( \nu_e \) and \( \nu_x \) energy spectra. The neutrino swapping can be explained by the conservation of the pendulum energy and of the lepton number \( \nu \).

The lepton number conservation is related to the constancy of \( D_z = J_z - \overline{J}_z \), that is a direct consequence of the equation of motion. For a detailed description of the pendulum analogy and of our reference model the reader is referred to our previous work \(^3\) and references therein.

### 3 Simulations

Figures 2 and 3 show the third component of \( \mathbf{P} \) and \( \overline{\mathbf{P}} \), as a function of the radius, for different energy values, for the single- and multi-angle simulations, respectively. Bipolar oscillations start at the same \( r \) and their periods are equal for both \( \nu \) and \( \overline{\nu} \) at any energy, confirming the appearance of a self-induced collective behavior, in the single- and in the multi-angle case. The behavior of each \( P_z \) and \( \overline{P}_z \) depends on its energy. For neutrinos, Figure 2, the spectral split starts around the critical energy \( E_c \approx 7 \text{ MeV} \): the curve relative to \( E < E_c \) ends up at the same initial value \( (P_{ee} = 1) \), while the curves for \( E > E_c \) show the \( P_z \) inversion \( (P_{ee} = 0) \). Neutrinos with an energy of \( \sim 19 \text{ MeV} \) do not oscillate much, because this is roughly the energy for which the initial \( \nu_e \) and \( \nu_x \) fluxes are equal. For antineutrinos, all curves show almost complete polarization reversal, with the exception of small energies (of few MeV, not shown in Figure 3). Figures 4 and 5 show the evolution of \( J \) and \( J_z \) for neutrinos and antineutrinos, in the single- and multi-angle cases. The behavior of these vectors can be related to the gyroscopic pendulum motion. At the beginning, in the synchronized regime, all the polarization vectors are aligned so that \( J = J_z \) and \( \overline{J} = \overline{J}_z \): the pendulum just spins in the upward position without falling. Around \( \sim 70 \text{ km} \) the pendulum falls for the first time and nutations appear. The nutation amplitude gradually decreases and bipolar oscillations eventually vanish for \( r \approx 100 \text{ km} \). At the same time, the spectral split builds up: antineutrinos tend to completely reverse their polarization, while this happens only partially for neutrinos. As said before, also for antineutrinos there is a partial swap of the spectra for \( E \sim 4 \text{ MeV} \). From Figure 5 it appears that bipolar oscillations of \( J \) and \( \overline{J} \) are largely smeared out in the multi-angle case. The bipolar regime starts somewhat later with respect to the single-angle case, since neutrino-neutrino interaction angles can be larger than the (single-angle) average one, leading to stronger self-interaction effects, that force the system.
in synchronized mode slightly longer. However, just as in the single-angle case, the spectral split builds up, $\mathcal{J}_z$ gets finally reversed, while the difference $D_z = J_z - \mathcal{J}_z$ remains constant. Figures 6 and 7 show the final neutrino and antineutrino fluxes, in the single- and multi-angle simulations. The neutrinos clearly show the spectral split effect and the corresponding sudden swap of $\nu_e$ and $\nu_x$ fluxes above $E_c \simeq 7$ MeV. In the right panel of Figure 6, the final antineutrino spectra are basically completely swapped with respect to the initial ones, except at very low energies, where there appears an “antineutrino” spectral split. This phenomenon can be related to the loss of $\mathcal{J}$ and of $|\mathcal{J}_z|^3$. Also in the multi-angle case of Figure 7, the neutrino spectral swap at $E > E_c \simeq 7$ MeV is rather evident, although less sharp with respect to the single-angle case, while the minor feature associated to the “antineutrino spectral split” is largely smeared out.

4 Conclusions

We have studied supernova neutrino oscillations in a model where the collective flavor transitions (synchronization, bipolar oscillations, and spectral split) are well separated from the MSW resonance. We have performed numerical simulations in both single- and multi-angle cases, using continuous energy spectra with significant $\nu-\overline{\nu}$ and $\nu_x-\nu_x$ asymmetry. The results of the single-angle simulation can be analytically understood to a large extent by means of a mechanical analogy with the spherical pendulum. The main observable effect is the swap of energy spectra, for inverted hierarchy, above a critical energy dictated by lepton number conservation. In the multi-angle simulation, the details of self-interaction effects change (e.g., the starting point of bipolar oscillations and their amplitude), but the spectral swap remains a robust, observable feature. In this sense, averaging over neutrino trajectories does not alter the main effect of the self interactions. The swapping of neutrino and antineutrino spectra could have an impact on $r$-process nucleosynthesis, on the energy transfer to the shock wave during the supernova explosion and on the propagation of the neutrinos through the shock wave. From the point of view of neutrino parameters, collective flavor oscillations in supernovae could be instrumental in identifying the inverse neutrino mass hierarchy, even for very small $\theta_{13}$.\textsuperscript{8}
Figure 7: Multi-angle simulation in inverted hierarchy: final fluxes (at $r = 200 \text{ km}$, in arbitrary units) for different neutrino species as a function of energy. Initial fluxes are shown as dotted lines to guide the eye.

References

V - Baryogenesis, Dark matter
and Astroparticle Physics

Cross-Reference: indirect dark matter searches are found under «Neutrino: astrophysics and astroparticle aspects», dark matter modes are also considered in the Standard Model and in the searches sections.
The quantum Boltzmann equations relevant for leptogenesis, obtained using non-equilibrium quantum field theory, are described. They manifest memory effects leading to a time-dependent CP asymmetry which depends upon the previous history of the system. This result is particularly relevant in resonant leptogenesis where the asymmetry is generated by the decays of nearly mass-degenerate right-handed neutrinos. The impact of the non-trivial time evolution of the CP asymmetry is discussed either in the generic resonant leptogenesis scenario or in the more specific Minimal Lepton Flavour Violation framework. Significant quantitative differences arise with respect to the usual approach in which the time dependence of the CP asymmetry is neglected.

1 Introduction

In our universe, the difference between the number densities of baryons and anti-baryons, per entropy density, is observed to be $Y_B \equiv (n_B - n_{\bar{B}})/s = (8.84 \pm 0.24) \times 10^{-11}$. This number, obtained from measurements of the Cosmic Microwave Background Radiation, is also in excellent agreement with the independent fit from Big Bang Nucleosynthesis (BBN). Thermal leptogenesis is a simple and well-motivated mechanism to explain this baryon asymmetry. The simplest implementation of this mechanism is realized by adding three right-handed (RH) Majorana neutrinos to the Standard Model (SM), i.e. the framework of type I see-saw. The fact that the same see-saw framework may simultaneously account for small neutrino masses and the baryon asymmetry of the universe makes it very attractive. In thermal leptogenesis, the heavy RH neutrinos are produced by thermal scatterings in the early universe after inflation, and subsequently decay out of equilibrium in a lepton number and CP violating way, thus satisfying Sakharov's conditions. A lepton asymmetry then arises, which is partially converted into a baryon asymmetry by electroweak sphaleron interactions.
In the case where the RH neutrinos masses are hierarchical, successful leptogenesis requires the RH neutrinos to be heavier than $10^9$ GeV. Since they need to be produced after inflation, the reheating temperature cannot be much lower than their mass. In supersymmetric scenarios, this may be in conflict with the upper bound on the reheating temperature necessary to avoid the overproduction of gravitinos during reheating, which may spoil the successful predictions of BBN. On the other hand, if the RH neutrinos are nearly degenerate in mass, the self-energy contribution to the CP asymmetries may be resonantly enhanced, thus making leptogenesis viable at temperatures as low as TeV. This interesting situation is called “resonant leptogenesis” 3.

In order to precisely quantify the lepton asymmetry generated by the leptogenesis mechanism, one needs to keep track of the abundances of the particles involved in the process by solving a set of coupled Boltzmann equations. The standard calculations employ a set of semi-classical equations. However, quantum Boltzmann equations for leptogenesis have been recently derived 4 (for an earlier study, see Ref. 5), using a Green’s function technique known as Closed Time-Path (CTP) — or Schwinger-Keldysh — formalism 6, which provides a complete description of non-equilibrium phenomena in field theory. While in the semi-classical setup every scattering in the plasma is independent of the previous one, in a full quantum approach the whole dynamical history of the system is taken into account. The quantum Boltzmann equations describe therefore a non-Markovian dynamics, manifesting the typical “memory” effects which are observed in quantum transport theory 7. The thermalization rate obtained from the quantum transport theory may be substantially longer than the one obtained from the classical kinetic theory.

Furthermore, and more importantly, the CP asymmetry turns out to be a function of time, even after taking the Markovian limit. Its value at a given instant depends upon the previous history of the system. If the timescale of the variation of the CP asymmetry is shorter than the relaxation time of the particles abundances, the CP asymmetry may be averaged over many scatterings and it reduces to its classical constant value; the solutions to the quantum and the classical Boltzmann equations are expected to differ only by terms of the order of the ratio of the timescale of the CP asymmetry and the relaxation timescale of the particle distributions. In thermal leptogenesis with hierarchical RH neutrinos this is typically the case. However, in the resonant leptogenesis scenario, where at least two RH neutrinos are almost degenerate in mass and their mass difference $\Delta M$ is of the order of their decay rates, the typical timescale to build up coherently the CP asymmetry (of the order of $1/\Delta M$) can be larger than the timescale for the change of the abundance of the RH neutrinos. Thus, in the case of resonant leptogenesis significant differences are expected between the classical and the quantum approach.

2 Quantum Boltzmann equations

The model I consider consists of the SM plus three RH neutrinos $N_i$ ($i = 1, 2, 3$), with Majorana masses $M_1 \leq M_2 \leq M_3$. The interactions among RH neutrinos, Higgs doublets $H$, lepton doublets $\ell_\alpha$ and singlets $e_\alpha$ ($\alpha = e, \mu, \tau$) are described by the Lagrangian

$$\mathcal{L}_{\text{int}} = \lambda_{i\alpha} N_i \ell_\alpha H + h_\alpha \bar{e}_\alpha \ell_\alpha H^c + \frac{1}{2} M_i N_i^2 + \text{h.c.},$$  

(1)

with summation over repeated indices. In the early universe, the quantum numbers conserved by sphaleron interactions are $\Delta_\alpha = B/3 - L_\alpha$, where $B, L_\alpha$ are the baryon asymmetry and the lepton asymmetry in the flavour $\alpha$, respectively.

The quantum Boltzmann equations describing the generation of the baryon asymmetry are obtained using the CTP formulation of non-equilibrium quantum field theory. The reader is referred to Ref. 4 for the technical details of the calculation. Here, I only summarize the main results. After taking the Markovian limit, the equations for the number densities of RH neutrinos
\[ Y_{N_i} \text{ and the asymmetries } Y_{\Delta \alpha} \text{ (per entropy density) read} \]
\[
\frac{dY_{N_i}}{dz} = -D_i \left( Y_{N_i} - Y_{N_i}^{eq} \right), \tag{2}
\]
\[
\frac{dY_{\Delta \alpha}}{dz} = -\sum_i \epsilon_{\alpha i} D_i \left( Y_{N_i} - Y_{N_i}^{eq} \right) - W_{\alpha} |A_{\alpha \alpha}| Y_{\Delta \alpha}, \tag{3}
\]

where the ratio between the mass of the lightest RH neutrino and the temperature \( z = M_1/T \) plays the role of the time variable. At equilibrium the \( N_i \) number density normalized to the entropy density of the universe is given by \( Y_{N_i}^{eq} = \frac{z_i^2 K_2(z_i)}{(4\pi^3)} \), where \( z_i = z\sqrt{x_i}, \ x_i = (M_i/M_1)^2, \ g_* = 106.75 \) and \( K_\alpha(z_i) \) is a modified Bessel function of the \( n \)-th kind. The decay and washout terms appearing in (2)-(3) are defined as
\[
D_i = K_i x_i z \frac{K_1(z_i)}{K_2(z_i)}, \quad W_{\alpha} = \sum_i \frac{1}{4} K_{i\alpha} \sqrt{x_i} K_1(z_i) z_i^3, \tag{4}
\]

where the washout parameters are given by the ratios between the decay rates and the Hubble parameter
\[
K_{i\alpha} = \frac{\Gamma(N_i \rightarrow \bar{\ell}_\alpha H)}{H(T = M_i)}, \quad K_i = \sum_{\alpha} K_{i\alpha}. \tag{5}
\]

The form of the matrix \( A \) depends on the number of lepton flavours which are effective in the dynamics of leptogenesis, and this in turn depends on the temperature at which leptogenesis takes place, which is roughly given by \( M_1 \). Indeed, for \( M_1 \gtrsim 10^{12} \text{ GeV} \) all lepton flavours are not distinguishable and the one-flavour regime holds; for \( 10^9 \text{GeV} \lesssim M_1 \lesssim 10^{12} \text{GeV} \) and \( M_1 \lesssim 10^9 \text{GeV} \), two and three lepton flavours become effective, respectively. For example, in the approximation where \( A \) is a diagonal matrix \( A = -\text{diag}(151/179, 344/537, 344/537) \), for \( M_1 \lesssim 10^9 \text{GeV} \). The complete expressions can be found in Refs.\(^8\). Finally, sphaleron interactions introduce a conversion factor for the final baryon asymmetry
\[
Y_B = \frac{12}{37} \sum_{\alpha} Y_{\Delta \alpha}(z \to \infty). \tag{6}
\]

The key quantities controlling the production of a net lepton number are the CP asymmetries in the \( N_i \) decays
\[
\epsilon_{\alpha i} = \frac{\Gamma(N_i \rightarrow \bar{\ell}_\alpha H) - \Gamma(N_i \rightarrow \bar{\ell}_\alpha H)}{\Gamma(N_i \rightarrow \bar{\ell}_\alpha H) + \Gamma(N_i \rightarrow \bar{\ell}_\alpha H)}. \tag{7}
\]

The Eqs. (2)-(3) reproduce exactly the usual Boltzmann equations obtained in the semiclas-sical approach, except for a crucial difference in the source term of (3). As mentioned above, the inclusion of quantum effects introduces a time dependence in the CP asymmetry
\[
\epsilon_{\alpha i}(z) = \sum_{j \neq i} \epsilon_{\alpha i}^{(ij)} m^{(i,j)}(z), \tag{8}
\]
\[
m^{(i,j)}(z) = 2 \sin^2 \left( \frac{1}{2} \frac{M_j - M_i}{2H(M_1)^2} z^2 \right) - \frac{\Gamma_j}{M_j - M_i} \sin \left( \frac{M_j - M_i}{2H(M_1)^2} z^2 \right), \tag{9}
\]
\[
\epsilon_{\alpha i}^{(ij)} = \frac{1}{8\pi} \text{Im} \left[ \lambda_{\alpha i}^* \lambda_{\alpha j} (\lambda \lambda^*)_{ij} \right] \left( g_{\alpha}^{(ij)} + g_{\nu}^{(ij)} \right), \tag{10}
\]
\[
g_{\alpha}^{(ij)} = \sqrt{\frac{x_j}{x_i}} \frac{1}{1 - \frac{x_j}{x_i}} \left( 1 + \frac{1 + x_j/x_i}{x_j/x_i} \right), \tag{11}
\]
\[
g_{\nu}^{(ij)} = \sqrt{\frac{x_j}{x_i}} \left( 1 - \left( 1 + \frac{x_j}{x_i} \right) \ln \frac{1 + x_j/x_i}{x_j/x_i} \right), \tag{12}
\]
where $\Gamma_j = \sum_\beta \Gamma(N_j \rightarrow \ell_\beta \bar{H}) = (\lambda \lambda^\dagger)_{jj} M_j / (8\pi)$ is the total decay rate of the $j$-th RH neutrino, $g_s$ and $g_v$ are the self-energy and the vertex correction functions, respectively.

In the quantum approach, the typical timescale for the variation of the CP asymmetry is

$$t = \frac{1}{2H(T)} = \frac{z^2}{2H(M_1)} \sim \frac{1}{M_j - M_i} = \frac{1}{\Delta M_{ji}}. \quad (13)$$

If the timescale for the variation of the particle abundances $1/\Gamma_j$ is much larger than $1/\Delta M_{ji}$, the CP asymmetry will average to its classical value $\bar{\epsilon}_{\alpha} = \sum_{i\neq j} \epsilon_{\alpha}^{(i,j)}$ and no significant quantum effect arises. On the other hand, quantum effects are expected to be sizeable if the timescale $1/\Delta M_{ji}$ for building up coherently the CP asymmetry is larger than the timescale $1/\Gamma_i$ for changing the abundances. In other words, the oscillation frequency $\Delta M_{ji}$ has to be sufficiently smaller than $\Gamma_i$, so that the factors $m_{\alpha}^{(i,j)}(z)$ do not effectively average to one. Under these conditions, the amplitude of the “sin” term in $m_{\alpha}^{(i,j)}(z)$ is also enhanced, which turns out to be a crucial effect.

The above discussion allows one to formulate a quantitative criterion for the importance of quantum effects, namely $\Delta M_{ji} \lesssim \Gamma_i$. This criterion can be naturally satisfied if RH neutrinos are nearly degenerate, such as in resonant leptogenesis and in models based on Minimal Lepton Flavour Violation.

### 3 Application to resonant leptogenesis

As anticipated in the Introduction, resonant leptogenesis relies on the fact that the CP asymmetries are resonantly enhanced when the mass differences among RH neutrinos are comparable to their decay widths. $\Delta M_{ji} \sim \Gamma_i \sim \Gamma_j$. Therefore, the criterion for the significance of the quantum effects is satisfied and one expects the quantum Boltzmann equations to provide appreciably different results with respect to their semi-classical counterparts.

For the sake of simplicity, let us restrict here to the simplest case where there are only two RH neutrinos with mass difference $\Delta M$ and similar decay rates $\Gamma_N$, in the one-flavour approximation where $\alpha$ has a single value. In this case, the Boltzmann equation for the lepton asymmetry in the single flavour $\alpha$ contains the CP asymmetry

$$\epsilon_{\alpha}(z) \simeq \bar{\epsilon}_{\alpha} \left[ 2 \sin^2 \left( \frac{K z^2 \Delta M}{4 \Gamma_N} \right) - \frac{\Gamma_N}{\Delta M} \sin \left( \frac{K z^2 \Delta M}{2 \Gamma_N} \right) \right], \quad (14)$$

where $\bar{\epsilon}_{\alpha} = \sum_i \sum_{j \neq i} \epsilon_{\alpha}^{(i,j)}$ is the constant value of the asymmetry in the classical limit, and $K$ is the total washout parameter.

The plot in Fig. 1 shows the absolute values of the final lepton asymmetry computed with and without the time-dependent factor in (14). For the strong washout regime $K \gtrsim 1$, the quantum and semi-classical methods give almost the same answers; instead, at small $K$, the discrepancy between the two approaches is sizeable, of an order of magnitude or more. This is easily understood from Eq. (14), where at large $K$ the “sin” functions average to constants while in the opposite case they determine an appreciable time-dependence of the CP asymmetry.

In Ref. 10, the reader can find a more detailed study of the more general “flavoured” case as well as analytical approximations for the lepton asymmetries in the different possible washout regimes. They reproduce fairly well the numerical solutions of the Boltzmann equations (2)-(3).

#### 3.1 MLFV leptogenesis

Nearly degenerate RH neutrinos naturally arise in the context of models satisfying the hypothesis of Minimal Flavour Violation (MFV)\textsuperscript{11,12}. In the quark sector, where the MFV hypothesis has
been formulated first, the MFV ansatz states that the two quark Yukawa couplings are the only irreducible breaking sources of the flavour-symmetry group defined by the gauge-kinetic lagrangian\textsuperscript{11}. In generic models satisfying this hypothesis, quark Flavour Changing Neutral Currents are naturally suppressed to a level comparable to experiments and new degrees of freedom can naturally appear in the TeV range. The extension of the MFV hypothesis to the lepton sector (MLFV) has been formulated in Ref.\textsuperscript{12}, where a number of possible scenarios, depending on the field content of the theory, have been identified. The case more similar to the quark sector and more interesting from the point of view of leptogenesis is the so-called extended field content scenario. The lepton field content is extended by three heavy RH neutrinos with degenerate masses at the tree level. The largest lepton flavour symmetry group of the gauge-kinetic term is $SU(3)_c \otimes SU(3)_e \otimes O(3)_N$ and, according to the MLFV hypothesis, it is assumed that this group is broken only by the charged-lepton and neutrino Yukawa couplings $h_\alpha$ and $\lambda_{i\alpha}$. In relation to leptogenesis, the key feature of this scenario is that the degeneracy of the RH neutrinos is lifted only by corrections induced by the Yukawa couplings, so that we end up with a highly constrained version of resonant leptogenesis.

Within this setup, the viability of leptogenesis has been considered either in the one-flavour approximation\textsuperscript{13} or in the flavoured case\textsuperscript{14}. However, in the light of the quantum version of the Boltzmann equations I discussed here, it turned out necessary to carry out an analysis to assess the impact of the quantum effect in this MLFV leptogenesis scenario. It has been shown in\textsuperscript{15}, both analytically and numerically, that neglecting the time dependence of the CP asymmetry may underestimate the baryon asymmetry by several orders of magnitude when a strong degeneracy among heavy RH neutrinos and small mass splittings in the light neutrino sectors are present. This is true both when the CP phases come from the RH sector and when they come entirely from the left-handed sector and may be identified with the low energy PMNS phases.

4 Conclusions

The simplest see-saw framework, where RH neutrinos are added to the particle content of the SM, may simultaneously account for the small neutrino masses and the baryon asymmetry of the universe, through the leptogenesis mechanism. Obtaining detailed results in leptogenesis requires solving the Boltzmann equations for the abundances of the particles involved. In this talk, I have presented a set of quantum Boltzmann equations which has been derived using non-
equilibrium quantum field theory. The main difference with respect to the usual semi-classical equations is that the CP asymmetry is time-dependent. A criterion to discriminate situations where one should expect quantum effects to be important has been discussed. In particular, this condition is satisfied in realistic models such as resonant and MLFV leptogenesis. In these scenarios, quantum effects play a significant role and should be taken into account.

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References

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The CRESST-II experiment in the Laboratori Nazionali del Gran Sasso uses scintillating crystals as a target to search for elastic scatterings of dark matter particles in a laboratory environment. The detectors are operated in a dilution cryostat at temperatures below 30 mK, and for each particle interaction, the phonon signal as well as the scintillation light signal are recorded. The current limit that can be placed on the spin-independent WIMP-nucleon scattering cross-section with this technique is below $6 \times 10^{-7}$ pb for WIMPs in the mass range from about 40 to 90 GeV/c$^2$.

1 Introduction

It remains one of the most pressing problems of physics today to clarify the nature of dark matter. So far, the evidence for a significant component of non-baryonic dark matter in the universe is compelling. It stems from a variety of independent experiments, ranging from measurements of the cosmic microwave background$^1$ and big bang nucleosynthesis$^2$ to astrophysical observations of large amounts of unseen matter in galaxy clusters$^3$ and galaxies$^4$. Yet, the nature of this component is still unknown, so experiments to unravel this mystery are more important than ever.

Highly motivated candidates for dark matter are weakly interacting, yet massive particles (WIMPs)$^5$. These particles are expected to be gravitationally bound in the Milky Way in a roughly isothermal halo, thus following a Maxwell-Boltzmann-distribution$^6,7$. The WIMP density at the position of the Earth is assumed to be around 0.3 GeV/cm$^3$. On our way around
the Galaxy we pass through this halo with a velocity of about 220 km/s, and we hope to detect the scattering of this WIMP wind on nuclei in an absorber.

Accelerator searches for new forms of matter place a lower bound on the mass of these particles of the order of 10 GeV/c². Since the de Broglie wavelength of particles with such mass and velocity is larger than the radius of a nucleus, one may expect the scattering to occur in a coherent way, the case we consider here. Then, for a given target material with mass number A, the scattering amplitude scales as $A^2$, favoring heavy nuclei for a search experiment.

2 Scintillating Crystals

Following the interaction of a WIMP with a target nucleus, the energy of the recoil is very small, typically only of the order of 10 keV. Together with the very small expected interaction rate of less than 10 events per kilogram of target and year of measuring time, available technologies are highly constrained. So the two main requirements to this kind of experiment are a very low energy threshold and the capability to reject backgrounds caused by known particle species.

Figure 1: A detector module as used in the CRESST-II experiment. Left picture: An open module with the light absorbing wafer on the left hand side and the scintillating crystal in its housing on the right hand side. Sketch: a tungsten superconducting phase transition thermometer is evaporated onto a CaWO₄ target crystal. A second thermometer is evaporated on a light absorber to measure the scintillation light. Both detectors are enclosed in a scintillating reflective housing.

In the second phase of the Cryogenic Rare Event Search with Superconducting Thermometers CRESST-II⁸ we use scintillating crystals as target material⁹. They are shaped as cylinders of $(4 \times 4$) cm and weigh about 300g each. We operate them as calorimeters at temperatures as low as 10 mK. In such dielectric crystals, most of the energy of a particle interaction goes into phonons. To measure these phonons, we evaporate a thin tungsten film onto the crystal. Tungsten becomes superconducting in this temperature range, and stabilizing the film in its superconducting phase transition (by means of a dedicated heater) makes an extremely sensitive thermometer: Any particle interaction warms up the film, thus changing its resistance, which can be measured with sensitive electronics¹⁰.

In addition to this phonon signal, a small fraction of the interaction energy (typically a few percent) is emitted as scintillation light. To detect this light, we use another tungsten phase transition thermometer on a separate light absorbing wafer. Thus, for each target crystal, we have two thermometers, one measuring the deposited energy, and the other measuring the scintillation light, see figure 1.

We define the light yield of an interaction as the amount of energy in the light detector divided by the energy in the phonon detector, and normalize it such that electron recoils have
a light yield of 1. Electron recoils are caused by electrons and gammas that impinge onto the crystal. Compared to such events, the light yield of alpha particles is reduced by a factor of 5. Neutrons are mainly seen when they scatter from oxygen due to the kinematics of the interaction, with a light yield reduced by a factor of 10 relative to that of the electron recoils. Coherent scatterings of WIMPs are expected to take place mainly on tungsten where the light yield is reduced by a factor of 40\(^{11}\). Thus, simultaneously measuring the light signal allows us to discriminate the (possibly WIMP induced) tungsten recoils from the dominant radioactive backgrounds.

3 The Upgraded Setup

In 2007, the experimental setup was extended to be capable of housing up to 33 detector modules. To shield them as much as possible from ambient radioactivity, we provide a variety of shielding layers. The experiment is hosted in the Laboratori Nazionali del Gran Sasso, Italy, to shield the detectors from cosmic ray induced backgrounds with the overburden of 1400 m of rock. During the upgrade, we added a muon veto to discriminate residual muon induced backgrounds, as well as a 45 cm thick wall of polyethylene to moderate neutrons to energies below detection threshold. The setup is constantly flushed with nitrogen vapor, in particular to keep radon contaminated air away from the inner parts. A 20 cm thick lead shield and a 14 cm copper shield absorb gamma radiation coming from outside. Also, the fivefold thermal shielding of the cryostat provides an additional 1.2 cm thick copper shield. All the materials in the vicinity of the detectors are selected for radiopurity and handled in a clean room environment.

![Figure 2: Setup of CRESST-II. The detectors are in a low-background environment in the center of the shielding. The cryostat (upper part of the figure) is kept away from this very clean environment. The polyethylene neutron shield is shown in yellow, the lead shield in gray, and the copper shield in orange. A muon veto made from plastic scintillator, shown in blue, helps to further reduce backgrounds induced by cosmic rays.](image-url)
4 Results

After a major upgrade we installed 10 detector modules and cooled the cryostat for commissioning of the new setup. The detectors were calibrated with a gamma source to set the energy scale, as well as a neutron source for validation of the light yield anticipated for nuclear recoils. The data reported here was taken in 2007 with two detector modules from March 27th through July 23rd.

We fit the pulses using a template fit, which is constructed from pulses from the gamma calibrations. In order to treat noise in an unbiased way, we allow the amplitude of the fit to take negative values. This guarantees that for no signal, the reconstructed amplitudes indeed scatter around zero with a width corresponding to the noise. Thus negative amplitudes can arise, resulting in events with negative light yields.

We perform only very basic quality cuts on the data, rejecting pulses only if they occur in one of the rare periods when the temperature of the cryostat is not as stable as desired, if they are direct hits of the light detector or the thermometer on the crystal, or if they are pile-up events. Events recorded from the two detectors with a cumulative exposure of 50 kg d are shown as scatter plot in figure 3.

![Figure 3: Low energy event distribution measured with two 300 g CaWO₄ detector modules during the commissioning run in 2007. The vertical axis represents the light yield expressed as the ratio (energy from the light channel)/(energy from the phonon channel), the horizontal axis is the total energy measured by the phonon channel. Below the red curve we expect 90% of all nuclear recoils, and below the black curve 90% of the tungsten recoils. The acceptance window is set below the black line and between 11 keV (above which we can discriminate electron recoils) and 40 keV (below which most of the WIMP induced recoils appear).]
5 Discussion

To derive a limit on the coherent WIMP-nucleus scattering cross section, standard assumptions on the dark matter halo are adopted. The finite extension of the nuclei is taken into account by assuming the Helm form factor, which basically limits the energy transfer to the tungsten nuclei to energies below 40 keV for all WIMP masses. In the energy region above 11 keV, where recoil discrimination becomes efficient, up to 40 keV, 4 tungsten recoil events were observed in the data of figure 3. Combining this data with data from the previous run, the upper limit for the WIMP scattering cross-section per nucleon is set using Yellin’s optimum interval method, shown as the red curve in figure 4. The minimum of this curve is below $6 \times 10^{-7}$ pb for WIMPs with masses between 40 and 90 GeV/c$^2$, obtained after a gross exposure (including down times due to refilling of cryogenic liquids etc.) of only 67 kgd.

![Figure 4: Exclusion limits on the WIMP-nucleon spin independent scattering cross section from a few experiments, from top to bottom from the KIMS experiment (green), EDELWEISS (yellow), XENON10 (blue) and CDMS (violet). The red curve is the limit from the CRESST experiment, derived from the data of figure 3 and that of the previous run. In cyan two theoretical expectations.](image)

The few events that we observe in our signal area need not be WIMP induced tungsten recoil events. During the commissioning, a weak point in the neutron shielding above the muon veto was identified, but patched only after data taking was completed. This background can be estimated to account for the observed number of events in the nucleon recoil band. In addition we might have a small neutron background induced by decays from radioactive contaminations in the other crystals that were not running during this data taking period, a situation unique to this commissioning phase. The upcoming run should clarify on these points.

In the hunt for a discovery one might want to compare the different technologies, and since exclusion limits are one way of doing so, figure 4 also shows the results from a few other experiments: From the KIMS experiment using a gross exposure of 3409 kgd on CsI(Tl), the EDELWEISS experiment ($\approx 180$ kgd, Ge), XENON10 ($1980$ kgd, Xe) and CDMS ($\approx 1250$ kgd, Ge). The figure also contains two expectations from theory.
References

Recent results in very high energy gamma ray astronomy

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The very high energy (E > 50 GeV) gamma ray astronomy is an emerging field. Three major Imaging Atmospheric Cherenkov Telescope collaborations, HESS, MAGIC and VERITAS are presently in operation. Many new results in astroparticle physics have been obtained by these instruments. In this paper, Galactic Center observations, dark matter searches and limits on Lorentz invariance violation are reported.

1 Introduction

The very high energy (VHE) gamma ray astronomy is a new emerging field. This paper focuses on selected results from the VHE gamma ray instruments which are of interest to the particle physics community. The first section reviews the experimental status. Next, recent results on the Galactic Centre from the HESS collaboration are reported. These results constrain the source of the VHE gamma ray emission. The next section is dedicated to indirect dark matter searches. More than 15 Active Galactic Nuclei (AGN) have been observed in the TeV regime. These distant variable sources allow to test for Lorentz invariance breaking. The final section reports constraints on violations of the Lorentz invariance obtained by the MAGIC collaboration.

2 The very high energy gamma ray instruments

Ground based very high energy gamma ray instruments detect atmospheric cascades initiated by astrophysical gamma-rays. These instruments are sensitive to photons in the range 50 GeV to 100 TeV. The two major classes of instruments are imaging atmospheric Cherenkov telescopes (IACT) and shower particle detectors. The characteristics of major currently operating IACTs are summarized in table 1. These instruments are the HESS, VERITAS and CANGAROO-III arrays and the MAGIC telescope. Each of their telescope is equipped with a finely pixelized camera at its focus. A large fraction of the cosmic ray background is rejected with the analysis of the Cherenkov image. In addition, Cherenkov arrays use stereoscopy to improve their background rejection and their energy/angular resolution.
Table 1: Principal characteristics of currently operating major IACTs

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<th>Instrument</th>
<th>lat. (deg)</th>
<th>long (deg)</th>
<th>Altitude (m)</th>
<th>Tels</th>
<th>Telescope Area (m²)</th>
<th>Camera Pixels</th>
<th>FOV (deg)</th>
<th>Threshold (TeV)</th>
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<td>4</td>
<td>107</td>
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<td>18</td>
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<td>1</td>
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<td>574</td>
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<td>160</td>
<td>3</td>
<td>57.3</td>
<td>427</td>
<td>4</td>
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</tr>
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</table>

HESS\(^a\) is an array of 4 12-meter diameter telescopes located in the Khomas highlands of Namibia. HESS was completed at the end of 2003. The typical angular resolution of the HESS instrument of 0.1 deg. allows to make detailed images of extended galactic objects such as the supernova remnant RXJ1713-3946\(^b\). The HESS collaboration is performing a galactic survey of TeV sources\(^7\) which takes advantage of both the angular resolution and the large 5 deg field of view (FOV) of the cameras. A fifth telescope with a 28-meter diameter is being built and should be completed in 2009.

The VERITAS\(^b\) array is an upgrade of the 10-meter Whipple telescope. The Whipple telescope was a major IACT of the 1980s and 1990s. VERITAS is very similar in characteristics to the HESS array and was completed in 2007.

The 17 meter diameter MAGIC\(^c\) telescope is the largest operating IACT. It is located on La Palma island in the Canaries and started its operations in 2004. It is optimised for low-energy photons (\(E \approx 50\,\text{GeV}\)) detection. The MAGIC collaboration is building a second similar telescope which should be completed in 2008. This second telescope will allow the MAGIC collaboration to use the stereoscopy technique.

The MILAGRO\(^d\) instrument is a shower particle detector. It is a water Cherenkov detector located at Los Alamos, at an altitude of 2630 meters. It is in operation since 2000. It has a large FOV (2 sr) which compensates for the high trigger threshold (\(\sim 1\,\text{TeV}\)) and the angular resolution of the order of 1 deg. MILAGRO has detected extended emission from the galactic plane (Abdo et al (2008)\(^2\)) and 4 “point sources” at a significance > 4\(\sigma\) (Abdo et al (2007)\(^1\)).

3 Galactic Centre

VHE gamma ray emission from the Galactic Centre has been reported in the TeV range by the CANGAROO, HESS, Whipple and MAGIC collaborations. The published\(^3\) position of the TeV source HESS J1745-290 is located within 20 arc seconds of the central black hole Sgr A*. The spectrum, which was obtained from 2004-2005 data is compatible with a pure power-law with spectral index \(\Gamma = 2.29 \pm 0.04\text{(stat)} \pm 0.1\text{(syst)}\) in the range 100 GeV to 10 TeV. The possible non-standard interpretation of the HESS J1745-290 signal as annihilation of dark matter particles has been investigated by the HESS collaboration (F.Aharonian et al.(2006))\(^3\). The observed spectrum is not well fitted by expected dark matter annihilation spectra, implying a mostly non-dark matter origin for the signal. Assuming that the observed signal is a blend of an astrophysical source and dark matter annihilations in a Navarro-Frank-White halo, the 95 % C.L. upper limits on the velocity-weighted cross-section are \(< \sigma v > \approx 10^{-24}\text{cm}^3\text{s}^{-1}\) for a neutralino mass in the range 100 GeV to 1 TeV.

The signal of HESS J1745-290 can be interpreted by a large variety of astrophysical models.

\(^{a}\)Web address http://www.mpi-hd.mpg.de/hfm/HESS/HESS.html
\(^{b}\)Web address http://veritas.sao.arizona.edu
\(^{c}\)Web address http://wwwmagic.mppu.mpg.de
\(^{d}\)Web address http://www.lanl.gov/milagro
Figure 1: Left panel: VLA image of the Galactic Centre region. The extended source on the left is the supernova remnant Sgr A East. The outer circle shows the error on the position of HESS J1745-290 from F.Aharonian et al (2006). The inner circle is the error on the position of the same source from C.VanEldik et al (2007). Right panel: Simultaneous Chandra and Hess observation on July 30 2005, from J.Hinton et al (2007). Top: Chandra 1-10 keV count rate (400 seconds bins). Bottom: HESS light curve in 15 minutes bins.

(see e.g. Hinton and Aharonian (2007)\textsuperscript{14} and references inside). Possible astrophysical sources for HESS J1745-290 include Sgr A*, the supernova remnant Sgr A East and the pulsar wind nebula G359.95-0.04. Recently, a careful study\textsuperscript{20} allowed to lower the pointing errors of the HESS experiments down to the level of 8 arc seconds. The preliminary position of the source is located at an angular distance of 7.3'' ± 8.7'' (stat) ± 8.5'' (syst) from the central galactic black hole Sgr A*. This position (see figure 1) is incompatible with the centroid of the radio emission of the supernova remnant Sgr A East, but still compatible with sources such as the pulsar wind nebula G359.95-0.04.

The Sgr A* black hole is well known to be a variable source in the infrared and X-ray passbands. The variability of the HESS J1745-290 source has been studied in M.Vivier et al.\textsuperscript{21}. No significant variability or periodicity was found between 30 seconds and one year. Simultaneous data were also taken with the Chandra satellite during a flare of SgrA* on July 30 2005 (see figure 1 and J.Hinton et al. (2007)\textsuperscript{15}). No significant TeV flare was seen during the X-Ray flare. This implies that the TeV emission is produced at a relatively large distance for the Sgr A* black hole.

4 Indirect dark matter searches

Popular particle physics models such as the Minimal Supersymmetric Standard Model (MSSM) or Universal Extra Dimensions ("Kaluza-Klein")\textsuperscript{18} predict WIMP (Weakly Interacting Massive Particles) dark matter annihilations in galactic halos. These annihilations could give observable signals in Cherenkov telescopes (for a review see Bertone, Hooper and Silk (2005)\textsuperscript{12}). The flux \(\frac{d\Phi}{dE_\gamma}\) of gamma rays is

\[
\frac{d\Phi}{dE_\gamma} = \frac{1}{4\pi} \frac{dN_\gamma < \sigma v >}{dE_\gamma} \frac{1}{M_X^2} J \Delta \Omega.
\]  

(1)
It is the product of an astrophysical term $\bar{J}$ and a particle physics term. The former depends on the mass density profile $\rho$ of the dark halo

$$\bar{J} = \langle \int_{\text{L.o.s.}} \rho^2 ds \rangle .$$

In equation 2, the average is taken over the solid angle $\Delta \Omega$ spanned by the Point Spread Function (PSF). The spatial resolution of H.E.S.S. is of the order of 5 arc minutes per event, giving $\Delta \Omega = 2 \times 10^{-5}$. The particle physics term depends on the velocity averaged annihilation cross section $< \sigma v >$ and the WIMP mass $M_\chi$.

The possible targets for WIMP annihilation searches can be ranked according to their values of $\bar{J}$. If the annihilation signal from a halo located at distance $D$ is “point-like”, then $\bar{J} \propto M^2/D^5$ where $M$ is the (often measured) dark mass inside the PSF. The best astrophysical targets are thus the Galactic Center and nearby dwarf galaxies. The results of the Galactic center search have already been mentioned (section 3). Data towards several dwarf galaxies have been taken by the Atmospheric Cherenkov collaborations. These galaxies are Sagittarius (HESS$^5$), Draco (MAGIC$^9$ and Whipple$^{22}$) and Ursa minor (Whipple$^{22}$). The expected flux from galaxy clusters such as Virgo or Coma is smaller by at least 3 orders of magnitude. It is also possible to look for dark matter clumps or dark matter annihilations in the vicinity of Intermediate Mass Black Holes$^4$.

The Sgr dwarf galaxy is a satellite of the Milky Way. It is located in the galactic plane in the direction of the Galactic Center, at a distance of 24 kpc. It is being torn apart by the tidal force of the Galaxy. The visible mass profile of the Sgr dwarf galaxy is difficult to obtain because of the contamination of galactic foreground stars. The center of the Sgr dwarf galaxy is coincident with the globular cluster M 54. The interpretation of velocity dispersion measurements is difficult because of the tidal interaction with the Milky Way. The central velocity dispersion has been measured by several groups (see e.g. Zaggia et al (2004)$^{23}$). The HESS collaboration (F.Aharonian et al (2008)$^5$) choose to describe Sagittarius galactic structure with two models. The first one is a Navarro-Frank-White model with parameters taken from Evans, Ferrer and Sarkar (2004)$^{13}$. The other model (“core model”) was fitted to the structural parameters (distribution of visible mass and central dispersion).

The Sgr dwarf galaxy has been observed by H.E.S.S. in June 2006. After quality cuts, a total exposure of 11 hours was obtained. No significant excess is seen at the position of M54. This translates into 95% C.L. upper limits of $< \sigma v > \approx 10^{-23}\text{cm}^3\text{s}^{-1}$ for the NFW model and $< \sigma v > \approx 2 \times 10^{-25}\text{cm}^3\text{s}^{-1}$ in the “core model”. These upper limits are valid in the 100 GeV - 1 TeV neutralino mass range.

The distance to the Draco dwarf galaxy is 80 kpc, more than 3 times the distance to the Sgr dwarf galaxy. On the other hand, the galactic structure of the Draco dwarf galaxy has been studied in details (see e.g Strigari et al (2007)$^{19}$, resulting in much smaller error bars on the astrophysical factor $\bar{J}$. The MAGIC and Whipple collaboration have published results on the Draco dwarf galaxy. The MAGIC collaboration has analyzed 7.8 hours of their data. They reach a sensitivity of $< \sigma v > \approx 10^{-22}\text{cm}^3\text{s}^{-1}$ for neutralino masses in the range 140 GeV - 500 GeV. The Whipple collaboration reaches a sensitivity of $< \sigma v > \approx 10^{-21}\text{cm}^3\text{s}^{-1}$ with 14.3 hours, for a neutralino mass in the range 500 GeV - 2 TeV. They have also analyzed 17.2 hours of data towards the Ursa Minor dwarf galaxy. Their Ursa Minor upper limit is worse than their Draco limit by a factor of 2.

5 Bounds on Lorentz invariance violation

More than 15 AGN have been detected in the TeV regime. Most of them have jets directed along the line of sight (they are member of the so-called “blazar” class). The only known exception
is the AGN associated to the M87 galaxy. This AGN has been observed in the TeV regime by the HEGRA, HESS, and VERITAS arrays.

The variability of the VHE emission of AGN can be used to test for Lorentz invariance in photon propagation. This is motivated by several quantum gravity theories (see Sarkar (2002)\(^\text{(17)}\) for a review) in which photons and neutrinos are expected to have an energy dependent velocity in vacuum. The photon velocity \(v(E)\) is parametrized by either a linear

\[
v(E) = 1 - \eta \left( \frac{E}{M_1} \right)
\]

or quadratic

\[
v(E) = 1 - \eta \left( \frac{E}{M_1} \right)^2
\]

ton of the energy \(E\). \(M_{1,2}\) is the scale of Lorentz invariance breaking. For quantum gravity theories, \(M_{1,2}\) are expected to be of the order of the Planck scale.

An experiment observing two photons with an energy difference of \(\Delta E\), arriving from a source at distance \(L\) (or redshift \(z = L/500\text{Mpc}\)) in a burst of duration \(\Delta t_{\text{burst}}\) is sensitive to a Lorentz invariance breaking scale

\[
M_1 = \frac{L \Delta E}{c \Delta t_{\text{burst}}} = 10^{18}\text{GeV} \left( \frac{z}{0.1} \right) \left( \frac{\Delta E}{1\text{TeV}} \right) \left( \frac{60\text{s}}{\Delta t_{\text{burst}}} \right)
\] (5)

Equation 5 shows that Lorentz invariance breaking effects can be observed in short TeV photon bursts from AGN. The MAGIC collaboration has observed flares from Mkn501\(^\text{(10)}\), an AGN at a redshift of \(z=0.034\). A very intense flare, with a VHE photon flux of more than 3.5 times the flux of the Crab nebula occurred on July 9 2005 (figure 3). The outburst lasted 15 minutes, with a flux doubling time of \(\sim 2\) minutes. The time of the flare maximum \(t_{\text{max}}\) was observed to depend on the energy. The slope of the time delay as a function of the energy is \(\tau_1 = 0.030 \pm 0.012\text{s/GeV}\), in the case of a linear dependence. In other words, there is a positive 2\(\sigma\) positive detection, (explainable by the astrophysical emission process) which gives\(^\text{11}\) a lower limit on \(M_1 > 0.26 10^{18}\ \text{GeV} (95\% \text{ C.L.})\). A similar analysis gives a lower limit on \(M_2 > 0.39 10^{11}\ \text{GeV} (95\% \text{ C.L.})\).

The HESS collaboration has observed an intense flare\(^\text{8}\) (more than 10 times the flux of the Crab nebula) from the blazar PKS2155-304, at a redshift of \(z=0.116\). The flare, shown on figure 3, occurred on July 28 2006, was roughly 1 hour long and composed of at least 5 smaller

![Figure 2: Results for HESS dark matter annihilations search towards the Sagittarius dwarf galaxy (Aharonian et al (2008)). Left panel: 95 % C.L. exclusion limits on MSSM models from HESS searches towards the Sgr dwarf galaxy. Right panel: 95 % C.L. exclusion limits on the Universal Extra Dimensions model of Servant and Tait.](image-url)
outbursts. The shortest rise time was measured to be $\Delta t_{\text{rise}} = 173 \pm 28$ s. An analysis similar to the MAGIC analysis of Mkn501 is undergoing to improve their bounds on Lorentz invariance breaking scales.

6 Conclusion and perspectives

The VHE instruments, mostly IACT, have published a number of interesting new results. The Galactic Center source, HESS J1745-290 is probably not dark matter annihilations. It is not associated to the Sgr A East remnant. If associated to Sgr A*, the TeV emission is produced at a larger distance from the black hole than the X-ray emission. New results on indirect dark matter searches include limits from HESS towards Sgr dwarf, MAGIC and Whipple towards Draco and Whipple towards Ursa Minor. The MAGIC collaboration has given lower limits on the Lorentz invariance breaking scale based on the observation of a flare of Mkn501.

Three new instruments are coming very soon. The GLAST satellite was successfully launched on June 11 2008. The second MAGIC telescope will be installed by the end of the year. Finally, the large 28-meter telescope of HESS will be completed in 2009. New astroparticle results from VHE instruments are likely to come soon.

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Solar axion search with the CAST experiment


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The CAST (CERN Axion Solar Telescope) experiment is searching for solar axions by their conversion into photons inside the magnet pipe of a LHC prototype dipole. In the phase II, CAST is operating with a buffer gas inside the magnet bore apertures in order to extend the sensitivity of the experiment to larger axion masses. Preliminary results from the analysis of this second phase with $^4$He inside the magnet pipes excludes axions down to $g_{a\gamma} < 2.2 \cdot 10^{-10}$ GeV$^{-1}$ at 95% C.L. for $0.02 \text{ eV} < m_a < 0.39 \text{ eV}$. The data analysis has resulted in the most restrictive experimental limit on the coupling constant of axions to photons and for the first time experimental result has entered the theory motivated axion parameter space. At the beginning of 2008, data started to be taken with $^3$He in the magnet pipes in order to extend the sensitivity to axion masses up to 1.2 eV.

1 Introduction to Axion Theory

In the Standard Model, the QCD Lagrangian \(^1\) includes a gluon interaction term that violates charge conjugation times parity (CP) and time reversal (T):

$$\mathcal{L}_{CP} = \bar{\theta} \frac{\alpha_s}{8\pi} G \tilde{G}$$

where $G$ is the color field strength tensor, $\tilde{G}$ its dual, $\bar{\theta}$ represents the effective QCD vacuum and it can assume any value between 0 and $2\pi$. A non vanishing $\bar{\theta}$ value would imply CP violation.

Evidence of CP violation would be observable by the electric dipole moment of neutrons, that has been theoretically estimated to be \(^2\):

$$d_n \approx \bar{\theta} \frac{e}{m_n} \frac{m^*}{\Lambda_{QCD}} \propto \bar{\theta} \cdot 3.6 \cdot 10^{-16} \text{ e cm}$$

where $\Lambda_{QCD}$ is QCD energy scale $\approx 1$ GeV and $m^*$ is the reduced mass of the up and down quark defined as $m^* = \frac{m_u m_d}{m_u + m_d}$; $m_n$ is the neutron mass, $e$ the unit electrical charge. The present experimental upper limit \(^3\) however, is smaller than

$$d_n < 2.9 \cdot 10^{-26} \text{ e cm (90\% CL)}$$

Consequently, the phase parameter $\bar{\theta}$ should be smaller than $10^{-10}$. This implies that CP is not very strongly broken in the strong interactions. So the solution of the $U(1)_A$ problem begets a different problem: why is CP not badly broken in QCD? This is known as the strong CP problem.

One solution was proposed by Peccei Quinn in 1977 \(^4\) introducing a new global chiral $U(1)_{PQ}$ symmetry. This symmetry is necessarily spontaneously broken at an unknown scale $f_a$, and its introduction into the theory effectively replaces the static CP-violating angle $\bar{\theta}$ with a dynamical CP-conserving field - the axion. Formally, to make the Lagrangian of the Standard Model $U(1)_{PQ}$ invariant this Lagrangian must be increased by axion interaction:

$$\mathcal{L}_a = -\frac{1}{2} (\partial a)^2 + \frac{\alpha_s}{8\pi f_a} a G \tilde{G}$$

\(^1\) Ref. 1
\(^2\) Ref. 2
\(^3\) Ref. 3
\(^4\) Ref. 4
where $\alpha_s$ is the strong coupling constant, $a$ is the axion field, $f_a$ the axion decay constant and $f_a \propto \frac{1}{g_a \gamma}$. Color anomaly factors have been absorbed in the normalization of $f_a$ which is defined by this Lagrangian. Non-perturbative effects induce a potential for $a$ whose minimum is at $a = \bar{\theta} f_a$, thereby canceling the $\bar{\theta}$ term in the QCD Lagrangian, and thus allowing for the dynamical restoration of the CP symmetry.

Weinberg\textsuperscript{5} and Wilczek\textsuperscript{6} realized that a consequence of this mechanism is a new pseudo-scalar boson, the axion, which is the Nambu-Goldstone boson of the PQ symmetry. The axion coupling to ordinary matter is proportional to the axion mass $m_a$ and, equivalently, to the inverse of the PQ scale $1/f_a$. The PQ symmetry is explicitly broken at low energies by instantaneous effects so that the axion acquires a small mass:

$$m_a = \frac{z^{1/2} f_\pi m_\pi}{1 + z} \approx 6 \text{ eV} \frac{10^9 \text{GeV}}{f_a}$$

where $f_\pi$ is the pion decay constant and $z = m_u/m_d$ is the mass ratio of up and down quarks (for this numerical estimate\textsuperscript{7} we used a canonical value of $z=0.56$). Depending of their density and mass, axions may constitute a candidate for the cold dark matter in the universe.

One generic property of the axion is a two-photons interaction that plays a key role for most searches:

$$\mathcal{L}_{a\gamma} = \frac{1}{4} g_{a\gamma} F \tilde{F} a = -g_{a\gamma} E \cdot B a$$

where $F$ is the electromagnetic field-strength tensor, $\tilde{F}$ its dual, $E$ and $B$ the electric and magnetic fields. The coupling constant is:

$$g_{a\gamma} = \frac{\alpha}{2\pi f_a} \left( \frac{E}{N} - \frac{24 + z}{31 + z} \right)$$

where $\alpha$ is the fine structure constant, $E$ and $N$ are the electromagnetic and color anomaly of the axial current associated with the axion field and $E/N$ is a model dependent parameter\textsuperscript{8}.

### 2 Solar Axions

Axions could be produced in stellar cores through their coupling to plasma photons, namely by Primakoff conversion\textsuperscript{9}, with energies in the range of keV. The expected solar axion flux on earth based on a Standard Solar Model (see fig.1) is well approximated by\textsuperscript{10}:

$$\frac{d\Phi}{dE} = \left( \frac{g_{a\gamma \gamma}}{10^{-10} \text{GeV}^{-1}} \right)^2 \Phi_0 E^{2.481} e^{E/(1.205)}$$

where $E$ is in keV and $\Phi_0 = 6.020 \cdot 10^{10} \text{ cm}^{-2} \text{s}^{-1} \text{keV}^{-1}$. The maximum of the distribution is at 3.0 keV and the average energy is 4.2 keV.

This flux can be searched with the inverse process described by (6) where an axion converts into a photon in a macroscopic magnetic field\textsuperscript{12}. The conversion probability in magnetic field region in vacuum is given by\textsuperscript{13}

$$P_{a \rightarrow \gamma} \approx (BL g_{a\gamma \gamma})^2 \frac{\sin^2 (qL/2)}{(qL)^2}$$

where $B$ is the magnetic field strength, $L$ is the path length, $q$ is the momentum transfer between the axion and the X-ray photon $q = \frac{m_a^2}{2E_0}$, $E_0$ is the axion energy. The conversion is coherent over a large propagation distance for $qL \ll 1$. 

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Figure 1: Axion surface luminosity seen from the Earth as a function of axion energy and the radius where the axion are produced in the sun normalized to the solar radius.\textsuperscript{11}

One could reach sensitivity for larger axion masses by filling the magnetic field region with a gas, consequently the photon acquires an effective mass, and the conversion probability is\textsuperscript{14}

\[ P_{a\rightarrow\gamma} \propto \left( \frac{BLd_{\text{eff}}}{2} \right)^2 \frac{1}{q^2 + \Gamma^2/4} \left[ 1 + e^{-\Gamma L} - 2e^{-\Gamma L/2} \cos(qL) \right] \]  \hspace{1cm} (10)

where $\Gamma$ is the inverse photon absorption length for the X-rays in the medium, $q$ is the momentum transfer $q = \frac{m_e^2 - m_a^2}{2E_a}$, $E_a$ is the axion energy, $m_\gamma$ is the photon effective mass in the buffer gas that depends on the type and density of gas:

\[ m_\gamma \approx \sqrt{\frac{4\pi \alpha N_e}{m_e}} = \sqrt{\frac{Z}{A}} \rho \]  \hspace{1cm} (11)

where $Z$ and $A$ are the atomic and mass number of the gas, $m_e$ is the electron mass, $N_e$ is the number of electron per cm$^3$ and $\rho$ is the gas density. The coherence $qL < \pi$ can be restored for a narrow mass range:

\[ \sqrt{m_e^2 - \frac{2\pi E_a}{L}} < m_a < \sqrt{m_\gamma^2 - \frac{2\pi E_a}{L}}. \]  \hspace{1cm} (12)

The choice of a specific gas pressure allows the test of a specific axion mass. Scanning over a range of pressure means scanning over a large range of axion mass.

3 CAST

The CAST\textsuperscript{15-16} (CERN Axion Solar Telescope) experiment uses a decommissioned LHC prototype superconducting magnet with a length of 9.26 m, a magnetic field of about 9 T inside two beam pipes. The magnet is mounted on a moving platform allowing a vertical movement of $\pm 8^\circ$ and an horizontal movement of $\pm 40^\circ$. Different detectors are mounted at the two ends of the magnet cryostat. The magnet can point to the sun core for about 1.5 hour at the sunrise and at the sunset, when the magnet is not aligned with the sun the time is dedicated to background measurements. The detectors cover about $1/10^{th}$ of the solar radius, searching for the axion potentially emitted by the solar core. The CAST tracking system has been accurately calibrated by geometric survey measurements and it is verified by a sun filming. The pointing precision is evaluated to be better than $0.01^\circ$.\textsuperscript{14}
On one end of the magnet, a conventional plexiglas Time Projection Chamber (TPC)\(^{17,18}\), covering both beam pipes, looks for X-rays coming from axion-to-photon conversion during the sunset. On the opposite magnet end, a Micro-Mesh Gaseous detector (MM)\(^{19}\) and a Charge Coupled Device (CCD)\(^{20}\) look for X-rays during the sunrise. The CCD is located at the end of X-ray telescope designed for the German X-ray satellite mission ABRIXAS\(^{21}\). This telescope focuses the photons from the magnet bore with a 14.5 cm\(^2\) aperture to a spot size of about 6 mm\(^2\) on the CCD, and thus improving the signal to background ratio by a factor of about 150.

4 Result

During the phase I (2003-2004 data taking) the CAST experiment was operated keeping the magnet pipes evacuated and therefore it was sensitive to the axion mass range up to \(m_a = 0.02\text{eV}\) due to coherence effects. The results of phase I improved the limit of the axion coupling constant by a factor of 7 with respect to the previous experimental searches and it went beyond the astrophysical limit of globular clusters for coherence masses\(^{11}\).

The phase II of CAST, where the magnet pipes are filled with a buffer gas, can extend the sensitivity to higher axion rest masses. During 2005, the experiment was upgraded to allow the injection of a buffer gas in the cold bores. A new gas system was built and has been operation since the end of 2005. This system gets a density stability, an accuracy of pressure of about 0.2 mbar and a reproducibility precision of 0.01 mbar. The data taking of phase II with \(^4\)He during 2005 and 2006 covered one different pressure per day for a total of 160 density steps (step size is about 0.083 mbar at 1.8 K) up to 13 mbar, near the condensation limit of \(^4\)He gas at the

![Figure 2: 95% C.L. exclusion line in the axion/photon coupling constant versus the axion mass plane obtained from the complete CAST phase I data\(^{11}\) (line labeled “CAST Phase I”), from phase II with \(^4\)He (blue line labeled \(^4\)He”) and the foreseen sensitivity for CAST phase II with \(^3\)He (red line labeled \(^3\)He”). These results are compared with other laboratory limits such as previous helioscopes Lazarus et al.\(^{22}\), the Tokyo\(^{23}\) helioscope and those obtained from axion experiments with crystalline detectors located underground (SOLAX\(^{24}\), COSME\(^{25}\) and DAMA\(^{26}\)) and other constraints like the horizontal branch (HB) stars\(^{27}\). The grey band labeled “CMB limit” represents the limit evaluated by the amount of Hot Dark Matter deduced by the cosmic microwave background data\(^{28}\). The yellow band represents typical theoretical models with \(|E/N - 1.95|\) in the range 0.07 – 7 while the green solid line corresponds to the case when \(E/N = 0\) is assumed.
operation temperature of 1.8 K that is about 16 mbar. This corresponds to a scan of axion mass range between 0.02 $eV$ up to 0.39 $eV$. The preliminary combined results of $^4$He data for all three detectors are shown in fig. 2. The average upper limit to the axion-photon coupling at 95% C.L., in the axion mass range 0.02 $eV$ up to 0.39 $eV$, is $g_{a\gamma} < 2.2 \cdot 10^{-10} GeV^{-1}$. For the first time experimental results entered in the QCD theoretically allowed axion models region.

During 2007 the experiment was upgraded to allow the injection of $^3$He as buffer gas. This gas can reach pressures up to about 135 mbar at 1.8 K, corresponding to axion masses up to about 1.2 $eV$. The new automatic gas system has the same accuracy and reproducibility of the previous one and moreover it has a high safety and reliability level required to have no leaks of this gas, considerable expensive. The new gas system is operational since end of 2007. In parallel with the above mentioned upgrade, CAST improved the detector system replacing the TPC and MM by MicroMegas detectors based on the new bulk and microbulk technology. These detectors have a better energy resolution and a better background rejection. The new detectors are covered by a shielding composed of copper, lead, cadmium, nitrogen and polyethylene allowing background reduction by a factor 3. The data taking for phase II with $^3$He started in March 2008. CAST plans to run during the next 3 years to fully exploit the region up to axion mass of 1.2 $eV$, entering deeper into the QCD theoretically axion models space up to the region excluded by the amount of Hot Dark Matter induced by the cosmic microwave background data$^{28}$ (CMB). The expected reachable region is shown in fig. 2.

5 Conclusion

The CAST experiment has been running in phase I (vacuum in the beam pipe in 2003 and 2004) yielding to a lower limit on $g_{a\gamma} < 8.8 \cdot 10^{-11} GeV^{-1}$ at 95% C.L. for axion masses $m_a < 0.02 eV$. During 2005, a major modification to the magnet pipe system was undertaken to fill the beam pipes with $^4$He during 2005 and 2006. A new system was commissioned and installed in 2007 to use as buffer gas $^3$He allowing to exploit the region up to axion mass of 1.2 $eV$. The preliminary results of phase II with $^4$He has been presented, the average upper limit to the axion-photon coupling is $g_{a\gamma} < 2.2 \cdot 10^{-10} GeV^{-1}$ at 95% C.L. for axion masses 0.02 $eV < m_a < 0.39 eV$. These results improve the previous constraints given by other experiments by a factor of 7 and have entered the theory motivated axion parameter space. The $^3$He phase, with data taking started in March 2008, allows to enter deeper in the theoretical axion model allowed phase space and will continue for about three years.

References

27. A. Buzzoni et al., Â 128, 94 (1983).
I discuss different aspects of the phenomenology of hypothetical sub eV mass particles arising in the context of extensions of the standard model. I focus on a simple extension based on an additional U(1) gauge symmetry and its corresponding gauge boson, called “hidden photon”. Kinetic mixing with the standard photon leads to photon-hidden photon oscillations that are searched for in laboratory experiments like ALPS at DESY. Hidden photons produced in the interior of the Sun could be also detected in axion helioscopes like CAST at CERN and could play an interesting role in late cosmology, where the presence of additional feebly interacting relativistic particles seems to be favored. All these effects disappear as the hidden photon mass decreases, allowing phenomenologically large kinetic mixings. However, in this case such a hidden photon will even play a role in gauge coupling unification.

In the days of exploring the TeV frontier, are we leaving something behind us?

It is a common opinion, and we will find numerous examples of it in this volume, that the standard model (SM) of particle physics is not completely satisfactory to describe certain aspects of nature. Extensions of the SM invoked to cure their diseases include generally many additional symmetries and fields. The corresponding particles have generally masses arranged to lay beyond the reach of our collider experiments (or just around the corner), namely beyond a TeV. It is clear that if these additional particles are very massive we have little chances of discover them in colliders, and we should rely on low energy precision experiments. But they could be additional light particles. On general grounds, low masses are related to some symmetry that prevents high radiative contributions from larger mass scales. It is clear that the knowledge of these hypothetical low energy particles will provide us with an understanding of their related symmetries, and guide us through the difficult task of extending the standard model to describe particle physics up to arbitrarily high energies.

Of course, when these particles couple directly to the SM its existence is severely constrained...
from laboratory searches and our current understanding of astrophysics and cosmology. However, there are certain models in which the powerful astrophysical constraints are evaded \(^1\).

1 Massive Hidden photons and the “meV valley”

In this contribution I focus on one of these models, whose only addition to the SM lagrangian consists in a new \(U(1)\) gauge symmetry and its corresponding gauge boson, here called “hidden photon”. The SM fields are assumed to be uncharged under this new gauge group, but nevertheless they can still interact with the hidden photon through kinetic mixing with the standard model photon. Therefore we will consider the low energy effective lagrangian

\[
L = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} + \frac{\sin \chi}{2} B_{\mu\nu} F^{\mu\nu} + \frac{\cos^2 \chi}{2} m^2 \gamma_{\mu} B_{\mu} B^\nu,
\]

(1)

where \(F_{\mu\nu}\) and \(B_{\mu\nu}\) are the photon (\(A^\nu\)) and hidden photon (\(B^\nu\)) field strengths. The dimensionless mixing parameter \(\sin \chi\) can be generated at an arbitrarily high energy scale and does not suffer from any kind of mass suppression from the messenger particles communicating between the visible and the hidden sector. This makes it an extremely powerful probe of high scale physics. The construction outlined here arises quite naturally in extensions of the SM based on string theory, where values in the range \(10^{-16} \lesssim \chi \lesssim 10^{-2}\) can be expected \(^2\).

The most prominent implication of the kinetic mixing term is that photons are no longer massless propagation modes. The kinetic mixing term can be removed by changing the basis \(\{A, B\} \rightarrow \{A_R, S\}\), where \(A_R = \cos \chi A\) is a renormalized photon field and \(S = B - \sin \chi A\) is the state orthogonal to it, and therefore completely sterile with respect to electromagnetic interactions. The renormalization is typically unobservable and will be discussed in section 2.

In this section we use \(A = A_R\). In the \(\{A, S\}\) basis the kinetic term is diagonal but kinetic mixing has provided an off-diagonal mass term which produces \(A - S\) (vacuum) oscillations with a probability

\[
P_{A-S} = \sin^2 2\chi \sin^2 \frac{m^2 \gamma_{\mu} L}{4\omega}.
\]

(2)

where \(\omega\) is the energy and \(L\) is the oscillation length. It also modifies the static Coulomb potential with a Yukawa-like contribution

\[
V(r) = -\frac{\alpha}{r} \left( \cos^2 \chi + e^{-m_{\gamma'} r} \sin^2 \chi \right).
\]

(3)

The phenomenology of such a model has been considered by Okun \(^3\) and others \(^4\). The stronger laboratory constraint comes from precision measurements of the Coulomb law \(^5\) and can be read off in Fig. 1. The sensitivity of this test has a clear maximum at distances of the order of the centimeter, corresponding to \(m_{\gamma'} \sim \mu\text{eV}\). For much smaller \(m_{\gamma'}\) the hidden photon contribution is indistinguishable from that of a single massless photon, and for much higher \(m_{\gamma'}\) the hidden photon contribution is exponentially suppressed.

Let us now consider the astrophysical bounds, focusing on the case of the Sun. Photons of the interior of the Sun can oscillate into the sterile component that can freely escape from the solar interior removing energy that otherwise will take much longer to drain. The response of the solar structure to such an exotic energy loss is to raise the temperature of the interior so that the thermonuclear reactions can provide this extra energy. The consequence is that Hydrogen is converted much faster into Helium and the duration of the Hydrogen burning period is reduced. Studies of Raffelt and Dearborn \(^6\) concluded that such an exotic luminosity cannot be higher than the actual visible luminosity of the Sun. Therefore, integrating eq. 2 over the thermal distribution of photons over the solar interior will provide us with a limit \(^4,\!^7\) on \(\chi\). To proceed we only have to take into account an important subtlety: namely that photons in a plasma
propagate as massive particles with a mass given by the plasma frequency $\omega_P = 4\pi\alpha n_e/m_e$ with $\alpha$ the fine structure constant and $m_e, n_e$ the electron mass and number density. In such a case, the essential modification of the transition probability is the introduction of an effective mixing angle given, in the small $\chi$ approximation, by

$$\chi^2 \rightarrow \chi_{\text{eff}}^2 = \frac{\chi^2 m_e^4}{(\omega_P^2 - m_e^2)^2 + (\omega \Gamma)^2}$$

which strongly suppresses $A - S$ transitions, and therefore energy drain, when $m_{\gamma'} \ll \omega_P$. Here $\Gamma$ is the photon absorption rate and cuts-off the effective mixing in the resonant regime where $m_{\gamma'} \approx \omega_P$ and the amplitude of the oscillations is maximum. The plasma frequency in the solar interior varies in the range $1 \text{ eV} \leq \omega_P \leq 300 \text{ eV}$ (and typically $\omega \Gamma$ is smaller), so hidden photons with masses below the eV can evade the solar luminosity bound even with relatively high values of the vacuum mixing parameter $\chi$ (See Fig. 1).

Even a low, harmless, hidden photon flux can be detectable at earth by a suitable detector. The CAST collaboration at CERN\(^8\) (See the review of Silvia Borgi in this same Proceedings) operates a search for solar axions of keV energies by tracking the Sun with a 10 m long LHC magnet, since axions emitted from the Sun can convert into photons by the inverse Primakoff effect.\(^9\) Such an experiment will be also sensitive to hidden photons, with the benefit that high vacuum conditions are kept in the conversion region and thus the effective mixing angle is not suppressed. This nearly background free experiment can measure a photon spectral flux generated inside the magnet of $10^{-5}$ photons per second, cm$^2$ and keV. This number was used\(^7\) to set the hidden photon limit labeled CAST in Fig. 1. A recent paper\(^10\) has pointed out that considerable improvement can be achieved by measuring hidden photons of lower energies $\sim$ eV where the flux is maximal since it mostly comes from the external shells of the Sun where the electron density (and hence the plasma frequency) is smallest.

The Coulomb and CAST limits leave a valley in the allowed parameter space around the suggestive mass scale of $m_{\gamma'} \approx \text{meV}$. Since photon-hidden photon oscillations are resonant when a plasma is present such that $\omega_P = m_{\gamma'}$, it would be advantageous to find environments with a huge number of photons and electron densities $\sim 10^{15} \text{ cm}^{-3}$. These conditions are found in the early universe when the temperature is of order $\sim \text{keV}$, i.e. after big bang nucleosynthesis (BBN) but before the cosmic microwave background (CMB) formation. In such a scenario a fraction of the photon background will be resonantly converted into hidden photons, forming a hidden cosmic microwave background (hCMB)\(^11\). This hCMB decouples much before than the standard CMB and from that moment on mimics the effect of additional neutrino species, $N_{\nu}^{\text{eff}}$. Since some of the CMB photons disappear, the baryon to photon ratio $\eta$ measured at decoupling also increases with respect to the value suggested by BBN. Therefore we can bound $\sim \text{meV}$ hidden photons from the agreement of the values of $N_{\nu}^{\text{eff}}$ and $\eta$ provided from BBN and CMB physics\(^11\). The CMB observations, combined with large scale structure data (LSS), slightly prefer\(^12\) $N_{\nu}^{\text{eff}} > 3$ but both frameworks can be made to coincide within the quoted errors. The preference of a high $N_{\nu}^{\text{eff}} > 3$ is supported by the SDSS and Ly-$\alpha$ data and might be likely due to systematics\(^13\). However, even in\(^13\) where a more careful treatment of the bias parameters is included, values slightly higher than 3 are still preferred, with a best global fit of $N_{\nu}^{\text{eff}} = 3.8^{+2.0}_{-1.6} \ (95\% \ C.L.)$. It is however premature to consider that such an excess has a physical interpretation in terms of new physics, but if eventually it is confirmed it may require new weakly interacting particles that are relativistic at CMB, namely sub eV particles, and hidden photons could certainly do the job.

Note that the conservative “suggested” excess $\Delta N_{\nu}^{\text{eff}} \simeq 0.8$ corresponds to a hidden photon with $\chi \simeq 2 \times 10^{-6}$. The mass should be then $\simeq 0.2 \text{ meV}$ to avoid distortions of the CMB Plack spectrum and the laboratory searches to be presented next. At the view of Fig. 1 this leads us to a clear goal in the parameter space!
Interestingly, such a scenario is going to be tested in the near future in the laboratory. The ALPS (Any Light Particle Search) experiment at DESY\textsuperscript{14} is currently setting up an upgraded “light-shinning-through-walls” experiment\textsuperscript{15} that will explore much of the relevant parameter space in the “meV valley”. The set up consists in a powerful laser beam which propagates under high vacuum conditions to end up blocked in an opaque wall. If hidden photons or any other weakly interacting low mass particles are produced before the wall they will go through it and can be reconverted after the wall in another high vacuum cavity in which a sensitive detector is placed. Some similar experiments have been already performed\textsuperscript{16,17,18}, the most recent motivated by the recent PVLAS episode\textsuperscript{19} and a recent paper has interpreted them in terms of hidden photons\textsuperscript{20}. The current ALPS proposal includes 300 W of laser power, conversion and reconversion lengths of $\sim$ 6 meters and a small background $\sim$ 50 mHz. The results will be presented in late fall of this same year, and immediately after several upgrades will be performed, including possibly higher laser power, a new detector and “phase shift plates”\textsuperscript{21} to enhance the coherence between photons and hidden photons.

Already with the first upgrade the ALPS experiment will be sensitive to part of the region of major cosmological interest, and will eventually cover it completely with subsequent improvements. On the long term, additional coverage could be also provided by the mentioned new solar hidden photon searches\textsuperscript{10} or by a photon regeneration experiment using radio waves instead of laser light\textsuperscript{22}.

Figure 1: The “meV Valley” in the mass-mixing plane of a hidden photon is bounded at low masses by searches from deviations of the Coulomb law and from searches of solar hidden photons with the CAST helioscope at higher masses. Light-shinning-through-walls (LSW) experiments have explored the peaceful realm around $m_{\gamma} \sim \text{meV}$ and an upgraded ALPS setup will penetrate even deeper in the near future. In the early universe a part of the CMB can resonantly oscillate into hidden photons contributing, as neutrinos do, to the radiation density at decoupling. Values higher than $N_{\nu} > 5$ can be excluded, but a value slightly higher than 3, $N_{\nu} \sim 3.8$ is still preferred (Red line). The precise determination of the CMB spectrum by FIRAS constrains the distortions that the creation of this hidden CMB would imprint on it. An experiment exploiting microwave cavities could be sensitive to most of the region of cosmological interest. See the text for references.
2 Massless hidden photons and Unification

All effects mentioned before are lost in the probably most natural case, a massless hidden photon. Well, not all of them. We have already mentioned that to get rid of the kinetic mixing and define fields with canonical kinetic terms we need to renormalize the photon field with a factor \( \cos \chi \). In the low energy lagrangian considered this is harmless, since a photon renormalization is simply reabsorbed in the definition of the electric charge as usual. Even when one considers the whole standard model gauge group and allows our new U(1) gauge boson to mix with the boson of hypercharge (kinetic mixing with non-abelian gauge fields will not respect gauge invariance) the corresponding gauge coupling \( g_1 \) will absorb again this factor and leave no trace in precision electroweak observables.

Since this shift will only affect \( g_1 \) but not \( g_2 \) or \( g_3 \), it could be detectable in a theory in which there is an a priori relation between the couplings, such as in grand unification. In this case we shall define the unification scale by the equality of the two couplings that are unchanged \( g_2(m_{\text{GUT}}) = g_3(m_{\text{GUT}}) \). Note that we measure the “renormalized” \( g_1 \) and this is always larger than the real value (the one we would expect to unify) in a factor \( 1/\cos \chi \), namely

\[
g_1^\text{measured} = \frac{g_1^\text{real}}{\cos \chi}.
\]

Interestingly, the measured value of \( g_1 \) in the standard model turns out to be also larger than the required to unify with \( g_2 \) and \( g_3 \) in a pure SU(5) model without supersymmetry. Therefore unification could be achieved at a scale \( \approx 10^{17} \) GeV (evading limits from proton decay) but being “masked” by the exotic hypercharge renormalization due to kinetic mixing with \( \chi \approx 0.4 \).

We have taken values of \( g_{1,2,3} \) at the Z-pole from and plotted the running in Fig. 2.

The case with supersymmetry (SUSY) is more complicated. Using the renormalization group equations at the one-loop level, a small value of \( \chi \approx 0.055 \) improves the already impressive unification, but this effect is of similar magnitude than the threshold corrections of SUSY particles of \( \sim \) TeV masses and particles at the scale of unification. When these corrections are included the measured value of \( g_1 \) seems to be a bit smaller than the required to unify perfectly (see for a recent discussion). While a more detailed study is under way, the two possible outcomes are clear: if \( g_1(m_{\text{GUT}}) < g_{2,3}(m_{\text{GUT}}) \) a bound on \( \chi \) of order \( 10^{-2} \) can be set, in the opposite case, a small value of \( \chi \) could be the responsible of the difference and unification could be achieved.

![Figure 2: One-loop running of the SM gauge couplings with an exotic renormalization of the hypercharge coupling \( g_1 \) due to kinetic mixing with an additional massless U(1) gauge boson. LEFT: standard model, RIGHT: with supersymmetry. Note that \( \alpha_{1,2,3} = g_{1,2,3}^2/(4\pi) \) and \( g_1 \) has been normalized with the usual SU(5) factor \( \sqrt{5}/3 \).](image-url)
References


DM-TPC: A DARK MATTER TIME PROJECTION CHAMBER

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The DM-TPC collaboration is developing an optical read-out time projection chamber for dark matter detection and identification. The detection material is low pressure CF4 gas, which has exceptional characteristics making it highly suitable for this purpose. Among these are excellent electron diffusion and scintillation performance, and the great sensitivity of fluorine to spin-dependent dark matter interactions. A summary of progress of the research and development program will be presented here.

1 An Imaging Time Projection Chamber

The concept of the Dark Matter Time Projection Chamber (DM-TPC) is illustrated in figure 1. The chamber contains CF4 gas at low pressure. An incoming dark matter particle collides elastically with an F nucleus, causing it to recoil with 10 to 100 keV of energy, with a range the order of one mm or more. The ionization electrons drift to the readout plane, where an electron avalanche occurs, accompanied by the emission of copious amounts of scintillation light. The scintillation light is imaged by the lens and CCD camera. Information on the component of the track perpendicular to the plane is provided by the time profile of a PMT signal recorded with a waveform digitizer. CF4 is an ideal gas due to its large scintillation efficiency and excellent electron diffusion characteristics. Furthermore, F has excellent properties for use in searches for dark matter through spin dependent interactions.

We have built two prototype detectors that have demonstrated this technique. An MIT prototype has used a readout plane consisting of a Multi-Wire-Proportional-Chamber (MWPC) to

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show that the direction of recoiling F nuclei can be determined. Examples of recoil tracks are shown in figure 2 for a CF4 pressure of 200 torr$^1$. The recoils were produced by elastic collisions with neutrons, and extend down to energies of about 200 keV. Since the stopping power decreases as a particle’s velocity approaches zero, the stopping end of the track is the one with reduced light yield. We expect that this effect will be observable below 100 keV when operated at lower pressure. Furthermore, we have recently developed a mesh readout plane that will also help with this goal. Tracks using the mesh in the MIT prototype are shown in figure 3. The stainless steel mesh, with wire diameter 30 $\mu$m and wire pitch 320 $\mu$m, was supported over a copper plate with spacers made of 300 $\mu$m plastic foil. With a CF4 pressure of 200 torr, the mesh operated with 2000 volts between the mesh and copper plate without sparking. The tracks shown in Figure 3 were produced by alpha particles. The Bragg peak from one alpha track is shown in figure 4. A second prototype has recently built at Boston University. This device has a 25 cm by 25 cm mesh and a cylindrical field cage 25 cm in diameter and 25 cm high.

![Schematic of DM-TPC](image)

**Figure 1: Schematic of DM-TPC.**

2 Dark Matter Search via Spin-Dependent Interactions: Majorana Neutrinos

The latest cosmic microwave background results show that the matter content of the universe is 5.7 times larger than exists in the form of ordinary baryonic material$^2$. Through observations of visible light, x-rays, and weak gravitational lensing of the collision of two clusters of galaxies, the dark matter has been found to interact very weakly, if at all, with itself and with ordinary matter by any means other than gravitation$^3$. Many hypothetical particles have been advanced as candidates for dark matter. The list includes axions, the neutralino from supersymmetry, technibaryons, and massive fourth generation neutrinos from many models. Data on neutrino oscillations and direct limits on the mass of the electron neutrino from beta decay experiments have shown that the neutrinos from the first three generations are too light to provide the cold dark matter required for galaxy formation. LEP excludes fourth generation neutrinos with mass less than half the Z mass, and the first direct search for dark matter excluded massive Dirac neutrinos as the dark matter of our Galaxy for neutrino masses between 20 GeV and 2 TeV$^4$. 
Massive Majorana neutrinos have not yet been excluded as the dark matter. We consider the search for these particles to illustrate the potential of the DM-TPC.

Figure 2: Tracks of recoiling F nuclei from collisions with neutrons. The neutrons came from the right. The direction of the neutron can be inferred from the track.

Figure 3: Alpha tracks from MIT prototype using mesh readout plane.

Majorana neutrinos are predicted by numerous theories. They have been shown to be viable as dark matter in a minimal technicolor theory$^5$. The authors of reference 5 have shown that the relic density of Majorana neutrinos can be in the observationally determined range if the Universe has a nonstandard history including an early dominance by a rolling scalar field as predicted by many models for dynamical dark energy. Majorana neutrinos also show up as Kaluza-Klein particles in a theory involving the universal extra dimension scenario$^6$. Their collision cross section with nuclei is determined by their mass and the properties of the nuclei they interact with. The total cross section is $\sigma = \frac{\alpha^2 \mu^2}{s \lambda^4} C^2 \lambda^2 J(J + 1)$, where $\mu$ is the reduced mass of the neutrino, $J$ is the nuclear angular momentum quantum number of the nucleus, and $\lambda^2 J(J + 1)$
is the "spin-factor" of the nucleus. C is related to the quark spin content of the unpaired nucleon and is 0.68 for the proton and -0.58 for the neutron. Estimates of the nuclear spin factor are given in reference 7 for a number of isotopes. The product of squared nucleus mass, isotopes per kg, and nuclear spin factor is a figure of merit as a Majorana neutrino detecting medium. This figure of merit is listed in parentheses (arbitrary units) for various promising nuclides: 1H (4.48), 19F (74.65), 23Na (5.68), 27Al (14.16), 43Ca (0.06), 73Ge (2.24), 93Nb (90.85), 127I (5.36), 129Xe (25.00), 131Xe (9.18). Form factor effects are difficult to calculate for spin-dependent interactions, but for F the effect is less than 20% for recoil energies less than 100 keV. Thus it is seen that F provides the best return per kg of detector mass as any other detecting material. Leaving out the form factor correction term and assuming the neutrino mass is large compared to the F mass, the cross section for a Majorana neutrino with a fluorine nucleus is 3.1 pb. For a local dark matter density of 0.4 GeV/cm³ there would be 300 interactions per year for 100 GeV neutrino mass in 1 kg of F. A few dozen of these would likely have recoil kinetic energies in excess of 100 keV, when one includes the effects of the tails of the neutrino velocity distribution. With our direction sensitive detector we can use the measured directions of the high energy recoils to test the hypothesis that they are produced in collisions with particles coming from the direction of the constellation Cygnus, which is believed to be the case for dark matter particles.

In the next year we plan to build a detector with a mass of 70 g of fluorine in a volume of 0.1 m³ of CF₄ at a gas pressure of 150 torr. The device will be operated at an appropriate underground site to reduce backgrounds from neutrons. The detector will have in a single vacuum vessel four cells, each of the order of 30 cm on a side based on diffusion and optical considerations. Each cell will be viewed by one CCD camera and two PMTs. This will permit 3-D track reconstruction and redundancy will enable the elimination of background events due to radioactivity events in the CCDs. The device would detect about one Majorana neutrino interaction per month if these particles are the dark matter.

References

Cosmic Rays
THE AIR-FLUORESCENCE YIELD

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Detection of the air-fluorescence radiation induced by the charged particles of extensive air showers is a well-established technique for the study of ultra-high energy cosmic rays. Fluorescence telescopes provide a nearly calorimetric measure of the primary energy. Presently the main source of systematic uncertainties comes from our limited accuracy in the fluorescence yield, that is, the number of fluorescence photons emitted per unit of energy deposited in the atmosphere by the shower particles. In this paper the current status of our knowledge on the fluorescence yield both experimental an theoretical will be discussed.

1 Introduction

Fluorescence telescopes have been successfully used for the detection of ultra-high energy cosmic rays ($> 10^{18}$ eV) since the pioneering Fly’s Eye experiment. In this technique the fluorescence radiation induced by the charged particles of the extensive air shower generated by a primary cosmic ray is registered at ground by wide-angle telescopes. Assuming that the intensity of the air-fluorescence light is proportional to the energy deposited in the atmosphere by the shower, this technique provides a nearly calorimetric measure of the energy of the primary cosmic ray. Therefore it has the advantage, as compared with methods relying on simulations (e.g. surface arrays working in standalone mode), of being nearly model independent. In spite of this advantage, fluorescence telescopes are presently limited by the uncertainty in the fluorescence yield, that is, the calibration parameter which converts number of fluorescence photons into absolute energy units. For instance in the Pierre Auger Observatory the uncertainty in the fluorescence yield contributes a 14% to the total systematic error in the energy calibration which is presently 22%.

In order to improve the accuracy of this parameter, dedicated laboratory experiments are carrying out precise measurements of the air-fluorescence emission. In these experiments an
2 The generation of air fluorescence excited by electrons

2.1 Physical processes

Air-fluorescence in the near UV range (300 - 400 nm) is basically produced by the de-excitation of atmospheric nitrogen molecules excited by the shower electrons. Most of the fluorescence light comes from the 2P System of N₂ and the 1N System of N₂⁺ (Fig. 1a). Excited molecules can also decay by collisions with other molecules (collisional quenching). This effect which grows with pressure $P$, reduces the fluorescence intensity by a factor $1 + P/P_\lambda$. The characteristic pressure $P_\lambda$ is defined, for a given $v - v'$ band of wavelength $\lambda$, as the one for which collisional quenching and radiative decay have the same probability.

Basically two different parameters are being used for the energy calibration of fluorescence telescopes. The first one $\varepsilon_\lambda$ is the number of photons of a given molecular band emitted per electron and unit path length, $\varepsilon_\lambda = N \times \sigma_\lambda/(1 + P/P_\lambda)$, where $N$ is the density of nitrogen molecules and $\sigma_\lambda$ is the cross section for the excitation of the molecular band. The second parameter is the fluorescence yield $Y_\lambda$, defined as the number of photons emitted per unit deposited energy.

$$Y_\lambda = Y_\lambda^0 \frac{1}{1 + P/P_\lambda}, \quad Y_\lambda^0 = \frac{A_\lambda}{(dE/dX)_{dep}}.$$  

$Y_\lambda^0$ is the fluorescence yield in the absence of quenching. $A_\lambda$ and $(dE/dX)_{dep}$ are respectively the number of emitted photons at zero pressure and the deposited energy both per unit mass thickness. The fluorescence yield as defined in (1) is more useful for calorimetric applications. Notice that for the determination of $Y_\lambda$, both photon number and deposited energy has to be measured in the same volume. This is particularly important for laboratory experiments carried out in small gas chambers. In this case secondary electrons ejected in ionization processes might escape the field of view of the optical system before depositing all the energy (Fig. 1b). In next section the role of secondary electrons in the generation of air-fluorescence light is described.
2.2 Secondary electrons

Secondary electrons from ionization processes are the main source of fluorescence light, since the excitation cross sections show a fast decrease with energy (Fig. 1a), in particular the one for the 2P system. A high energy electron loses energy as a result of collisions with air molecules. Ionization processes give rise to the ejection of secondary electrons which deposit their energy within a certain distance from the interaction point (Fig. 1b). The average energy deposited per unit mass thickness inside a given volume around the interaction point can be expressed as

$$\rho \frac{dE_{\text{dep}}}{dX} = N_{\text{air}} \left\{ < E^{0}_{\text{dep}} > + < E_{\text{dep}} > \right\} \sigma_{\text{ion}}(E), \quad < E^{0}_{\text{dep}} > = < E_{\text{exc}} > \frac{\sigma_{\text{exc}}}{\sigma_{\text{ion}}} + I + < E^{\text{ion}}_{\text{exc}} >,$$

where \(\rho\) is the air density, \(N_{\text{air}}\) is the number of air molecules per unit volume and \(\sigma_{\text{ion}}\) is the ionization cross section. The average energy deposited in the medium by the primary electron per primary ionization process \(< E^{0}_{\text{dep}} >\) is obtained from several molecular parameters\(^a\). The energy deposited in the volume by the secondary electrons \(< E_{\text{dep}} >\) is calculated by a dedicated simulation\(^4\). Figure 2a) shows the result for a sphere of radius \(R\) (Fig. 1b). As expected, the deposited energy depends on \(PR\) and for an unlimited medium, \(PR \to \infty\), equals the energy loss predicted by the Bethe-Bloch theory.

Neglecting the collisional quenching, the number of photons emitted per electron and per unit path length can be expressed by \(\varepsilon_{\lambda}(P) = \rho A_{\lambda} = N\{\sigma_{\lambda}(E) + \alpha_{\lambda}(E,P)\sigma_{\text{ion}}(E)\}\), where \(\alpha_{\lambda}(E,P)\) is the average number of photons generated inside the volume per secondary electron, also calculated in the simulation. A very simple expression for \(Y^{0}_{\lambda}\) can be obtained from the above equations

$$Y^{0}_{\lambda} = \frac{N}{N_{\text{air}}} \times \frac{\sigma_{\lambda}}{\sigma_{\text{rest}}} + \frac{\alpha_{\lambda}}{< E^{0}_{\text{dep}} > + < E_{\text{dep}} >},$$

This procedure allows theoretical predictions on the absolute value of \(Y^{0}_{\lambda}\) and its dependence on the electron energy as shown below.

2.3 Fluorescence emission versus deposited energy

The energy calibration of fluorescence telescopes relies on the assumption that the intensity of fluorescence light is proportional to the energy deposited in the atmosphere, that is, the

\(^a\)ionization potential \(I\), total excitation cross section \(\sigma_{\text{exc}}\), average excitation energy of neutral molecules \(< E_{\text{exc}} >\) and of ionized molecules \(< E^{\text{ion}}_{\text{exc}} >\).
fluorescence yield is assumed to be independent on the electrons energy. The validity of this assumption can be theoretically checked by means of the model described above. Fig. 2b) shows $Y^0$ versus $E$ for the most intense band of the 2P system (0-0 transition at 337 nm). The results shown in this plot can be summarized as follows. The fluorescence yield decreases with $E$ about a 10% in the range 1 keV - 1 MeV and about 4% in the interval 1 MeV - 20 GeV. This smooth dependence of the fluorescence yield on $E$ has no impact on the energy calibration of fluorescence telescopes. The proportionality assumption has been also verified experimentally by several groups.²

3 The dependence of the fluorescence yield on atmospheric parameters

Fluorescence yield depends on pressure, temperature $T$ and humidity. Thus for a precise energy calibration of fluorescence telescopes these dependencies have to be determined accurately.

As mentioned above collisional quenching reduces the fluorescence emission by a factor $1 + P/P_x$. In the general case, for a mixture of gases (e.g. nitrogen, oxygen, water vapor, etc.), the characteristic pressure obeys the law

$$\frac{1}{P_x} = \sum_i \frac{f_i}{P'_i} = \frac{kT}{\sigma_{N_i} \bar{v}_{N_i}}, \quad \bar{v}_{N_i} = \sqrt{\frac{8kT}{\pi \mu_{N_i}}},$$

where $f_i$ is the fraction of molecules of type $i$ in the mixture, $\sigma_{N_i}$ is the collisional cross section which depends on the particular band, and $\bar{v}_{N_i}$ and $\mu_{N_i}$ are the relative velocity and reduced mass of the two body system N-i respectively.

The experimental procedure for the determination of the dependence of fluorescence yield on the above parameters is the following. At a fixed temperature the dependence of fluorescence intensity on pressure is measured for dry air. This measure, if properly carried out, allows a determination of $P'$ and therefore the dependence of the fluorescence yield on pressure at a fixed temperature. Experimental values of $P'$ for the molecular bands of the 2P and 1N systems in dry air at room temperature have been reported by many authors. The most complete set of $P'$ values have been reported very recently by AIRFLY improving the accuracy of previous measurements. This set of values are being used by the Pierre Auger Observatory for the calculation of the dependence of the fluorescence yield versus altitude.

The $P'$ parameter depends on temperature because the collision frequency grows with $\sqrt{T}$ as predicted by the kinetic theory of gases. In addition the collisional cross section depends on the kinetic energy of the encounters following a power law ($\sim T^\alpha$). Assuming this effect is negligible, the temperature dependence of the fluorescence yield can be easily predicted by equation (4). Recently some experimental works have found a noticeable variation of the collisional cross section with temperature. According to the preliminary values reported by AIRFLY, neglecting this effect results in an overestimation of the fluorescence yield by an amount going up to $\approx 20\%$ for the 1N (0-0) 391 nm band.

Water molecules have a significant cross section for the air-fluorescence quenching and therefore humidity modifies the value of $P'$. Several authors have measured the dependence of fluorescence intensity on humidity. A decrease of the fluorescence yield up to a 20% is found (at 100% relative humidity). From these measurements, values of the characteristic pressure for the quenching with water molecules $P'_{H_2O}$ have been determined for the main molecular bands of nitrogen.

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4 the effect of secondary electrons escaping the field of view might introduce systematic errors.

5 the contribution of the pressure dependence to the total uncertainty in the energy determination has been reduced to a 1%.
Figure 3: Comparison of fluorescence signal generated by a) an electron beam with b) that from Rayleigh scattering of a nitrogen laser. This procedure allows the absolute calibration of the optical system.

4 Absolute value

The accurate measurement of the absolute value of the fluorescence yield is an experimental challenge. The value is obtained as the ratio \( Y_\lambda = N_\lambda / (N_e \times EDEP) \). For the measurement of the absolute number of fluorescence photons in the wavelength interval of interest \( N_\lambda \), the efficiency of the various elements of the collection and detection system has to be known accurately. The number of electrons traversing the observation region \( N_e \) has to be absolutely measured as well. Finally the total energy \( EDEP \) deposited in the volume from where the registered fluorescence was emitted has to be determined (usually by means of a Monte Carlo simulation). In order to reduce systematic errors in the optical calibration (e.g. PMT quantum efficiency, transmission of optical elements, geometrical factors, etc.) some techniques have been developed, based on the comparison with well known physical processes like Cherenkov emission or Rayleigh scattering (Fig. 3).

Several measurements of \( Y_\lambda \) are presently available\(^3\). Unfortunately the comparison is not simple since some authors report the experimental result of \( \varepsilon_\lambda \) (i.e. photons/m) while others provide \( Y_\lambda \) (i.e. photons/MeV). In addition the spectral intervals of the various experiments use to be different. A detailed summary of the available results can be found elsewhere\(^5\). Here we will compare some representative experimental data (Tab. 1). For this comparison, measured values of \( \varepsilon_\lambda \) are converted into fluorescence yields using our results on deposited energy. Notice that deposited energy is weakly dependent on the size of the region and therefore a rough estimate of the equivalent \( R \) value is sufficient. From these results the fluorescence yield \( Y_{337} \) for the most intense band, 2P (0-0) at 337 nm, is calculated using the experimental relative intensities reported by AIRFLY\(^6\). Finally the \( Y_{337} \) values have been normalized to 293 K temperature and 1013 hPa pressure using equations (1) and (4). This procedure is appropriate for a comparison of measurements with typical uncertainties of about 13% or higher. Results are shown in last column of Tab. 1.

Firstly, the \( \varepsilon_\lambda \) values of Kakimoto et al. in the range 300-400 nm at several energies have been superimposed in Fig. 2a) to the energy deposited at atmospheric pressure assuming an observation volume with \( R \) ranging between 5 and 15 cm. The comparison of fluorescence intensity (photons/m) with deposited energy has allowed the determination of the fluorescence yield (photons/MeV) in that wavelength interval.

The \( \varepsilon_{337} \) value of 1.021 photons/m from Nagano et al. has been combined with the deposited energy for \( R \approx 5 \) cm giving the corresponding \( Y_{337} \) value. For the determination of the fluorescence yield, both MACFLY and FLASH calculate the deposited energy from a MC simulation. For these experiments only the conversion for wavelength intervals as well as minor \( T \) and \( P \) corrections were necessary. Finally AIRFLY reports a preliminary value of \( Y_{337} \) determined from the ratio of the absolute number of photons and the energy deposited according to a GEANT4 simulation.
Table 1: Comparison of data on fluorescence yields. Experimental results are used to infer the value of the fluorescence yield for the 337 nm band at \( T = 293 \) K and \( P = 1013 \) hPa (last column). See text for details.

<table>
<thead>
<tr>
<th>Experiment</th>
<th>( \Delta \lambda )</th>
<th>( T )</th>
<th>( P )</th>
<th>Experimental result</th>
<th>( I_{337}/I_{\Delta \lambda} )</th>
<th>( Y_{337} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>Kakimoto et al.</td>
<td>300 - 400</td>
<td>288</td>
<td>1013</td>
<td>see text</td>
<td>0.278</td>
<td>5.4</td>
</tr>
<tr>
<td>Nagano et al.</td>
<td>337</td>
<td>293</td>
<td>1013</td>
<td>1.021 ph./m</td>
<td>1</td>
<td>5.5</td>
</tr>
<tr>
<td>MACFLY</td>
<td>290 - 440</td>
<td>296</td>
<td>1013</td>
<td>17.6 ph./MeV</td>
<td>0.261</td>
<td>4.6</td>
</tr>
<tr>
<td>FLASH 07</td>
<td>300 - 420</td>
<td>304</td>
<td>1013</td>
<td>20.8 ph./MeV</td>
<td>0.276</td>
<td>5.6</td>
</tr>
<tr>
<td>AIRFLY (prelim.)</td>
<td>337</td>
<td>291</td>
<td>993</td>
<td>4.12 ph./MeV</td>
<td>1</td>
<td>4.0</td>
</tr>
</tbody>
</table>

5 Conclusions

Our understanding on the processes leading to generation of air fluorescence has increased significantly in the last years\(^5\). The world-wide campaign for the experimental determination of the fluorescence yield has achieved remarkable results, in particular in the measurement of the various dependencies with atmospheric parameters. The fundamental assumption of proportionality between fluorescence intensity and deposited energy has been verified both theoretically and experimentally.

In regard with the determination of the absolute value of the fluorescence yield new data are available. However the interpretation of the results is not straightforward. A comparison using the procedure discussed here shows a general agreement with typical differences of about 15%. For a real improvement in the accuracy of fluorescence telescopes an uncertainty better than 10% in the fluorescence yield is necessary. Several experiments claim high accuracy, for instance, the reported uncertainty of the FLASH experiment is of about 8%. In addition the AIRFLY collaboration will publish soon a final absolute value with an error below 10%. A discussion on these and other high accuracy measurements have been presented elsewhere\(^5\). Discrepancies between these experiments go beyond the reported accuracies and therefore some experimental effort is still necessary to clarify the situation.

Acknowledgments

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THE FLUX AND THE COMPOSITION OF ULTRA HIGH ENERGY COSMIC RAYS MEASURED BY THE PIERRE AUGER OBSERVATORY

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The cosmic ray flux above $10^{18}$ eV has been measured with high statistics by the Pierre Auger Observatory. At high energies the flux is suppressed and the hypothesis of a single power-law behavior as obtained in the lower energy range is rejected with a significance of more than 6 sigma. The measurement of the shape of the energy spectrum of ultra high energy cosmic rays can constrain acceleration models only when combined with a composition determination. The fluorescence detector measurement of the longitudinal development of air showers is used to determine the cosmic ray composition and the surface detector data are used to derive upper limits on the flux of photons and tau neutrinos.

Introduction

In the near future the wealth of data recorded by the Pierre Auger Observatory will help to answer some of the main questions in ultra high energy cosmic ray (UHECR) physics, such as their origin and composition. Due to the Greisen-Zatsepin-Kuzmin effect (GZK) a flux suppression is expected in the highest energy range, mainly caused by the energy loss of cosmic rays interacting with the microwave background radiation. This has been seen both by the Pierre Auger Observatory and by the HiRes collaboration. Through the hybrid technique of observing the same air showers with two different detectors, the nearly calorimetric estimation of the energy of the primary particle as obtained from the fluorescence technique is transferred to the large number of events recorded by the surface detector. This method is described below in section 1 and the resulting energy spectrum in section 2.

The excess of cosmic rays above the energy threshold given by the GZK effect, as reported by the AGASA experiment might be explained in the so-called top down models. In some of these scenarios the origin of cosmic rays is at relatively close distances to the Earth and therefore
a large photon content in the cosmic ray flux is predicted at the highest energies. Limits on photon content are presented in section 3 below.

1 Energy calibration

The Pierre Auger Observatory, located in the province of Mendoza (Argentina), is used to measure the properties of extensive air showers by observing their longitudinal development in the atmosphere as well as their lateral spread at ground level. The Observatory consists of 1600 water-Cherenkov detectors (SD), filled with 12 tonnes of water each and equipped with three photomultipliers to detect secondary photons and charged particles. The tanks are spread over 3000 km$^2$ on a triangular grid of 1.5 km spacing. The atmosphere above the array is viewed by 4 fluorescence detectors (FD), each housing 6 telescopes, located on the border of the area. The field of view of each telescope is 30° in azimuth, and 1.5 – 30° in elevation. Light is focused with a spherical mirror of 11 m$^2$ effective area on a camera of 440 hexagonal pixels. Each pixel is a photomultiplier tube with 18 cm$^2$ detection area. More details on detector setup and calibration can be found in $^6,^7$.

After entering the atmosphere, cosmic rays interact with nuclei in the air and start creating extensive air showers. The muons, electrons and photons that reach the ground are detected with the SD, their lateral spread from the air shower axis at primary energies above $10^{18}$ eV being in the order of a few kilometers. On the way through the atmosphere charged particles excite nitrogen molecules, which afterwards emit fluorescence light in the ultra-violet band. The amount of light is proportional to the energy deposited by the air shower in the atmosphere and is detected with the FD.

The SD has a high duty cycle of almost 100 %, but the energy calibration can be inferred in a model-independent way only from the FD energy assignment$^8$. The detected signal at 1000 m from the shower axis on the ground level, $S(1000)$, is a good estimator for the energy of the cosmic ray. Due to the attenuation in the atmosphere, $S(1000)$ depends on the zenith angle: an air shower developing vertically produces a smaller signal than an inclined shower produced by

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Figure 1: (left hand side) $S(1000)$ attenuation in the atmosphere. The line represents an empirical fit that is used for the conversion to $S_{38}$. (right hand side) Energy calibration. The events at low energies (below the dashed line) have been rejected to avoid threshold effects. The relation between $S_{38}$ and energy is almost linear and is shown with the continuous line$^3$. 

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a cosmic ray with the same energy. The constant intensity method \(^3\) is applied to obtain the zenith angle correction: it assumes that the cosmic ray flux is isotropic in local coordinates, i.e. the number of events above a certain threshold energy is constant as a function of \(\cos^2 \theta\). This hypothesis leads to the correction function for \(S(1000)\) shown in Fig. 1(left). A new variable obtained by using the empirical fit shown in the same figure, \(S_{38}\), represents the signal at 1000 m the very same shower would have produced if it had arrived from a zenith angle of 38\(^\circ\). This angle corresponds to the median of the zenith angle distribution of the SD data. The number of events above a certain \(S_{38}\) is zenith angle independent. In principle the attenuation might be energy dependent, because showers with higher energies develop deeper in the atmosphere and can be observed before their maximum development. This effect was found to be negligible.

The transformation from \(S_{38}\) to energy is obtained from high quality hybrid events. These are air showers that triggered both SD and FD, so \(S(1000)\) and FD energy have been reconstructed with good accuracy. The relation between the two variables, shown in Fig. 1 (right), exhibits a power law correlation with a relative dispersion of \(19 \pm 1\%\). The uncertainties in the determination of both FD energy and SD signal are assigned on event by event basis.

2 Ultra high energy cosmic ray flux

The data collected at the Pierre Auger are divided in three sets. The first set consists of air showers with zenith angle of less than 60\(^\circ\) detected by SD. The energy calibration described above is applied in this case. The integrated exposure reported here is 5165 km\(^2\) sr yr after some quality cuts (through February 2007), more than a factor of three larger than the exposure obtained by the largest forerunner experiment AGASA \(^5\). (Elsewhere in these Proceedings, Bonino reports other Auger results using an exposure through August 2007 of 9000 km\(^2\) sr yr). The acceptance is computed by simple geometrical considerations and from the continuous monitoring of the configuration of the array \(^9\). The data set used for obtaining the energy spectrum contains only events with energies greater than \(3 \cdot 10^{18}\) eV; above this energy the array is fully efficient.

The second set contains air showers measured by the SD with a zenith angle between 60\(^\circ\) and 80\(^\circ\). The procedure to derive the energy is equivalent to the vertical events, but instead of using \(S_{38}\) the shower size is determined from the relative distributions of the two-dimensional muon number densities at ground level. The normalization factor of the muon map, \(N_{19}\), is the estimator to be related to the hybrid energy. It gives the total number of muons relative
to a shower initiated by a proton with an energy of $10^{19}$ eV. The acceptance calculation is purely geometrical and the threshold energy above which the trigger efficiency is more than 98% is $6.3 \cdot 10^{18}$ eV. Above this energy the integrated exposure until the end of February 2007 is $1510 \text{ km}^2 \text{ sr} \text{ yr}$; 29% of the equivalent acceptance for vertical events.

The remaining set comprises of showers detected by the fluorescence detector and at least one SD unit. The hybrid exposure calculation relies on the simulation of the FD and SD response and it is energy dependent. A large sample of Monte Carlo simulations are performed to reproduce the exact conditions of the experiment and the entire sequence of given configurations, for example the rapidly growing array, as well as the seasonal and instrumental effects. The advantage of the hybrid measurement of the energy spectrum is the coverage of the energy range between $10^{18}$ eV and $3 \cdot 10^{18}$ eV.

The energy spectra obtained with the three methods are illustrated in Fig. 2 (left). The agreement is well within the independent systematic uncertainties, the difference between the overall normalizations is at a level of less than 4%. All spectra are affected by the 22% uncertainty in the FD energy scale, the main contributions coming from the determination of the fluorescence yield (14%), from the energy reconstruction itself (10%) and from the absolute calibration of the detector (9.5%). This systematic uncertainty does not affect the relative comparison of the three spectra. In order to obtain the Auger energy spectrum extending over the widest energy range possible, a maximum likelihood method is applied taking into account the independent uncertainties of each measurement.

In Fig. 2 (right) is illustrated the fractional difference between the Auger spectrum and a power-law $\propto E^{-2.6}$ which corresponds to the behavior of the energy spectrum between 18.6 and 19.6 in log(E/eV). Two spectral features are clearly visible: the so-called ankle at energies of $\approx 10^{18.5}$ eV and a flux suppression at energies above $\approx 10^{19.6}$ eV. The spectral index changes from $\gamma_1 = -3.30 \pm 0.06$ to $\gamma_2 = -2.62 \pm 0.03 \text{(stat)} \pm 0.02 \text{(sys)}$ at $\log(E_{\text{ankle}}/\text{eV}) = 18.65 \pm 0.04$, and above $10^{19.6}$ eV to $\gamma_3 = -4.14 \pm 0.42 \text{(stat)}$. A continuation of the energy spectrum as a power law with index $\gamma_2$ predicts 132 $\pm$ 9 events above $10^{19.6}$ eV and 30 $\pm$ 2.5 above $10^{20}$ eV, whereas we observe only 51 events and 2 events. The hypothesis of a pure power-law can be rejected with a significance of 6$\sigma$. 

Figure 3: The mean of $X_{\text{max}}$ distribution as a function of energy together with predictions from Monte Carlo simulations. Event numbers are indicated below each data point.
3 Mass composition

Spectral features alone cannot constrain acceleration models. Additionally the mass of the arriving particles has to be determined. Observables related to the shower development are used to identify the nature of the primary particles. At a given energy showers from a heavier nuclei develop earlier in the atmosphere, leading to different footprints in the array or the camera.

One variable sensitive to the composition of the cosmic rays is the slant depth position $X_{\text{max}}$ at which the maximum of the longitudinal profile occurs. Its average value is related linearly to the mean logarithmic mass, $\langle \ln A \rangle$, at a certain energy $E$: $\langle X_{\text{max}} \rangle = D_p [\ln (E/E_0) - \ln (\langle A \rangle)] + c_p$. $D_p$ is referred to as elongation rate and $c_p$ is the average depth of a proton with energy $E_0$.

Showers observed by at least one fluorescence detector and with at least one triggered tank were used to derive the mass composition of the cosmic rays. The mean $X_{\text{max}}$ as a function of energy together with predictions from air shower simulations is shown in Fig. 3 (left). The $X_{\text{max}}$ uncertainty is less than 20 g/cm². A moderate lightening of the primary composition is observed up to $\approx 2.5 \cdot 10^{19}$ eV, indicated by a slope larger than that of either pure proton or pure iron, whereas a constant mixed composition is present at high energies. The highest energy cosmic rays seem to develop higher in the atmosphere, indicating a heavy composition, or at least not a pure proton one. Larger statistics or independent analysis of the fluctuations of $X_{\text{max}}$ and SD mass composition estimators are needed to strengthen these results.

Photon induced showers have a strong individuality compared to the showers induced by nuclei. They develop slower in the atmosphere having a larger $X_{\text{max}}$. The main reason is the smaller multiplicity in the electromagnetic interactions compared to the hadronic ones combined with the LPM effect. Photon showers also contain fewer secondary muons, which combined with the deep penetration in the atmosphere leads to large rise times of the signal in the SD tanks and to a shower front with a larger curvature than hadronic showers. These two parameters are combined into a single SD observable through the principal component analysis to maximize the discrimination power. The Pierre Auger collaboration has derived a direct limit on the flux of photons for the first time by searching for photon candidates and relating their number to the exposure of the surface array. No photons have been found in the Auger data and therefore only limits on the photon fraction are shown in Fig. 4 (left). These limits improve significantly upon bounds from previous experiments, excluding some top down models as the super-heavy dark matter scenario. The flux expected for GZK photons will be reached with data accumulating.
over the next years.

In the case of the topological defects models, the UHECRs sources are distributed all over the universe and most of the high energy photons interact with the cosmic microwave background. The photon flux at Earth would be low in this scenario while the neutrino flux is not attenuated. The upper limit on the diffuse flux of ultra high energy tau neutrinos, presented in Fig. 4 (right), was built based on the search for neutrinos with the characteristics of extremely inclined, deeply penetrating events with a large electromagnetic component. The Pierre Auger Observatory is most sensitive to Earth-skimming \( \nu_\tau \) in the energy range where the GZK neutrinos are expected. The derived limit in the energy range \( 2 \cdot 10^{17} - 2 \cdot 10^{19} \text{eV} \) is at present the most sensitive bound.

4 Conclusions

The southern Pierre Auger Observatory will be completed at the end of 2008. Already with a data set that is comparable to the statistics of one year fully operational array, the hypothesis of a continuation of the energy spectrum in the form of a power law above an energy of \( 10^{19.6} \text{eV} \) is rejected with 6 sigma significance. This result is independent of the energy scale uncertainties. Combined with the directional correlation of the highest energetic cosmic rays with nearby active galactic nuclei the observed flux suppression suggests the existence of a GZK-effect. A mixed composition of the cosmic rays is present over the whole energy range and upper limits on photon and neutrino fluxes are given. The nature of the highest energies will be determined more precisely within the following year with increased statistics and different observables.

References

STUDY OF THE ULTRA HIGH ENERGY COSMIC RAYS ARRIVAL DIRECTIONS WITH THE PIERRE AUGER OBSERVATORY

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We present the first results about the studies of Ultra High Energy Cosmic Rays arrival directions using the early data collected at the Pierre Auger Observatory (corresponding to $\sim 1$ year of data taking of the complete southern array). We discuss in particular:
- the analysis of large-scale patterns in the arrival directions of cosmic rays;
- a search for an excess of events from the direction of the Galactic Center region and from some extragalactic objects;
- the observed correlation between cosmic rays with energies above 60 EeV ($1 \text{ EeV} = 10^{18} \text{ eV}$) and the directions of nearby active galactic nuclei (AGN).

1 Introduction

The identification of the sources of ultra-high energy cosmic rays and the comprehension of the mechanisms by which they acquire their energies have been great challenges since the detection of the first $10^{20} \text{ eV}$ event in 1962 at Volcano Ranch.

The maximum energy attainable in an accelerator with characteristic magnetic field $B$ and size $L$ is of order $E_{\text{max}} \sim Z e B L$. Only a few types of astronomical objects appear able to accelerate protons to $10^{20} \text{ eV}$; these include Active Galactic Nuclei, galaxy clusters, and objects with large radio lobes.

Furthermore, particles with energies above about $6 \cdot 10^{19} \text{ eV}$ are expected to interact inelastically with cosmic microwave background photons, losing energy at each interaction. As a consequence the cosmic ray flux may be significantly reduced above 100 EeV. Particles exceeding the interaction energy threshold and originating at distances greater than 100 Mpc should never be observed on Earth. This effect, known as the "GZK effect" \cite{3,4}, requires the sources of the cosmic rays observed at Earth to be relatively nearby, within about 100 Mpc at most, further reducing the number of possible candidates.

Among the observables that might help to solve the puzzle of the sources, one of the most effective is the study of anisotropy in the UHECR arrival directions. In air-shower experiments the incoming directions of the highest energy cosmic rays are determined well and hence it is possible to estimate whether or not they are isotropically distributed on the sky. At the highest energies ($> 5 \cdot 10^{19} \text{ eV}$) the arrival directions point back to the sources because these particles should be only slightly deflected by magnetic fields.

In anisotropy studies, especially on small angular scales, it is fundamental to determine the arrival direction of cosmic rays with great precision. Consequently, an accurate knowledge of
the angular resolution of the Auger Surface Detector (SD) is required. We discuss this in section 2, followed by a presentation of results on large- and small-scale anisotropy. The first specific targets chosen by the Auger Collaboration have been the Galactic Center at EeV energies and BL-Lacs and AGN at higher energies.

2 Angular resolution of the Surface Detector

The arrival direction of a SD event is determined by fitting the arrival time of the first particle in each station to a shower front model (see fig. 1). The precision achieved in the arrival direction reconstruction depends therefore on the uncertainty in the time measurement and on the effectiveness of the shower front model adopted.\(^5\)

![Figure 1: Sketch of the shower front arrival.](image)

![Figure 2: SD angular resolution as a function of zenith angle for different station multiplicities.](image)

The angular resolution is calculated on an event by event basis, from the zenith (\(\theta\)) and azimuth (\(\phi\)) uncertainties of the geometrical reconstruction. It is defined as the angular radius that would contain 68% of showers coming from a point source.

In fig. 2 the angular resolution is shown as a function of the zenith angles for various station multiplicities.\(^6\) It is better than 2° in the worst case of vertical showers with only 3 stations hit and improves significantly with the number of stations. For events with 6 or more stations, corresponding to events with energies above 10 EeV, it is always better than 1°.

3 Large scale anisotropy studies

Lower energy cosmic rays likely originate within our Galaxy, while higher energy particles are believed to be extragalactic. At the transition the large scale angular distribution might change significantly. Large scale anisotropy, especially its evolution with primary energy, represents one of the main tools for discerning between the galactic and extragalactic origin of cosmic rays and for understanding their mechanisms of propagation.

If the transition to extra-galactic sources occurs at the ankle of the spectrum,\(^7\) then at \(10^{18}\) eV cosmic rays are still mainly galactic and their diffusive escape from the Galaxy may be efficient enough so that the sky distribution of their arrival directions is not isotropic. The predictions for the shape and amplitude of the corresponding anisotropy are very model-dependent, but a %-level modulation is plausible.\(^8\)

On the other hand, if the transition occurs at lower energy, i.e. around \(5 \cdot 10^{17}\) eV, then \(10^{18}\) eV cosmic rays are already extragalactic and their sources may be cosmologically distributed. If
so then no large-scale pattern would be detectable except for the CMB-like dipole anisotropy\textsuperscript{10}. In this case anisotropy amplitudes of the order of $\sim 0.6\%$ are expected.

\subsection{Auger results}

The statistics accumulated so far by the Auger Observatory permits the study of $\%$-level large-scale patterns, but this is challenging due to the difficulty of controlling the sky exposure of the detector and various acceptance effects, such as detector instabilities and weather modulations.

In order to avoid such problems three complementary analyses have been performed. All show that at EeV energies the Right Ascension (RA) distribution is remarkably compatible with an isotropic sky; an upper limit on the first harmonic modulation of 1.4\% in the energy range $1 < E < 3$ EeV has been set\textsuperscript{11} (see fig.3 for more details). This result does not confirm the 4\% RA modulation found by the AGASA experiment\textsuperscript{12} (although the sky regions covered by the two experiments are different) and already sets some constraints on the galactic hypothesis (further statistics and analysis are in any case necessary).

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{figure3.png}
\caption{Overview on the results of large scale anisotropy studies; Auger upper limits are drawn in red.}
\end{figure}

\section{The Galactic Center region}

The Galactic Center is one of the most interesting targets in the study of small scale anisotropies at EeV energies because it contains a super massive black hole, a good candidate accelerator of high-energy cosmic rays. This black hole is believed to be associated with the radio emissions from Sagittarius A*. The H.E.S.S. collaboration has recently observed TeV $\gamma$-ray emissions close to this radio source\textsuperscript{13}. A further reason of interest for this region is the privileged position of the Pierre Auger Observatory: the GC passes only 6\degree away from the observatory zenith.

In the past there have been claims of excesses of cosmic rays from the GC region from the AGASA\textsuperscript{12} and SUGAR\textsuperscript{14} experiments. Both the excesses are located in regions near the GC but not coincident with it (in the case of AGASA the GC is not in its field of view).

\subsection{Auger results}

Besides the privileged position, another advantage for Auger comes from the exposure of the array: the number of EeV cosmic rays accumulated so far from this part of the sky greatly exceeds that from previous experiments.
The claims of the forerunner experiments are periodically tested by the Auger experiment in different energy ranges and window sizes. In the most recent analysis two different energy ranges have been considered, 0.1-1 EeV and 1-10 EeV, but no significant flux excess has been found in the region around the GC (see tab.1: the numbers of observed events are always compatible with the expected ones)\textsuperscript{15}. The distribution of Li-Ma significances for overdensities in this region is consistent with an isotropic sky for both energy ranges.

Table 1: Summary of excesses searches for $0.1 < E < 1$ EeV (top) and $1 < E < 10$ EeV (bottom) around the GC in the form of both extended and point-like source.

<table>
<thead>
<tr>
<th>search</th>
<th>window size</th>
<th>$n_{\text{obs}}/n_{\text{exp}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>extended</td>
<td>$10^\circ$ (TH)</td>
<td>5663 / 5657 = 1.00 ± 0.02(stat) ± 0.01(syst)</td>
</tr>
<tr>
<td></td>
<td>$20^\circ$ (TH)</td>
<td>22274 / 22440 = 0.99 ± 0.01(stat) ± 0.01(syst)</td>
</tr>
<tr>
<td>point-like</td>
<td>$1.3^\circ$ (G)</td>
<td>192.1 / 191.2 = 1.00 ± 0.07(stat) ± 0.01(syst)</td>
</tr>
</tbody>
</table>

5 Correlation of UHECR with nearby extra-galactic objects

At the highest energies, above a few $\times 10^{19}$ eV, cosmic rays should be only slightly deflected by magnetic fields. A direct way to search for sources of UHECR is to analyze the distribution of their arrival directions for small-scale clustering and specifically to search for correlations with known astronomical objects that are candidate sources.

5.1 Active Galactic Nuclei

AGN have long been considered to be sites where energetic particle production might take place, and where protons and heavier nuclei could be accelerated up to the highest energies measured so far.

The Auger collaboration searched for a correlation of its highest energy events with these astronomical objects (the selected AGN come from the 12th edition of Véron-Cetty/Véron catalogue\textsuperscript{16}). The data set analyzed has been acquired by the surface array during the first 3.5 years of data taking and corresponds to an integrated exposure of about 9000 km$^2$ sr yr.

Under the assumption of isotropy, it’s possible to calculate the probability that any given cosmic ray falls within a fixed distance from any AGN, i.e. the probability $P$ for a set of $N$ events from an isotropic flux to contain $k$ or more events at a maximum angular distance $\psi$ from any member of a collection of candidate point sources. $P$ is given by the cumulative binomial distribution $\sum_{j=k}^{N} C_{N}^{j} p^{j}(1-p)^{N-j}$, where $p$ is the fraction of the sky (weighted by the exposure) defined by the regions at angular separation less than $\psi$ from the selected sources. The degree of correlation has been computed as a function of three parameters: the maximum angular separations $\psi_{\text{max}}$, the maximum AGN distance $D_{\text{max}}$ and energy thresholds $E_{\text{th}}$.

The strategy adopted in this analysis requires as a first step an exploratory scan for the minimum of $P$, aimed to identify the configuration of parameters that maximizes the correlation. The absolute minimum of $P$ was found for $\psi_{\text{max}} = 3.1^\circ$, $z_{\text{max}} = 0.018$ ($D_{\text{max}} = 75$ Mpc) and $E_{\text{th}} = 56$ EeV. In this optimized configuration 12 events among 15 correlated with the selected AGN, while only 3.2 were expected by chance if the flux were isotropic.
To avoid the negative impact of trial factors in a posteriori search, the Auger collaboration decided to test this hint of anisotropy on an independent data set with parameters specified a priori. A prescription was written down, fixing the set of parameters and sources; new data would be analyzed sequentially until the probability to incorrectly reject isotropy was 1% and the probability to incorrectly reject correlation was 5%.

This prescription was tested on an independent data set collected after 27 May 2006 (when the prescription started), with exactly the same reconstruction and calibration algorithms as in the exploratory scan. On 25 May 2007, 6 out of 8 events were found to fulfill the prescription.

After the successful result of this test, a re-scan of the full data set (from 1 January 2004 to 31 August 2007) was performed, adopting newer and somewhat more accurate reconstruction and calibration algorithms. A similar result is obtained, with the correlation maximized for the 27 events with energies above 57 EeV: 20 of these events correlate with at least one AGN for a maximum angular separation $\psi_{\text{max}} = 3.2^\circ$ and a maximum distance to AGN $D_{\text{max}} \sim 71$ Mpc. The results of this analysis are shown in fig.4.

Figure 4: Projection on the celestial sphere in galactic coordinates with circles of radius $3.2^\circ$ centered at the arrival directions of the 27 cosmic rays with highest energy detected by Auger. The positions of the selected 472 AGN are indicated by red asterisks.

Summarizing, the most important results of this analysis are:

- The anisotropy of UHECR (above 57 EeV) has been confirmed at 99% CL with an a priori test on an independent data set (in a detailed report is given). This is the first time that a so strong signal of correlation is revealed and the hypothesis of an isotropic distribution of these cosmic rays is rejected at such a confidence level.

- The observed correlation is compatible with the hypothesis that UHECR originate from extra-galactic sources within the GZK horizon (i.e. compatible with the flux suppression observed in the spectrum starting at $\sim 60$ EeV).

- The angular scale of the correlation is a few degrees, suggesting a predominantly light composition.

- AGN are the tracers but cannot be identified unambiguously as the sources: objects with a similar spatial distribution (GRB, quasar remnants, ...) are not excluded. It is also plausible that only a subclass of AGN are the sources.
Several events lie close to the super-galactic plane (particularly close to Cen A) whereas a paucity of events has been recorded from Virgo.

The Véron-Cetty/Véron catalogue chosen for the correlation search is one of the largest collection of such objects but it is not an unbiased statistical sample. It is incomplete around the Galactic plane and for objects distances greater than 100 Mpc. It is important to note that these flaws are not an obstacle to the limited aim of demonstrating anisotropy.

A significant increase in ultra-high energy cosmic-ray statistics, combined with the future northern site of Auger, should lead to an unambiguous identification of the sources and their characteristics.

5.2 BL-Lacs

Active Galactic Nuclei include different sub-classes of astronomical objects. One of the more attractive candidate classes for UHECR sources is that of BL-Lacs. These are blazars in which the relativistic jet axis of the active galaxy is aligned with our line of sight.

Significant correlations of arrival directions of UHECR with positions of BL-Lacs were found by forerunner experiments with different subsets of BL Lacs and setting different energy thresholds.

A test on all these correlations has been performed with the present Auger data set which is already 6 times larger than those used in preceding cross-correlation searches for energies above 10 EeV. The results of this test\(^{19}\) do not support previously reported excesses of correlation since the number of correlations found is compatible with that expected for an isotropic flux.

References

A COMMENT ON CATALOG SEARCHES

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We illustrate in a concrete example that a mere positional correlation of highest-energy cosmic rays with active galactic nuclei (AGN), although suggests, does not necessarily imply that the latter are sources of the cosmic rays. Different interpretations of this correlation are possible, and signatures other than positional correlations are needed to discriminate between them. We point out that some of these signatures seem to disfavor the AGN interpretation with already existing data.

In this talk I would like to clarify two points related to the correlations between the ultra-high energy cosmic rays (UHECR) and nearby active galactic nuclei (AGN) from the catalog, which were recently found by the Pierre Auger Observatory (PAO).

1 Contrary to naive expectation, a correlation of cosmic rays with AGN (or any other objects) does not automatically imply that the latter are cosmic ray sources. This is not related to the significance of the correlation, but follows from the very nature of the statistical test performed to establish the correlation. In positional correlation analysis one compares the distribution of the data events over the sky with the isotropic distribution. If the two distributions are found to be incompatible, this means simply that the data are not isotropic. The actual sources should be identified by different methods.

To illustrate the relevance of this point consider a concrete example. The same set of cosmic ray events which correlate with AGN in the PAO analysis may be cross-correlated, by the same method, with just one object for which we take Cen A, an active galaxy in the direction of the Centaurus supercluster. Cen A is a radio-galaxy which is exceptionally close to us: the distance to Cen A is about 3.5 Mpc. It possesses jets and radio lobes, the usual attributes of a potential acceleration site.

There is an excess of events in the data in the direction of Cen A. The significance of the excess at a given angular scale δ can be characterized by the probability P(δ) that equal or larger excess occurs by chance as a result of a fluctuation in the uniform distribution. The smaller is the probability to obtain a given excess by chance, the more significant it is. This probability may be determined by
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Figure 1: The probability $P$ that the observed excess of events within angular distance $\delta$ around Cen A has occurred by chance. The values of $P$ are indicative only since their calculation accounts neither for the statistical penalty associated with the choice of angular scale nor for the bias in the sample.

the Monte-Carlo simulation. The result of the simulation is shown in Fig. 1. One can see that the excess is most significant at about $20^\circ$. Out of 27 events in total, 9 events fall within $20^\circ$ from Cen A while only 1.5 are expected for the uniform distribution. Note that the events contributing to this correlation with Cen A are the same events that contribute to the correlation with AGN if the latter are assumed to be sources.

Such a situation is explained in the following way. The distribution of the nearby AGN is rather inhomogeneous. Moreover, Cen A is projected onto one of the largest nearby structures, the Centaurus supercluster, as can be seen on Fig. 2. For this reason, the same data show correlations with both Cen A and AGN. Importantly, if either AGN or Cen A are indeed sources of highest-energy cosmic rays, both correlation signals will increase with the accumulation of statistics. So, a mere increase of significance will not allow to discriminate between the two possibilities.

It follows from the above that alternative signatures are needed to distinguish between the two cases. We present here one of such signatures.

The idea is that the cosmic ray flux predicted by the AGN hypothesis can be computed and compared to the observed one. In this way the AGN hypothesis itself will be subject to a test, not the hypothesis of the isotropic distribution.

The computation can be performed in a straightforward way taking into account the distance to AGN and the attenuation of protons of different energies (see Refs.\textsuperscript{7,8} for details). The results are presented in Fig. 2 in the form of red crosses which show the positions of the nearby AGN. The intensity of a cross represents its expected contribution to the flux. This figure should be understood in a statistical sense: the fluxes of individual sources cannot, of course, be predicted without the detailed modeling of corresponding AGN (for which modeling there is probably not enough information anyway). However, in large groups of galaxies like galaxy clusters individual differences in luminosity will average away and only the common factors determined by the distance will remain. The relative contributions to the total flux from such groups can thus be reliably predicted.

One can observe the overdensity of the events in the direction of the Centaurus supercluster. The second region where a high flux is expected, the Virgo cluster, is completely devoid of events. This is a strange feature that does not look compatible with the AGN hypothesis.

The latter statement can be quantified by comparing the expected and observed distributions of events in the angular distance from the center of the Virgo cluster, as well as their distributions in Galactic and supergalactic longitudes and latitudes. The comparison may be performed by the Kolmogorov-Smirnov test. The results of different tests show different degree of incompatibility between the predicted and observed distributions with the probability that it has occurred as a result of a
fluctuation varying from 10% to $10^{-4}$. Taking into account the strongest discrepancy and the number of tests performed, we estimate the significance of the tension between the AGN hypothesis and the data to be of order 99%.

One of the drawbacks of the analysis just described is the incompleteness of the AGN catalog. To check how much our results depend on this incompleteness we have replaced the catalog of AGN by a complete catalog of galaxies containing objects up to 270 Mpc. The above tests performed with the AGN catalog replaced by the complete galaxy catalog show similar results. We think therefore that incompleteness of the catalog is not an issue.

Another drawback, which unfortunately cannot be avoided at present, is the a posteriori nature of the tests performed. To avoid this problem, the tests which we have described will have to be repeated with the new independent data. This is why now, before the new data arrive, it is particularly important to formulate other hypotheses and procedures to test them which may then be performed in a more reliable a priori way with independent data sets.

**Acknowledgments**

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XLIId RENCONTRES DE MORIOND

Electroweak Interactions and Unified Theories

VI - Young Scientists Forum
The Z Transverse Momentum Distribution

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The DØ collaboration has measured the transverse momentum of the Z boson in the process $Z/\gamma^* \rightarrow ee$. This is compared to a theoretical model of the boson transverse momentum and found to agree well with the existing parametrization. This parametrization is a critical input in the measurement of the W boson mass measurement, which is discussed briefly.

1 Introduction

At leading order W and Z bosons are produced at the Tevatron predominately through the Drell-Yan diagrams shown in figures 1 and 2, where the only particle in the final state is the boson. Without anything to recoil against the boson is necessarily produced at rest in the $q\bar{q}$ frame. The partons themselves are assumed to have very little momentum transverse to the beam direction (from the confinement of the parton within the proton). Therefore the boson is produced with effectively no $p_T$.

Figure 1: W boson production and decay at tree level.

Figure 2: Z/γ* production and decay at tree level.

Higher order quantum chromodynamics (QCD) processes involve additional particles in the final state. The most straightforward to consider are the initial state radiation of a gluon (ISR)
and the Compton radiation of a gluon (figure 3). This additional particle gives the boson something to recoil against. Typically this results in a boson $p_T$ of several GeV.

![Figure 3: W boson production at next to leading order. Here the boson recoils against the gluon, resulting in non-zero boson transverse momentum.](image)

## 2 Boson Production

A complete discussion of the $p_T$ distribution of electroweak bosons at production is a significant undertaking, the full details of which are beyond the scope of this discussion. However, because of the critical role the description of boson production plays in the measurement of the W boson mass, as well as searches for production of the Higgs boson and other beyond the Standard Model physics, we describe several of the important phenomenological details. The differential cross section with respect to the boson transverse momentum is written as

$$\frac{d^2\sigma}{dp_T^2} = \sum_{ij} \int dx_1 dx_2 f_i(x_1)f_j(x_2) \frac{d^2\sigma(ij \rightarrow V)}{dp_T^2}. \tag{1}$$

For the process $q\bar{q} \rightarrow W\bar{q}$ shown in figures 3 one can accurately calculate the transverse momentum distribution at high $p_T$ ($p_T \approx M_{\text{Boson}}$) using perturbative QCD. In the low $p_T$ region the results are divergent. A framework was developed by Collins, Soper and Sterman to re-sum the perturbation series by grouping the divergent, non-perturbative terms together. This resummation works for all orders in perturbation theory but requires a correction at low-$p_T$ where non-perturbative physics becomes important. The correction is parameterized with a function $W$ that smoothly turns off as the boson $p_T$ increases. The parton level cross section is then

$$\frac{d^2\sigma(q\bar{q} \rightarrow V)}{dp_T^2} \sim \int_0^\infty \left[d^2b \exp(i\mathbf{p}_T \cdot \mathbf{q}) \times W(b,Q)\right] + Y(P_T,Q) \tag{2}$$

where $Y$ is the perturbative piece and the impact parameter $b$ is the conjugate variable of $p_T$ (as $b$ increases $p_T$ decreases). The form of the $W$ function is phenomenologically motivated and is determined by fitting data from several experiments. Recently the form

$$W_{NP}(b) = \exp\left(-\left(g_1 + g_2 \ln\left(\frac{Q}{2Q_0}\right) + g_1g_3 \ln(100x_1x_2)\right) b^2\right) \tag{3}$$

was proposed by Brock, Landry, Nadolsky and Yuan, where in this case

$$Q \sim 91 \text{GeV}, \quad Q_0 = 1.6 \text{GeV}, \quad x_{i,j} \sim 0.05 \tag{4}$$

and the parameters

$$g_1 = 0.21 \pm 0.01 \text{ GeV}^2, \quad g_2 = 0.68^{+0.01}_{-0.02} \text{ GeV}^2, \quad g_3 = -0.60^{+0.05}_{-0.04} \tag{5}$$

are determined through global fits. In the kinematic region we are interested in the $g_2$ term dominates. Independent of the global fits the DØ experiment uses the transverse momentum spectrum of $Z \rightarrow ee$ decays to determine $g_2$.  

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3 \textit{Z Transverse Momentum Analysis}

The event selection used to collect \textit{Z} boson candidates requires two isolated electromagnetic (EM) clusters that have a shower shape consistent with that of an electron and are away from the module boundaries of the calorimeters. Electron candidates are required to have transverse momentum greater than 25 GeV. The pairs of electrons must have an invariant mass $70 < M(ee) < 110$ GeV. If both electrons are in the central calorimeter, then each electron must be spatially matched to a reconstructed track. The central section of the calorimeter has coverage for electrons of $|\eta| < 1.1$ and two endcap calorimeters have an approximate coverage of $1.5 < |\eta| < 3.2$ for electrons ($\eta = \ln[\tan(\theta/2)]$). In figure 4 the Z boson $p_T$ distribution is shown corrected for efficiencies, background and acceptance and then unfolded to obtain the true differential cross section. The curve is the ResBos prediction using the BLNY parametrization.

We fit for the parameter $g_2$ using the distribution in figure 4 and find the best fit value to be $g_2 = 0.77 \pm 0.06$. This is then used as an input in the measurement of the W mass.

Figure 5: The $\chi^2$ as a function of the $g_2$ parameter using the unfolded data and the BLNY model. The solid line is a polynomial fit.
4 Implications for W Mass Measurement

To measure the W boson mass using the $W \rightarrow e\nu$ decay channel a theoretical description of the boson’s transverse momentum is necessary. This is because the invariant mass cannot be reconstructed from the electron/neutrino final state. Instead, the sensitivity of the electron transverse momentum and the transverse mass (analogous to the invariant mass) distributions to the W boson mass are exploited. The electron $p_T$ and the transverse mass distributions cannot be described analytically so a theoretical description of the boson production and decay is employed, in conjunction with a parameterized detector simulation. Using the parametrization in [5] with the value for the $g_2$ parameter determined above we estimate the contribution to the overall uncertainty on the W boson mass at DØ (using approximately 1 fb$^{-1}$ of data) to be 5 MeV for the transverse mass distribution and 16 MeV for the transverse momentum distribution. With a total estimated statistical uncertainty of 17 MeV using the transverse mass distribution (23 MeV using the electron $p_T$ distribution) for this sample this aspect of the theoretical description is not the dominate uncertainty.

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References

Precision Top Quark Mass From a Simultaneous Fit in Lepton + Jets and Dilepton Channels Using 2 fb$^{-1}$ of data collected by the CDFII detector

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We present a preliminary measurement of the top quark mass employing the template method with data sample collected by the CDF Run II detector corresponding to integrated luminosity of 2 fb$^{-1}$. Lepton + Jets and Dilepton final states are selected. For each event in the Lepton + Jets channel we apply kinematic constraints on the pair of top quarks and their decay products to determine a reconstructed top quark mass. We simultaneously determine the invariant mass of the decaying W boson to calibrate the energy response of the detector. The events in the Dilepton sample are reconstructed using the Neutrino Weighting Algorithm. To improve the precision, for each Dilepton event we calculate $H_T$ - the linear sum of missing transverse energy and transverse momenta of jets and leptons. The reconstructed top quark mass and W boson invariant mass distributions from the Lepton + Jets channel and reconstructed top quark mass and $H_T$ distributions from the Dilepton channel are fit to Monte Carlo derived templates in a likelihood fit to extract the top quark mass and an in-situ measurement of the jet energy scale. We measure $M_{\text{top}} = 171.9 \pm 2.0$ GeV/$c^2$.

1 Introduction

Since it’s discovery$^{1,2}$ at the Tevatron the top quark has been one of the most studied fundamental particles. It is more than an order of magnitude heavier than the next heaviest Standard Model fermion. This points to it’s crucial role in a puzzle of the origin of mass. The top quark and the Higgs boson contribute in the loop corrections to the W boson mass, therefore knowing precisely the top quark mass and the W boson mass allows to constrain indirectly the Higgs boson mass$^3$. Once the Higgs boson is discovered the knowledge of it’s mass together with measurements of the top quark mass and the W boson mass will provide a sensitive test of the Standard Model$^4$. Current measurements of the top quark mass and the W boson mass may be giving us hints on the nature of physics beyond the Standard Model$^5$. In this letter we present a preliminary top quark mass measurement using the Lepton + Jets and Dilepton decay channels simultaneously. This approach is applied for the first time in a top quark mass measurement. More details on this analysis can be found in$^6$.

2 Top Quark Production and Decay

Top quarks are produced at the Tevatron mainly in quark-antiquark annihilation events where a gluon is produced, splitting into a $t\bar{t}$ pair. Each of the top quarks then decays into a $W$ boson and a $b$ quark with essentially 100% branching fraction. The $W$ bosons can decay into a quark pair or a charged lepton-neutrino pair, giving rise to classification of the $tt$ decays into three
classes. Thus we have an All-hadronic decay channel with six jets in the final state, a Lepton + Jets decay channel where we find four jets, one lepton and missing transverse energy and a Dilepton decay channel characterized by two leptons, two jets and missing transverse energy. Due to difficulty of reconstructing $\tau$ leptons we restrict the meaning of lepton to an electron or a muon.

3 Combination Strategy

Measurements in all decay topologies are valuable as statistically independent cross-checks and are all needed to obtain best precision possible. Traditionally a dedicated analysis is performed in each channel and the results are combined using an averaging technique\(^7\). In any such combination one must assume the values of correlations in systematic effects between the measurements in different channels. A form of the likelihood shape is also required as an input and is usually assumed to be Gaussian. In this letter we present a preliminary top quark mass measurement using two decay channels simultaneously. The analysis presented here allows us not to make any assumptions mentioned above, yielding a more robust measurement.

4 The Jet Energy Scale

In the top quark mass measurements a major source of uncertainty is the modelling of the jet calibration or the jet energy scale (JES). Multiple effects contribute to the uncertainty on the jets\(^8\). A major uncertainty arises from modelling of nonlinearities of the calorimeter and energy loss in uninstrumented regions (absolute energy scale). Flow of particles outside of the jet cone (out of cone energy scale) gives large uncertainty especially for low energy jets. Another large systematic uncertainty arises from detector nonuniformity as a function of the pseudorapidity (relative energy scale). Interactions of the spectator partons (underlying event energy scale) and additional soft $p\bar{p}$ interactions in the same bunch crossing are sources of small systematics. We measure the offset from a nominal calibration in units of the total systematic uncertainty on the calibration $\sigma_c$. In Lepton + Jets channel presence of a hadronically decaying $W$ boson allows us to calibrate in-situ the value of the shift $\Delta_{JES}$ from the nominal JES. Since the measurement is performed in two channels simultaneously this calibration will be applied uniformly to the two decay channels used.

5 Event Selection

5.1 Lepton + Jets Channel

To select the Lepton + Jets sample we require at least four jets with high transverse energies. At least one of the jets has to be identified as a $b$ quark jet based on a secondary vertex or a “$b$-tag”. We separate the Lepton + Jets sample into 1-tag and 2-tag subsamples. In the 1-tag sample we require that there are exactly four jets with transverse energies greater than 20 GeV when corrected to the particle level. In the 2-tag samples we relax the energy for the fourth most energetic jet to have $E_T > 12$ GeV. We also allow additional jets in the event. We require a central electron or a muon with $E_T$ or $p_T > 20$ GeV. The missing transverse energy must be greater than 20 GeV.

The background estimate for the Lepton + Jets samples is obtained form combination data - Monte Carlo technique. The major backgrounds arise from production of $W$ boson in association with heavy flavour jets and light flavor jets where the light flavour jet is tagged (so called “mistag”) and from QCD events where one of the jets is misidentified as a lepton. We expect in the 1-tag sample $42.7 \pm 12.5$ and in the 2-tag $4.2 \pm 1.9$ background events.
5.2 Dilepton Channel

We require at least two jets with $E_T > 15 \text{ GeV}$. Two leptons of opposite charge must be present with transverse energies of at least 20 GeV. If the leptons are of the same flavor we impose the requirement that their invariant mass lies at least 15 GeV/$c^2$ from the $Z$ boson mass. Additionally we require that $H_T > 200 \text{ GeV}$, $E_T > 25 \text{ GeV}$ where $H_T$ is a linear sum of $E_T$ and transverse energies of jets and leptons. Topological cuts designed to remove events where $E_T$ arises due to instrumentalational effects or $\tau$ production are applied. The Dilepton sample is divided into two subsets: a 0-tag sample and a 1-tag sample.

The background contributions to the Dilepton channel include events where a lepton is produced in association with jets and one of the jets is reconstructed as a lepton (“Fakes”), Drell-Yan production and diboson production. The Fakes background is estimated from data while other backgrounds are estimated using data-Monte Carlo and Monte Carlo only techniques. In the non-tagged sample we expect $31.1 \pm 5.6$ and in the tagged sample $2.4 \pm 0.6$ background events.

5.3 Event Reconstruction

In each event we form a reconstructed top quark mass, a variable which is highly sensitive to the true top quark mass. In the Lepton + Jets channel we use a $\chi^2$ fit where the magnitudes of lepton and jet momenta and the transverse components of unclustered energy are allowed to float within their resolution around the observed values. We impose a constraint that the invariant masses of the neutrino-lepton system and the light quark system are close to the measured $W$ boson mass. The invariant mass of the leptonically decaying top quark daughters is constrained to be within the theoretical top quark width from the invariant mass of the hadronically decaying top quark daughters. The constraint is imposed through fit parameter taken to be the reconstructed top quark mass. The $\chi^2$ minimization is performed for all jet-to-quark assignments consistent with b-tagging combination and the combination with lowest minimum $\chi^2$ is used. To form a reconstructed top quark mass in the Dilepton channel events we use the Neutrino Weighting Algorithm. We scan a range of top quark masses. At each point in the scan we integrate over the pseudorapidities of the two neutrinos and sum over the two possible jet-to-quark assignments. Knowing the top quark mass, neutrino pseudorapidities and masses of all particles in the decay cascade we solve for the neutrino transverse momenta. The integrand is formed by a Gaussian weight that compares the measured $E_T$ value to the solution obtained for the neutrino transverse momenta. The top quark mass in the scan that yields the highest weight is taken as the reconstructed top quark mass in this event. Additionally in each Dilepton event we calculate the $H_T$. In the Lepton + Jets channel we reconstruct also the invariant mass of the hadronically decaying $W$ boson. As mentioned above this variable captures the shifts in JES.

5.4 Mass Fitting

We employ a template approach in this analysis. We generate $t\bar{t}$ Monte Carlo samples with a range of top quark mass and JES shifts. We also construct background models using data and Monte Carlo samples. We form probability density functions (pdf) for the observables mentioned above and compare them to the distributions of the observables obtained from data in an extended likelihood fit, to obtain a measurement of the top quark mass $M_{t_{\text{top}}}$ and the jet energy scale shift $\Delta_{\text{JES}}$. The probability density functions are constructed using the Kernel Density Estimation (KDE) techniques. In this approach the probability for an event to have certain values of the observables is calculated as a sum of values of kernel functions from all events in a given Monte Carlo sample. This technique treats intrinsically the correlations.
between observables. KDE gives the value of signal pdf at distinct values of $M_{\text{top}}$ and $\Delta_{\text{JES}}$ where Monte Carlo samples were generated. To obtain a pdf that varies smoothly as a function of those two parameters we use Local Polynomial Smoothing (LPS) \(^{12}\). LPS performs a fit to a parabolic function using the KDE estimates from Monte Carlo templates with the $M_{\text{top}}, \Delta_{\text{JES}}$ parameters lying close to the point where the estimate is desired. The value of the parabola at that point is interpreted as the pdf.

Using the distributions of observables in data, the negative log-likelihood is minimized for the top quark mass of $171.9 \pm 1.7$ (stat.+JES) GeV/$c^2$. Fitted $\Delta_{\text{JES}}$ value is consistent with nominal calibration of 0 $\sigma_c$.

6 Systematics

The largest systematic (0.6 GeV/$c^2$), $b$ quark jet energy scale arises due to differences in modelling $b$ and light flavour jets. As described in section 4 many effects contribute to uncertainty on JES. Modelling the offset from the nominal calibration as just one number $\Delta_{\text{JES}}$ gives source to the residual JES uncertainty of 0.5 GeV. Another large systematic (0.5 GeV/$c^2$) is due to the modelling of the initial and final state radiation. Additional systematics include generator differences (0.2 GeV/$c^2$), background shape (0.1 GeV/$c^2$), Monte Carlo sample statistics (0.1 GeV/$c^2$), lepton energy scale (0.1 GeV/$c^2$) and multiple $p\bar{p}$ interactions (0.1 GeV/$c^2$). Total systematic uncertainty is 1.0 GeV/$c^2$.

7 Conclusions

We performed the first top quark mass measurement simultaneously in two decay channels treating the correlations in the systematic effects intrinsically. No assumptions on the form of the likelihood needed to be made and the JES calibration was applied uniformly to both channels. The result obtained is:

$$M_{\text{top}} = 171.9 \pm 1.7 \text{ (stat. + JES)} \pm 1.0 \text{ (other syst.) GeV}/c^2 = 171.9 \pm 2.0 \text{ GeV}/c^2$$

References

BiPo PROTOTYPE FOR SuperNEMO RADIOPURITY MEASUREMENTS

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The BiPo project is dedicated to the measurement of extremely low radioactive contaminations of SuperNEMO $\beta\beta$ source foils ($^{208}\text{Tl} < 2 \mu\text{Bq/kg}$ and $^{214}\text{Bi} < 10 \mu\text{Bq/kg}$). A modular BiPo1 prototype with its 20 modules and its shielding test facility is running in the Modane Underground Laboratory since February, 2008. The goal of this prototype is to study the backgrounds and particularly the surface contamination of plastic scintillators. After 2 months, a preliminary upper limit on the sensitivity of a 10 m$^2$ BiPo detector in $^{208}\text{Tl}$ contamination of selenium source foils can be extrapolated to: $A(^{208}\text{Tl}) < 7.5 \mu\text{Bq/kg}$ (90 % C.L.).

1 Principle

The BiPo detector is dedicated to the measurement of the high radiopurity levels in $^{214}\text{Bi}$ and $^{208}\text{Tl}$ of very thin materials and especially the double beta source foils of the SuperNEMO detector$^1$. The expected sensitivity is $A(^{208}\text{Tl}) < 2 \mu\text{Bq/kg}$ and $A(^{214}\text{Bi}) < 10 \mu\text{Bq/kg}$.

In order to measure $^{208}\text{Tl}$ and $^{214}\text{Bi}$ contaminations, the original idea of the BiPo detector is to detect the so-called BiPo process, which corresponds to the detection by organic scintillators of an electron followed by a delayed alpha particle. The $^{214}\text{Bi}$ isotope is nearly a pure $\beta$ emitter ($Q_\beta = 3.27$ MeV) decaying into $^{214}\text{Po}$, an $\alpha$ emitter with a half-life of 164 µs (Fig. 1). The $^{208}\text{Tl}$ isotope is measured by detecting its parent the $^{212}\text{Bi}$ isotope. $^{212}\text{Bi}$ decays with a branching ratio of 64 % via a $\beta$ emission towards $^{212}\text{Po}$ ($Q_\beta = 2.25$ MeV) which is again an $\alpha$ emitter with a short half-life of 300 ns. So, for these two chains a BiPo signature is an electron associated to a delayed $\alpha$ with a delay time depending on the isotope contamination we want to measure.

![Figure 1: BiPo processes for $^{214}\text{Bi}$ and $^{208}\text{Tl}$.](image)

The particles emitted by the source foil are detected with plastic scintillators coupled to low radioactivity photomultipliers (Fig. 2). Plastic scintillators are very radiopure and reduce the backscattering of electrons. The detection efficiency is dominated by the capacity for an $\alpha$ particle to escape the foil. GEANT4 simulations give a total efficiency of 6.5 % for contaminations.
in selenium foils (40 mg/cm²) with 1 MeV threshold for α. Therefore the energy threshold of the detector must be as low as possible. Moreover the energy converted into scintillation light is much lower for α compared to electrons. This quenching factor depends on the energy of the α and has been measured with a dedicated test bench. For example, a 1 MeV α will produce same amount of light than a 40 keV electron.

![Figure 2: BiPo detection principle with plastic scintillators and time signal seen with PMTs. Dots represent the contamination and crosses represent energy depositions in scintillators (trigger in blue and delayed in red).](image)

2 Backgrounds

The BiPo measurement consists in the detection of the electron in one scintillator and the detection of the delayed α particle in the other scintillator. This strong BiPo signature constrains the background of the detector to only 3 processes:

- bismuth (²¹²Bi or ²¹⁴Bi) contaminations in the volume of the scintillator. In such decay the electron deposits part of its energy in the first scintillator before crossing the foil to reach the other one. The delayed α is detectable only in the first scintillator because it can’t cross the foil. This background can be rejected because two hits in time are observed in the two scintillators: this is not a BiPo event. (Fig. 2a)

- bismuth contaminations on the surface of the scintillator. In this case the electron doesn’t deposit enough energy in the first scintillator to be detected. The delayed α particle is still detectable only in this first scintillator. This contamination is not distinguishable from a BiPo signal because this signature exactly corresponds to a BiPo event coming from the foil. (Fig. 2b)

- random coincidences due to external γ. To reduce this background, the BiPo detector uses low background materials, is shielded and installed in underground lab. The single counting rate of each scintillator has to be less than 40 mHz to measure ²⁰⁸Tl and less than 10 mHz for ²¹⁴Bi because of ²¹⁴Po longer half-life. Pulse shape discrimination also reduces this background. (Fig. 2c)

![Figure 3: Surface and volume bismuth contaminations of the scintillators and random coincidences backgrounds.](image)

3 BiPo1 prototype

BiPo1 prototype is divided in 20 modules. Each module is a black box containing two polystyren based scintillators coupled to low background 5" photomultipliers with PMMA light guides. Scintillator dimensions are 20×20×1 or 20×20×0.3 cm³, the entrance window is covered with 200 nm of ultra-pure aluminum to isolate optically each scintillator and to improve the light collection. The sides of scintillators and light guides are covered with 0.2 mm of Teflon for

*Small thickness where the energy deposited by the electron is below the threshold (~ 100 µm for 150 keV).*
light diffusion. The prototype is installed in the Modane Underground Laboratory (LSM) under 4800 m.w.e. Surrounding the modules, a shielding of 15 cm of low activity lead reduces external $\gamma$ and 3 cm of pure iron stops bremsstrahlung $\gamma$ emitted in the lead by the decay of the $^{210}\text{Bi}$ from long half-life $^{210}\text{Pb}$. Radon-free air flushes the volume of each module and the inner volume of the shielding (Fig. 4).

![Figure 4: BiPo1 prototype in its shielding in LSM.](image)

Photomultipliers signals are sampled with VME digitizing board, during 2.5 $\mu$s with a high sampling rate (1 GS/s) and a 12 bit high dynamic range (1 V). The acquisition is triggered each time a pulse reaches the 150 keV energy threshold and the 2 photomultipliers signals from a module are stored. The delayed hit research is performed later by the analysis of the signals (Fig. 5).

4 BiPo1 calibration

The first BiPo1 module has been dedicated to the validation of the detection principle for $^{212}\text{Bi}$ contaminations inside a foil. A 150 $\mu$m aluminum foil (40 mg/cm$^2$) with a contamination measured with HPGe detectors of $A(^{212}\text{Bi} \rightarrow ^{212}\text{Po}) = 0.19 \pm 0.03$ Bq/kg, has been installed between the two scintillators. After 141 days of data taking, 1501 BiPo events have been detected. Taking into account the efficiency calculated by GEANT4 simulations, it corresponds to a reconstructed activity of $A(^{212}\text{Bi} \rightarrow ^{212}\text{Po}) = 0.22 \pm 0.01$ Bq/kg, in good agreement with initial HPGe measurement. The delay between the two hits is also measured, and the fit of the decay law perfectly corresponds to the $^{212}\text{Po}$ half-life (Fig. 5). These results are a strong validation of the measurement principle and the calculated efficiency.

![Figure 5: Example of a BiPo event observed in BiPo1 and delay distribution between the $\beta$ and the $\alpha$ decays.](image)

Particle identification in the plastic scintillator could help to sign the BiPo process and reject random coincidences of $\gamma$. Indeed, longer half-life states in scintillators are excited by $\alpha$ particles.
but not by electrons. More light is therefore observable in the tail of the pulse for $\alpha$ particles. A pulse shape discrimination, using tail-to-total charge ratio as discrimination factor, has been applied on the photomultiplier signals from the data of the calibration foil. A good separation has been observed for prompt ($e^-$) and delayed ($\alpha$) signals (Fig. 6). Using this discrimination, it is possible to reject 80% of the random coincidence background and to keep 90% of the true BiPo events.

Figure 6: Distributions of the tail-to-total charge ratio and the tail charge $q$ as a function of the total charge $Q$ for the prompt and delayed signal from the aluminum foil in the first module of BiPo1.

5 BiPo1 surface radiopurity

The other modules of BiPo1 are dedicated to the measurement of the surface radiopurity of scintillators. BiPo events, coming from the contact surface between the two scintillators, are observed by a hit in one scintillator and a delayed hit in the other one. After 2 months of data taking, 7 BiPo events have been observed on a statistics equivalent to 0.8 m$^2$×month. It corresponds to an activity of about 2 $\mu$Bq/m$^2$. Extrapolating this background to a final 10 m$^2$ BiPo detector for a one month measurement of 4 kg selenium source foil for SuperNEMO, the thallium preliminary sensitivity is: $A^{({}^{208}\text{Tl})} < 7.5 \mu\text{Bq/kg (90% C.L.)}$.

6 Conclusion

The BiPo1 prototype demonstrated the validity of the experimental technique to measure $^{208}\text{Tl}$ contaminations in thin materials. The preliminary sensitivity achieved is 10 times better than standard HPGe measurements. Particles identification with simple plastic scintillators enhance the performances of the prototype. New modules using “phoswich” scintillators, a compound of a thin fast scintillator to detect $\alpha$ and a thick slow scintillator for electrons, will improve the discrimination between the 2 particles. A second prototype, BiPo2, using large scintillator plates (0.56 m$^2$) will also be tested in LSM before this summer.

References

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SEARCH FOR THE STANDARD MODEL HIGGS BOSON IN THE $Hz \rightarrow b\bar{b} \nu \bar{\nu}$ CHANNEL AT DØ

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A search for the standard model Higgs boson has been performed in 2.1 fb$^{-1}$ of $p\bar{p}$ collisions at 1.96 TeV, collected with the DØ detector at the Fermilab Tevatron. The final state considered is a pair of $b$ jets with large missing transverse energy, as expected from the reaction $p\bar{p} \rightarrow Hz \rightarrow b\bar{b} \nu \bar{\nu}$. The search is also sensitive to the $HW \rightarrow b\bar{b} \nu \bar{\nu}$ channel, when the charged lepton is not identified. Boosted decision trees were used to discriminate the signal from the backgrounds, dominated by $Wb\bar{b}$ and $Zb\bar{b}$ final states. For a Higgs boson mass of 115 GeV, a limit has been set at 95% C.L. on the cross section times branching fraction of $(p\bar{p} \rightarrow Hz(W/Z) \times (H \rightarrow b\bar{b}))$, which is 7.5 times larger than the standard model value.

1 Introduction

The $p\bar{p} \rightarrow Hz$ reaction, with $H \rightarrow b\bar{b}$ and $Z \rightarrow \nu \bar{\nu}$, is among the most promising for the discovery of a low mass Higgs boson at the Fermilab Tevatron. A search with 2.1fb$^{-1}$ of data collected with the DØ detector is presentented here. A lower mass limit of 114.4 GeV was set by the LEP experiments for the Higgs boson from analyses of the reaction $e^+ e^- \rightarrow Hz$, while an upper limit of 144 GeV can be inferred from precision electroweak data. Here and in the following, all limits quoted are at the 95% confidence level.

The final state topology considered in this analysis is a pair of $b$ jets from the decay of the Higgs boson, with missing transverse energy ($E_T$) due to the neutrinos from the $Z$ decay. The search is therefore also sensitive to the $HW$ channel, with $W \rightarrow \ell\nu$ when the charged lepton from the $W$ decay is not detected. The main backgrounds arise from $(W/Z)_j$+ jets, from top quark and diboson production and from multijet events produced by the strong interaction, with fake $E_T$ resulting from fluctuations in jet energy measurements and with real $b$ or mistagged light parton jets.
A kinematic selection is first applied to reject most of the multijet events. The two jets expected from the Higgs boson decay are next required to be tagged as $b$ jets, using a neural network $b$-tagging algorithm. Finally, discrimination between the signal and the remaining backgrounds is achieved by means of a boosted decision tree technique.

2 Data and Simulated samples

For this analysis, the data were recorded using a set of triggers designed to select events with jets and missing transverse energy. As the trigger conditions are not included into the full simulation, the calibration of the trigger response was performed on $Z \rightarrow \mu^+\mu^-+\text{jets}$ events. Due to the muons which only deposit energy in the calorimeter at the minimum of ionization, these events have the same topology as the signal. Except for the background from multijet production, which was estimated from data, all backgrounds from the standard model (SM) were determined by Monte Carlo simulation.

3 Event Selection

3.1 Pre-Tagging Selection

At this stage, most of the cuts are applied to specifically reject the multijet background. The selected events are required to have:

- exactly 2 or 3 jets with $p_T > 20$ GeV and within $|\eta_{\text{jet}}| < 2.5^a$
- $\Delta \phi(jet_1, jet_2) < 165^\circ$ to avoid back-to-back jets in the plane transverse to the beam direction.
- $E_T > 50$ GeV. This criterion is tightened in case the direction of the missing $E_T$ is close to the direction of one of the jets in the transverse plane: $E_T$(GeV) > 80 - 40 × min$\Delta \phi$($E_T$; any jet), where the angle is measured in radians.
- The asymmetry $A = (E_T - H_T)/(E_T + H_T)$, where $H_T = | - \Sigma \vec{p}_{T \text{jet}} |$, is required to lie between $-0.1$ and $0.2$.
- In signal events, the missing track-$p_T$, $\vec{p}_T$, defined as the opposite of the vectorial sum of the charged particle transverse momenta, is expected to point in a direction similar to that of the $E_T$. Advantage is taken of this feature by requiring $\Delta \phi(E_T, \vec{p}_T) < \pi/2$.

To reject backgrounds from $W+\text{jets}$, top, and diboson production, events containing an isolated electron or muon are rejected.

3.2 Heavy-Flavor Tagging

Advantage is next taken of the large branching fraction for $H \rightarrow t\bar{b}$ by requiring the two leading jets to be $b$-tagged.

As can be seen in Fig. 1 for the invariant mass of the dijet system built from the two leading jets, the simulation provides a good description of the data both before (40 340 events expected and observed) and after (439 events observed and 442.8±1.1 expected) $b$-tagging. After $b$-tagging, the $W+\text{jets}$ and $Z+\text{jets}$ backgrounds are dominated by heavy flavor jet production, which contributes $O(50\%)$.

4 Discriminant Analysis

In order to take full advantage of the different kinematic characteristics of the signal and background processes, a boosted decision tree (DT) technique was used. A dedicated DT was trained for each of the Higgs boson masses probed and for each of the two data taking periods (these are Run IIA and Run IIB with different in experimental conditions).

The DT discriminants are shown in Fig. 2 for a Higgs boson mass of 115 GeV.

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$a_{\text{jet}}$ is the pseudorapidity measured from the detector center.
Figure 1: Distributions of the invariant mass of the two leading jets before (left) and after $b$-tagging (right). The data are shown as points with error bars. The various background contributions (SM and multijet) are shown as histograms, with color codes as indicated in the frames. Distributions for a signal with a Higgs boson mass of 115 GeV are also shown, multiplied by 500 before and by 10 after $b$-tagging.

Figure 2: Distributions of the decision tree discriminants for a Higgs boson mass of 115 GeV ($HZ$ and $HW$ signals combined) for Run Ia (left) and Run Iib (right). The data are shown as points with error bars. The various background contributions (SM and multijet) are shown as histograms, with color codes as indicated in the frames. The distributions for the signal are multiplied by a factor of 25.

5 Systematic Uncertainties

Systematic uncertainties originate from various sources. Experimental uncertainties arise from the trigger simulation (5.5%), from the jet energy calibration (from 2% to 3%), resolution (about 1%), and reconstruction efficiency (2%), and from $b$-tagging (about 6%). Furthermore, a 6% error is assigned to the luminosity determination.

Moreover, the cross sections of the various SM and signal processes suffer from theoretical uncertainties. They were found to be at a level of 6 to 16%. They were estimated from mcfm $^4$ or from Refs. $^5$ and Ref. $^6$. Finally, uncertainties on the heavy flavor fractions of $W/Z+jets$ background(50%) are quoted.

6 Results

Agreement between data and expectation from SM and multijet backgrounds is observed both in terms of numbers of events selected and of DT discriminant shapes (Fig. 2). To set limits, based on the DT output, on the SM Higgs boson production cross section, a modified frequentist approach $^8$ was used.

The results obtained are shown as a function of the Higgs boson mass in Fig. 3 and in Table 1, in terms of the ratio of the excluded cross section times branching fraction for $H \to \bar{b}b$ to the SM Higgs prediction. The LLRs (Log Likelihood Ratios) are also shown in Fig. 3. For a 115 GeV Higgs boson mass, the observed and expected limits on the cross section of combined $HZ$ and $HW$ production times branching fraction for $H \to b\bar{b}$ are 7.5 and 8.4 times larger than the SM value, respectively.
Figure 3: As a function of the Higgs boson mass, limit on the cross section of combined $HZ$ and $HW$ production times branching fraction for $H \rightarrow b\bar{b}$ (left), relative to the SM value, and log likelihood ratio (right). On the right, the observed and expected limits are shown as solid and dashed curves, respectively. On the left, the observed LLRs are shown as black and red dashed curves for the background-only and signal+background hypotheses, respectively, and the green and yellow areas correspond to the one and two $\sigma$ deviations around the expected background-only LLR.

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Table 1: For various Higgs boson masses, observed and expected ratios of excluded to SM production cross sections times branching fraction for $H \rightarrow b\bar{b}$.

7 Summary

A search for the SM Higgs boson has been performed in 2.1 fb$^{-1}$ of $p\bar{p}$ collisions at 1.96 TeV. The topology analysed consists of a pair of $b$ jets with large $E_T$, as expected from $p\bar{p} \rightarrow HZ \rightarrow b\bar{b}a\bar{r}$. The search is also sensitive to $HW$ production, where the $W$ decays leptonically and the charged lepton is undetected. No deviation from the expectation from SM backgrounds was observed. A boosted decision tree technique was used to derive an upper limit on the cross section of the $p\bar{p} \rightarrow HZ$ and $p\bar{p} \rightarrow HW$ processes combined, as a function of the Higgs boson mass. For a mass of 115 GeV, this limit is a factor 7.5 larger than the SM cross section.

References

Search for SM Higgs in the $WH \to l\nu b\bar{b}$ Channel using $\sim 2\text{fb}^{-1}$

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We report a search for Standard Model (SM) Higgs boson production in association with a $W^\pm$ boson. This search uses data corresponding to an integrated luminosity of $1.9 \text{fb}^{-1}$ collected with the CDF detector at Tevatron. We select events matching the $W + \text{jets}$ signature and require at least one jets to be identified as $b$-quark jets. To further increase discrimination between signal and background, we use kinematic information in an artificial neural network. The number of tagged events and the resulting neural network output distributions are consistent with the Standard Model expectations, and we set an upper limit on the $WH$ production cross section times branching ratio $\sigma(p\bar{p} \to W^\pm H) \times BR(H \to b\bar{b}) < 1.1$ to 1.0 pb for Higgs masses from 110 GeV/c$^2$ to 150 GeV/c$^2$ at 95% confidence level.

1 Introduction

The success of the Standard Model in explaining and predicting experimental data provides strong motivation for the existence of a neutral Higgs boson. Current electroweak fits combined with direct searches from LEP2 indicate the mass of the Higgs boson is less than 190 GeV/c$^2$ at 95% confidence level $^1$. In proton-antiproton collisions of $\sqrt{s} = 1.96\text{TeV}$ at the Tevatron, the Standard Model Higgs boson may be produced in association with a $W$ boson $^2$. For low Higgs masses (below 140 GeV/c$^2$) the dominant decay mode is $H \to b\bar{b}$. The final state from the $WH$ production is therefore $l\nu b\bar{b}$, where the high-$p_T$ lepton from the $W$ decay provides an ideal trigger signature at CDF. The analysis strategies make use of $b$-tagging algorithm to suppress the $W + \text{jets}$ backgrounds and apply artificial neural network to discriminate signal to remaining backgrounds.
2 Event Selection

Events are collected by the CDF II detector with high-$p_T$ central electron or muon triggers which have an 18 GeV threshold. The central electron or muon is further required to be isolated with $E_T$ (or $p_T$) > 20 GeV in offline. Events having the $W$+jets signature are confirmed with a missing transverse energy requirement ($E_T > 20$ GeV).

We use forward (plug) electron events with a trigger intended for $W$ candidate events. This trigger requires both a forward electron candidate and missing transverse energy. Plug electron events are further required to have $E_T > 20$ GeV and $E_T > 25$ GeV. For plug electron events, additional selection is required to suppress QCD background.

Events are required to have 2 jets with $E_T > 20$ GeV and $|\eta| < 2.0$. Because the Higgs boson decays to $b\bar{b}$ pairs, we employ $b$-tagging algorithms which rely on the long lifetime and large mass of the $b$ quark to suppress enormous $W$+jets backgrounds.

2.1 Bottom Quark Tagging Algorithm

To greatly reduce the backgrounds to this Higgs search, we require that at least one jet in the event be identified as containing $b$-quarks by one of three $b$-tagging algorithms. The secondary vertex tagging algorithm identifies $b$ quarks by fitting tracks displaced from the primary vertex. In addition, we add jet probability tagging algorithm that identifies $b$ quarks by requiring a low probability that all tracks contained in a jet originated from the primary vertex, based on the track impact parameters. To be considered for double tag category, an event is required to have either two secondary vertex tags, or one secondary vertex tag and one jet probability tag.

Furthermore we also make use of exactly one $b$-tagged events with the secondary vertex tagging algorithm. To improve signal-to-background ratio for one tag events, we employ neural network (NN) $b$-tagging algorithm applied. This neural network is tuned for only jets tagged by the secondary vertex tagging algorithm. The purity of $b$-jets tagged by this algorithm is improved.

Finally we categorize events in three $b$-tagged conditions, double secondary vertex tagged events, one secondary vertex tagged plus one jet probability tagged events and one NN tagged events. These two categories of double-tagged events and one category of one neural network tagged events are defined exclusively.

2.2 Expected Signal Events and Systematics

The acceptance in Higgs mass of 120 GeV/$c^2$ in central region is $0.48\pm0.05\%$, $0.38\pm0.04\%$, $0.93\pm0.05\%$ for the double secondary vertex tagged, the secondary vertex plus jet probability and one neural network tagged category, respectively. The expected signal events in Higgs mass 120 GeV/$c^2$ are about 3.9 events in central plus plug data.

The uncertainties on the signal acceptance currently have the largest effect on the Higgs sensitivity. The $b$-tagging uncertainty is dominated by the uncertainty on the data/MC scale factor. 3.5-9.1% systematic is assigned for each $b$-tagging category. The uncertainties due to initial state radiation and final state radiation (2.9-5.2%) are estimated as difference from the nominal. Other uncertainties on parton distribution functions ($\sim 2\%$), jet energy scale (2-3%), trigger efficiencies ($< 1\%$) and lepton identification contribute (2%) are taken into account.

3 Backgrounds

This analysis builds on the method of background estimation detailed in Ref. 4. In particular, the contributions from the following individual backgrounds are calculated: falsely $b$-tagged events,
$W$ production with heavy flavor quark pairs, QCD events with false $W$ signatures, top quark pair production, and electroweak production (diboson, single top).

We estimate the number of falsely $b$-tagged events (mistags) by counting the number of negatively-tagged events, that is, events in which the measured displacement of the secondary vertex is opposite the $b$ jet direction. Such negative tags are due to tracking resolution limitations, but they provide a reasonable estimate of the number of false positive tags after a correction for material interactions and long-lived light flavor particles.

The number of events from $W +$ heavy flavor is calculated using information from both data and Monte Carlo programs. We calculate the fraction of $W$ events with associated heavy flavor production in the ALPGEN Monte Carlo program interfaced with the PYTHIA parton shower code. This fraction and the $b$-tagging efficiency for such events are applied to the number of events in the original $W+\text{jets}$ sample after correcting for the $t\bar{t}$ and electroweak contributions.

We constrain the number of QCD events with false $W$ signatures by assuming the lepton isolation is independent of $\not{E}_T$ and measuring the ratio of isolated to non-isolated leptons in a $\not{E}_T$ sideband region. The result in the tagged sample can be calculated in two ways: by applying the method directly to the tagged sample, or by estimating the number of non-$W$ QCD events in the pretag sample and applying an average $b$-tagging rate.

In this analysis, 83, 90 and 805 events for each $b$-tagging category are observed against $80.62 \pm 18.75$, $86.99 \pm 17.99$ and $809.61 \pm 159.38$ expected in signal region. The good agreement is also obtained in plug region.

4 Artificial Neural Network

To further improve signal to background separation we employ an artificial neural network. This neural network combines six kinematic variables into a single function with better discrimination between the Higgs signal and the background processes than any of the variables individually.

To train the neural network, JETNET package\textsuperscript{5}. The input variables are defined below:

**Dijet invariant mass**: The invariant mass reconstructed from the two jets. If there are additional looser jets, the loose jet that is closest to one of the two jets is included in this invariant mass calculation.

**Total System $p_T$**: The vector sum of the transverse momenta of the lepton, the $\not{E}_T$, and the two jets.

**$p_T$ Imbalance**: The scalar sum of the lepton and jet transverse momenta minus the $\not{E}_T$.

**$\sum E_T$ (loose jets)**: The scalar sum of the loose jet transverse energy.

**$M_{\text{min}}$**: The invariant mass of the lepton, $\not{E}_T$ and one of the two jets, where the jet is chosen to give the minimum invariant mass. The $p_z$ of neutrino is ignored for this quantity.

**$\Delta R$ (lepton-$\nu$)**: The distance between the direction of lepton and neutrino in $\eta - \phi$ plane, where the $p_z$ of the neutrino is taken from largest $|p_z|$ calculated from $W$ mass constraint.

The training is defined such that the neural network attempts to produce an output as close to 1.0 as possible for Higgs signal events and as close to 0.0 as possible for background events. For optimal sensitivity, a different neural network is trained for each Higgs mass considered. Fig.1 show the neural network results for each $b$-tagging category.

5 Results

We perform a direct search for an excess in the signal region of the neural network output distribution from single-tagged and double-tagged $W+\text{jets}$ events. A binned maximum likelihood technique is used to estimate upper limits on Higgs production by constraining the number
Figure 1: Neural network results calculated from six variables for each lepton region (Top: central lepton, Bottom: plug electron). The plot on the left shows exactly one $b$-tagged events, the plots on the center and right show double secondary vertex $b$-tagged event and one secondary vertex plus one jet probability $b$-tagged events, respectively.

of backgrounds to the estimates within uncertainties. Each $b$-tagging category is combined to obtain best sensitivity for both central and forward lepton region. We set an upper limit on the production cross section times branching ratio as a function of $m_H$, plotted in Fig. 2. The results are also collected in Table 3.

Figure 2: Observed and expected limits as a function of the Higgs mass hypothesis. The solid line shows observed 95% C.L. upper limit and dash line shows expected limit obtained by the assumption of null signal hypothesis.

References

STUDY OF VECTOR BOSON FUSION HIGGS IN ATLAS-LHC

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Within the framework of Standard Model, the production mode of Higgs boson through the fusion of the vector bosons $W$ or $Z$ (Vector Boson Fusion) is one of the most important production mechanisms, providing a specific detection signature. A detailed study regarding this issue is being undergone for ATLAS detector in LHC and some general features of this analysis are being presented in this note emphasizing in the study of Central Jet Veto.

1 VBF Topology and its characteristics

The four different production mechanisms of a Standard Model Higgs boson\(^1\) at the LHC are (Fig.1b): the $gg$ fusion, the Vector Boson Fusion, the Higgs boson production associated with vector bosons $W$ or $Z$ and the Higgs boson production associated with the production of a $t\bar{t}$ pair. Although VBF cross section is one order of magnitude lower than the dominant one of $gg$ fusion (Fig.1a), its topological characteristics provide a signature which makes it an important discovery channel for a low mass Higgs boson.

As it is illustrated in Fig.1b the VBF\(^a\) Higgs (or $WZ$ Fusion) occurs when two quarks originating from the proton beam are being scattered through the exchange of a $W$ or $Z$ boson and their fusion produces the Higgs boson. The important aspect of this topology is the fact the two scattered quarks give two energetic jets highly separated in rapidity in the forward regions of the detector.

Because of the electroweak character of the VBF Higgs boson production ($W, Z$ colorless exchange), a low gluon radiation activity is expected which is translated in detection terms as a low jet activity. Consequently, no jets are expected in the central region of the detector above a certain $p_T$ limit, in addition to the Higgs boson decay products. This aspect is motivating a central jet veto: the event is rejected if a third jet is detected in the central region above a given threshold of $p_T$. This cut is efficient against the main backgrounds and especially against $t\bar{t}$ and can be a useful tool to the improvement of the $S$ Vs $B$ separation.

\(^a\)Vector Boson Fusion
Figure 1: (a) Cross Section of Higgs boson for (b) the four different production mechanisms in LHC, within the frame of Standard Model.

2 VBF $H \rightarrow \tau^+\tau^-$ in ATLAS

The most promising part of the VBF Higgs analysis is the one which studies its decay into a pair of $\tau^+\tau^-$ leptons. This channel was first studied in ref.\textsuperscript{2} for ATLAS\textsuperscript{3} and now, a more detailed study is being undergone using state of the art Monte Carlo generators and a fully detailed simulation of the detector.

Due to the weak branching ratio of Higgs boson decaying to $\tau^+\tau^-$ for masses $m_h > 140\text{GeV}/c^2$, the analysis is performed in the mass range of $110\text{GeV}/c^2 < m_h < 140\text{GeV}/c^2$ and it is divided into three different categories depending on the decays of $\tau$ leptons (leptonic or hadronic decay). As a consequence, three sub-channels are obtained at the end: The lepton-lepton channel with a B.R. $\approx 12\%$, the lepton-hadron channel with a B.R. $\approx 46\%$ and the hadron-hadron channel with a B.R. $\approx 42\%$.

For the processes of background the generated events correspond to:

- $Z \rightarrow \tau^+\tau^- + jets$
- $t\bar{t} + jets$
- $W^\pm + jets$

Among them the dominant background is the $Z \rightarrow \tau^+\tau^- + jets$ which has a quite similar topology with the signal and therefore is the most difficult to suppress.

The final goal is to reconstruct the invariant mass of the couple $\tau^+\tau^-$ which equals to the mass of Higgs boson. The existence of at least two neutrinos in the event, requires the use of a collinear approximation to estimate the $m_{\tau^+\tau^-}$ invariant mass from the measure of the missing $p_T$ vector and the $p_T$ of the visible tau decay products. This approximation implies that the tau decay products are collinear to the direction of taus. This is an acceptable assumption considering the high $p_T$ of the $\tau$ leptons. However it provides a good mass resolution only if the tau decay products are not back to back, which is the case if the parent Higgs boson has reasonably high $p_T$.

The first step of the analysis consists of the identification of leptons and hadrons following the general ATLAS methods. Then, several cuts are applied in order to separate background from signal. This cuts can be categorized in three groups. The first group is related to the tau decay and different constraints concerning the kinematic variables of leptons or jets originating from...
the hadronic decay of tau are being applied. These cuts differ in lepton-lepton, lepton-hadron
and hadron-hadron cases. Next, the following set of cuts is based on collinear approximation
method used for the mass reconstruction and finally, several constraints deriving from the special
jet VBF behavior described in the previous paragraph are being applied such as, a given $p_T$
threshold for the two forward jets, a large separation in $\eta$, the decay products to lie between the
forward jets in $\eta$, high invariant mass of the two forward jets and also a central jet veto.

At Fig.2 we see the mass distribution after having applied all the different cuts. It is obvious
that the $Z \rightarrow \tau^+ \tau^-$ is the dominant background as expected, but still there is a clear peak related
to the mass of Higgs boson. Nevertheless, the importance of the mass resolution is highlighted
through this plot, which shows that a worse resolution would result to a contamination of the
signal by the $Z$ resonance.

3 Central Jet Veto

Jets play an important role in this channel and therefore, many studies have been done towards
this issue. In this analysis, the jets are reconstructed with a cone algorithm of a radius of 0.4
and part of these jet studies related to VBF Higgs topology, is the optimization of the central
jet veto cut. This cut rejects an event if a third jet is found in a region of $|\eta| < 3.2$ having a
$p_T > 20 GeV/c$. Other alternatives of this cut are also studied by changing this $|\eta|$ region. A
jet veto was studied without applying any constraints on $\eta$ and also applying the cut if a third
jet is only found between the two forward jets in $\eta$. Another interesting feature seen, was that
in many cases of signal events rejected by this cut, the vetoing jet was very close in $\Delta R$ to
one of the two forward jets. This was due to the splitting of the initial jet originating from
the scattered parton. Using a larger jet reconstruction cone, such as 0.7 for instance which is
the other alternative in ATLAS algorithms, the jet splitting effect might have been reduced but
the overall signal significance would still be lower. Therefore, a better option is to introduce
an additional constraint, which is to not take into consideration the third jet if it is close to
one of the forward jets, and more precisely, if $\Delta R$ between these two jets is less than 1. In
order to compare the three different approaches the cut efficiency is calculated for Signal and
Background. The results are shown in Fig.3.

With the combination of the $\Delta R$ condition to eliminate the vetoing jets close to the forward
jets and searching for a third jet between the two forwards jets in terms of $\eta$, a gain of $\sim 10\%$

$\Delta R = \sqrt{\Delta \eta^2 + \Delta \phi^2}$
Figure 3: Jet veto cut efficiency for signal and background. The white triangles represent the primary jet veto method where no $\Delta R$ cut is applied and the third jet is searched in $|\eta| < 3.2$. The green triangles represent the primary jet veto method but searching the third jet only between the two forward jets in term of $\eta$ and finally, the red squares represent the previous jet veto method having applied the $\Delta R$ cut as well. Every point in all curves corresponds to the efficiencies of a given $p_T$ threshold of the third jet starting from 10 GeV/c with a step of 5 GeV/c.

for the signal jet veto cut efficiency is achieved, as it is seen in Fig.3.

4 Outlook

The main characteristics of VBF Higgs boson topology and some general aspects of the analysis of its decay in a couple of $\tau^+\tau^-$ in ATLAS were presented. An example of improving a cut efficiency was addressed and it was shown that a gain of 10 % is accomplished modifying the jet veto cut. This analysis is not yet finalized since many studies are still being undergone, in particular concerning the effect of pile-up, which causes the creation of more jets into both signal and background events for the high LHC luminosities. Nevertheless, it was shown that VBF Higgs boson is one of the most promising channels for a low mass Higgs boson.

References

Yukawa corrections to Higgs production associated with two bottom quarks at the LHC

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We investigate the leading one-loop Yukawa corrections to the process $pp \rightarrow b\bar{b}H$ in the Standard Model. We find that the next-to-leading order correction to the cross section is small about $-4\%$ if the Higgs mass is 120GeV. However, the appearance of leading Landau singularity when $M_H \geq 2M_W$ can lead to a large correction at the next-to-next-to-leading order level for a Higgs mass around 160GeV.

1 Introduction

The cross section for $b\bar{b}H$ production at the LHC is very small compared to the gluon fusion channel. However, it is important to study that because of the following reasons:

- It can provide a direct measurement of the bottom-Higgs Yukawa coupling $(\lambda_{bbH})$ which can be strongly enhanced in the MSSM.

- We can identify the final state in experiment by tagging b-jets with high $p_T$. This reduces greatly the QCD background.

- Theoretically, it is a $2 \rightarrow 3$ process at the LHC which is a good example of one-loop multi-leg calculations. Moreover, the process $gg \rightarrow b\bar{b}H$ is, to the best of our knowledge, the most beautiful example where the leading Landau singularity (LLS) occurs in an electroweak box Feynman diagram. Considering that one rarely encounters such a singularity, studying its effect is very important.

The next-to-leading order (NLO) QCD correction to the exclusive process $pp \rightarrow b\bar{b}H$ with high $p_T$ bottom quarks has been calculated by two groups $^1$. The QCD correction is about $-22\%$ for $M_H = 120$GeV and $\mu = M_Z$ (renormalisation/factorisation scale). No leading Landau singularity occurs in any QCD one-loop diagrams.

The aim of our work is to calculate the Yukawa corrections, which are the leading electroweak corrections in this case, to the exclusive $bbH$ final state with high $p_T$ bottom quarks at the LHC $^2$. These corrections are triggered by top-charged Goldstone loops whereby, in effect, an external $b$ quark turns into a top quark. Such type of transitions can even trigger $gg \rightarrow b\bar{b}H$ even with vanishing $\lambda_{bbH}$, in which case the process is generated solely at one-loop level.

2 Calculation and results

At the LHC, the entirely dominant contribution comes from the sub-process $gg \rightarrow b\bar{b}H$. The contribution from the light quarks in the initial state is therefore neglected in our calculation.
Typical Feynman diagrams at the tree and one-loop levels are shown in Fig. 1. All the relevant

couplings are:

\[
\lambda_{bbH} = -\frac{m_b}{v}, \quad \lambda_{tH} = -\frac{m_t}{v}, \\
\lambda_{tb\chi} = -i\sqrt{2}\lambda_{tH}(P_L - \frac{m_b}{m_t}P_R), \quad \lambda_{\chi^-H} = \frac{M_H^2}{v},
\]

where \(v\) is the vacuum expectation value and \(P_{L,R} = (1 \mp \gamma_5)/2\). The cross section as a function of \(\lambda_{bbH}\) can be written in the form

\[
\sigma(\lambda_{bbH}) = \sigma(\lambda_{bbH} = 0) + \lambda_{bbH}^2\sigma'(\lambda_{bbH} = 0) + \cdots,
\]

\[
\lambda_{bbH}^2\sigma'(\lambda_{bbH} = 0) \approx \sigma_{NLO} = \sigma_{LO}[1 + \delta_{NLO}(m_t, M_H)],
\]

where \(\sigma(\lambda_{bbH} = 0)\) is shown in Fig. 2 (right), \(\sigma_{LO}\) and \(\sigma_{NLO}\) are shown in the same figure on the left.

\(\sigma(\lambda_{bbH} = 0)\) is generated solely at one-loop level and gets large when \(M_H\) is close to \(2M_W\). This is due to the leading Landau singularity related to the scalar loop integral associated to the box diagram in the class (c) of Fig. 1. This divergence, which occurs when \(M_H \geq 2M_W\), is not integrable at the level of loop amplitude squared and must be regulated by introducing a width for the unstable particles in the loops. Mathematically, the width effect is to move the LLS into the complex plane so that they do not occur in the physical region. The solution is shown in Fig. 3. The important point here is that the LLS, even after being regulated, can lead to a large correction to the cross section, up to 49% for \(M_H = 163\text{GeV}, \Gamma_W = 2.1\text{GeV}\) and \(\Gamma_t = 1.5\text{GeV}\).

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Figure 2: Left: the leading order (LO) and NLO cross sections as functions of $M_H$. Right: the cross section in the limit of vanishing $\lambda_{bbH}$.

Figure 3: The real and imaginary parts of the scalar box integral associated with the LLS diagram in the class (c) of Fig. 1.

References

The discovery potential for the MSSM with heavy scalars at the LHC in the case of light inos is examined. We discuss the phenomenology of the model and the observables to determine the parameters. We show that for light gauginos, the model parameters can be constrained with a precision of the order of 15%.

1 Introduction

Assuming a large soft–breaking scale for the MSSM scalars \(^{1,2,3,4,5}\) pushes squarks, sfermions and heavy Higgses out of the kinematic reach of the LHC without affecting the gaugino sector. The hierarchy problem will not be solved without an additional logarithmic fine tuning of the Higgs sector. Nevertheless, a model can be constructed to provide a good candidate for dark matter and realize grand unification while minimizing proton decay and FCNCs. We investigate the LHC phenomenology of the model, where all scalars are decoupled from the low energy spectrum. We focus on gaugino–related signatures to estimate the accuracy with which its underlying parameters can be determined.

2 Phenomenology

The spectrum at the LHC is reduced to the gauginos, Higgsinos and the light Higgs. At the intermediate scale \(M_S\) the effective theory is matched to the full theory and the usual MSSM renormalization group equations apply. The Higgsino mass parameter \(\mu\) and the ratio \(\tan \beta\) in the Higgs sector correspond to their MSSM counter parts. The gauginos masses \(M_{1,2,3}\) and the Higgs-sfermion-sfermion couplings unify, and \(M_S\) replaces the sfermion and the heavy Higgs’

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\(^a\)This work was done in collaboration with D. Zerwas (LAL-Orsay, Fr), N. Bernal and A. Djouadi (LPT-Orsay, Fr), M. Rauch and T. Plehn (U. of Edinburg, UK) and R. Lafaye (LAPP-Annecy, Fr)
mass parameters. This set resembles the mSUGRA parameter set except for $\tan \beta$ now playing the role of a matching parameter (with the heavy Higgses being decoupled) rather than that of an actual vev ratio.

We select our parameter point according to three constraints: first, we minimize the amount of fine tuning necessary to bring the light Higgs mass into the 100 to 200 GeV range and reduce $M_2$ to 10 TeV, still well outside the LHC mass range. The main reason for this low breaking scale is that we want the gluino to decay inside the detector (preferably at the interaction point) instead of being long–lived.

Secondly, we obtain the correct relic dark–matter density $\Omega h^2 = 0.111^{+0.006}_{-0.008}$ by setting $\mu = 290$ GeV and $M_2(M_{\text{GUT}}) = 132.4$ GeV or $M_2(M_{\text{weak}}) = 129$ GeV. This corresponds to the light–Higgs funnel $m_{\text{LSP}} \approx M_2/2 \approx M_h/2$, where the s-channel Higgs exchange enhances the LSP annihilation rate. And finally, $m_h$ needs to be well above the LEP limit, which we achieve by choosing $\tan \beta = 30$. We obtain $m_h = 129$ GeV, $m_\tilde{g} = 438$ GeV, chargino masses of 117 and 313 GeV, and neutralino masses of 60, 117, 296, and 310 GeV with a modified version of SuSpect, decoupling the heavy scalars from the MSSM RGEs. $\tilde{\chi}_0^2$ and $\tilde{\chi}_1^\pm$ as well as $\tilde{\chi}_0^0$ and $\tilde{\chi}_2^\pm$ are degenerate in mass. All supersymmetric particles and most notably the gluino are much lighter than in the SPS1a parameter point It is important to note that this feature is specific to our choice of parameters and not generic in heavy–scalar models. As a consequence, all LHC production cross sections are greatly enhanced with respect to SPS1a.

Table 1 shows the main (NLO) cross sections at the LHC from Prospino$^{12,13,14}$. The SUSY production is dominated by gluino pairs whose rate is eight times that of the SPS1a point: the lower gluino mass enlarges the available phase space, while in addition the destructive interference between s and t–channel diagrams is absent. The second largest process is the $\tilde{\chi}_1^\pm \tilde{\chi}_0^0$ production, which gives rise to a 145 fb of hard-jet free, $e$ and $\mu$ trilepton signal, more than a hundred times that of the SPS1a point.

### 3 OBSERVABLES

The first obvious observable is the light Higgs mass $m_h$. Although slightly higher than in most MSSM points, $m_h$ can still be measured in the Higgs to two photons decay$^{15}$ ($m_h < 150$ GeV). The systematic error on this measurement is mainly due to the uncertainty on the knowledge of the electromagnetic energy scale.

A measurement of the gluino pair production cross section appears feasible and could be very helpful to determine $M_3$. The branching ratio of gluinos decaying through a virtual squark into a chargino or a neutralino along with two jets is 85%. The chargino will in turn decay mostly into the LSP plus two leptons or jets. Such events would feature at least 4 high-$p_T$ jets, a large amount of missing energy due to the two $\tilde{\chi}_0^0$ in the final state and possibly leptons. The main backgrounds for such signatures are $t\bar{t}$ pairs, $W$+jets and $Z$+jets with respective production rates of 830 pb, 4640 pb and 220 pb. Despite these large cross sections, most of the background can be eliminated by applying standard cuts on $H_T$, the number of high-$p_T$ jets as well as the effective mass$^b$ which we checked using a fast LHC-like simulation. The main

\[ M_{\text{eff}} = E_T + \sum p_T(\text{jets}). \]
source of systematic errors for this observable is the 5% error on the knowledge of the luminosity. We take the theoretical error on the calculation of the cross section to be roughly 20%.

The next observable is the trilepton signal. After gluino pairs, the next dominant channel is the direct production of $\tilde{\chi}_1^+\tilde{\chi}_0^0$. 22% of $\tilde{\chi}_1^+$s decay through a virtual $W$ into an electron or muon and a neutrino and the LSP. Similarly, 7% of $\tilde{\chi}_0^0$s decay through a virtual $Z$ into an Opposite-Sign-Same-Flavour lepton pair (OSSF) and the LSP. The resulting signal features three leptons among which two are OSSF, a large amount of missing transverse energy due to the two LSPs plus the neutrino and no jet in the hard process. The background for this signature is mainly $WZ$ and $ZZ$ in which one of the leptons was non-identified or outside acceptance. According to PYTHIA the lepton production ($e$ and $\mu$) rates are 386 fb for $WZ$ and 73 fb for $ZZ$. The trilepton signal has a rate of 145 fb, using SDECAY for the calculation of the branching ratios. Including identification efficiencies of 65% for electrons and of 80% for muons gives rates of 110 to 211 fb for the background and 40 to 74 fb for the signal before any cut. A study with full detector simulation and reconstruction would provide a better understanding of signal and background. As in the previous case, the main source of systematic errors is the uncertainty on the luminosity. We also take the theoretical error on the value of the trilepton cross section to be roughly 20%.

Within this trilepton signal lies another observable. 10% of $\tilde{\chi}_0^0$s decay into an OSSF lepton pair and the LSP. The distribution of the invariant mass of the pair features a kinematic upper edge whose value is $m_{\tilde{\chi}_0^0} - m_{\tilde{\chi}_1^0}$. Such an observable gives precious information on the neutralino sector and hence on $M_1$. The systematic error is dominated by the lepton energy scale. The statistical error was extracted from a $\text{ROOT}$ fit of the $M_{\ell\ell}$ distribution and we estimate the theoretical accuracy to be of the order of 1%.

The last observable we use in this study is the ratio of gluino decays including a $b$ quark to those not including a $b$. A systematic error of 5% due to the tagging of $b$-jets and a theoretical uncertainties of 20% are assumed.

<table>
<thead>
<tr>
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<th>Statistical errors</th>
<th>Theoretical</th>
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</thead>
<tbody>
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<td>Error 0.1%</td>
<td>Source energy scale</td>
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<td>energy scale</td>
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<td>$R(\tilde{g} \rightarrow b/lb)$</td>
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<td>$b$-tagging</td>
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<tr>
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</table>

Table 2: Summary of the observables and the corresponding errors.

Table 2 summarises the value and error of the observables assumed in this study. The third and fourth columns give the experimental systematic errors and there source. The fifth column gives the statistical errors for an integrated luminosity of 100 fb$^{-1}$ corresponding to one year of data-taking at the LHC nominal luminosity. The last column gives an estimation of the theoretical uncertainties.

4 PARAMETER DETERMINATION

We use different sets of errors for the fits. First we determine the parameters in the low statistic scenario ignoring theoretical uncertainties. Second we assume an infinite statistic and therefore assume negligible statistical errors to estimate the ultimate precision barrier imposed by experimental systematic errors. Finally the effect of theoretical uncertainties is estimated by including them into the previous set. We expect these to dominate.

With no information on the squark and sfermion sector at all, except for non-observation, we are
forced to fix $M_S$ and $A_t$ and set $M_2$ to be equal to $M_1$. We fit the parameters to the observables using the Minuit fitter. The minimum of the $\chi^2$ is found by MIGRAD. We start from a point far from the nominal values ($\{M_1, M_3, \tan \beta, \mu\} = \{100, 200, 10, 320\}$) and reach the values reported in table 3. Errors are determined with MINOS. Theoretical errors are treated as Gaussian.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Nom. values</th>
<th>Fit values</th>
<th>Low stat.</th>
<th>$\infty$ stat.</th>
<th>$\infty$ stat.+th</th>
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<tr>
<td>$M_S$</td>
<td>10 TeV</td>
<td>fixed</td>
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<td></td>
<td></td>
</tr>
<tr>
<td>$A_t$</td>
<td>0</td>
<td>fixed</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$M_1$</td>
<td>132.4 GeV</td>
<td>132.8 GeV</td>
<td>6</td>
<td>0.24</td>
<td>0.2%</td>
</tr>
<tr>
<td>$M_2$</td>
<td>132.4 GeV</td>
<td>132.8 GeV</td>
<td>6</td>
<td>0.24</td>
<td>0.2%</td>
</tr>
<tr>
<td>$M_3$</td>
<td>132.4 GeV</td>
<td>132.7 GeV</td>
<td>60</td>
<td>1.24</td>
<td>4%</td>
</tr>
<tr>
<td>$\tan \beta$</td>
<td>30</td>
<td>28.3</td>
<td>60</td>
<td>1.24</td>
<td>4%</td>
</tr>
<tr>
<td>$\mu$</td>
<td>290 GeV</td>
<td>288 GeV</td>
<td>3.8</td>
<td>1.1</td>
<td>0.4%</td>
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</tbody>
</table>

Table 3: Result of the fits. Errors on the determination of the parameter are given for the three error sets. Both absolute and relative values are given.

Table 3 shows the result of the fits in both absolute and relative values. It is interesting to note that $\tan \beta$ in undetermined except in the case of infinite statistical and theoretical accuracies. The quality of the trilepton and gluino signals gives very good precision on the determination of $M_1$ and $M_3$ even with low statistics. The inclusion of theoretical uncertainties indeed decreases the accuracy but still allows for a determination. $M_3$ only depends on the large gluino signal and its decays, explaining its relative stability. $M_1$ and $M_2$ see the largest impact of theoretical errors. This is because they depend on first order on the trilepton cross-section and on second order on the $b$ to non $b$ gluino decays ratio both of which bear a large theoretical error.

5 CONCLUSION

The MSSM with heavy scalars can very well satisfy current experimental and theoretical limits on physics beyond the standard model and also solve a good number of issues present in the traditional MSSM. We described its phenomenology at the LHC in the case of light inos and showed that such a simple and light spectrum could lead to very high production rates making the model discoverable. The main observable channels are gluino pairs and the trilepton channel whose hard-jet free channel makes it well distinct from SM and SUSY backgrounds. Other observables such as the light Higgs mass, the $|m_{\tilde{\chi}^0_2} - m_{\tilde{\chi}^0_1}|$ kinematic edge and the $b$ to non $b$ producing gluino decays could lead to a determination of most parameters to the level of a few percent with 100 fb$^{-1}$ ignoring theoretical errors. In a more realistic picture where we assumed non-zero theoretical errors, we saw that most parameters can be determined with a precision of 15%. We also saw that the scalar section including $\tan \beta$ could only be poorly determined if at all.

New complementary observables could help determine better the scalar sector. Equally, a look at other parameter points will provide a more complete view of the discovery potential of a MSSM with decoupled scalars at the LHC.

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Search for the decay $K_S \to e^+ e^-$

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We present results of a direct search for the decay $K_S \to e^+ e^-$ with the KLOE detector, obtained with a sample of $e^+ e^- \to \phi \to K_SK_L$ events produced at DAΦNE, the Frascati $\phi$–factory, for an integrated luminosity of 1.9 fb$^{-1}$. The Standard Model prediction for this decay is BR($K_S \to e^+ e^-$) = $2 \times 10^{-14}$. The search has been performed by tagging the $K_S$ decays with simultaneous detection of a $K_L$ interaction in the calorimeter. Background rejection has been optimized by using both kinematic cuts and particle identification. At the end of the analysis chain we find BR($K_S \to e^+ e^-$) < $9.3 \times 10^{-9}$ at 90% CL, which improves by a factor of $\sim 15$ on the previous best result, obtained by CPLEAR experiment.

1 Introduction

The decay $K_S \to e^+ e^-$, like the decay $K_L \to e^+ e^-$ or $K_L \to \mu^+ \mu^-$, is a flavour-changing neutral-current process, suppressed in the Standard Model and dominated by the two-photon intermediate state. For both $K_S$ and $K_L$, the $e^+ e^-$ channel is much more suppressed than the $\mu^+ \mu^-$ one (by a factor of $\sim 250$) because of the $e^- - \mu$ mass difference. The diagram corresponding to the process $K_S \to \gamma^* \gamma^* \to \ell^+ \ell^-$ is shown in Fig. 1. Using Chiral Perturbation Theory ($\chi$PT) to order $O(p^4)$, the Standard Model prediction BR($K_S \to e^+ e^-$) is evaluated to be $\sim 2 \times 10^{-14}$. A value significantly higher than expected would point to new physics. The best experimental limit for $BR(K_S \to e^+ e^-)$ has been measured by CPLEAR, and it is equal to $1.4 \times 10^{-7}$, at 90% CL. Here we present a new measurement of this channel, which improves on the previous result by a factor of $\sim 15$.

Figure 1: Long distance contribution to $K_S \rightarrow \ell^+ \ell^-$ process, mediated by two-photon rescattering.
2 Experimental setup

The data were collected with KLOE detector at DAΦNE, the Frascati φ-factory. DAΦNE is an $e^+e^-$ collider that operates at a center-of-mass energy of $\sim 1020$ MeV, the mass of the φ meson. φ mesons decay $\sim 34\%$ of the time into nearly collinear $K^0\bar{K}^0$ pairs. Because $J^{PC}(\phi) = 1^{--}$, the kaon pair is in an antisymmetric state, so that the final state is always $K_S K_L$. Therefore, the detection of a $K_L$ signals the presence of a $K_S$ of known momentum and direction, independently of its decay mode. This technique is called $K_S$ tagging. A total of $\sim 4$ billion φ were produced, yielding $\sim 1.4$ billion of $K_S K_L$ pairs.

The KLOE detector consists of a large cylindrical drift chamber (DC), surrounded by a lead/scintillating-fiber sampling calorimeter (EMC). A superconducting coil surrounding the calorimeter provides a 0.52 T magnetic field. The drift chamber $^3$, 4 m in diameter and 3.3 m long, is made of carbon-fibers/epoxy and filled with a light gas mixture, 90$\%$ He-10$\%$C$_4$H$_{10}$. The DC position resolutions are $\sigma_{xy} \approx 150\mu$m and $\sigma_z \approx 2$ mm. DC momentum resolution is $\sigma(p_{1\perp})/p_{1\perp} \approx 0.4\%$. Vertices are reconstructed with a spatial resolution of $\sim 3\,$mm.

The calorimeter $^4$ is divided into a barrel and two endcaps and covers 98$\%$ of the solid angle. The energy and time resolutions are $\sigma_E/E = 5.7\%/\sqrt{E(\text{GeV})}$ and $\sigma_t = 57 \text{ ps}/\sqrt{E(\text{GeV})} \pm 100 \text{ ps}$, respectively.

To study the background rejection, a MC sample of φ decays to all possible final states has been used, for an integrated luminosity of $\sim 1.9\,$fb$^{-1}$. A MC sample of $\sim 45000$ signal events has been also produced, to measure the analysis efficiency.

3 Data analysis

The identification of $K_L$-interaction in the EMC is used to tag the presence of $K_S$ mesons. The mean decay lengths of $K_S$ and $K_L$ are $\lambda_S \sim 0.6\,$cm and $\lambda_L \sim 350\,$cm, respectively. About 50$\%$ of $K_L$'s therefore reach the calorimeter before decaying. The $K_L$ interaction in the calorimeter barrel ($K_{\text{crash}}$) is identified by requiring a cluster of energy greater than 125 MeV not associated with any track, and whose time corresponds to a velocity $\beta = v_{cl}/c \sigma_{cl}$ compatible with the kaon velocity in the φ center of mass, $\beta^* \sim 0.216$, after the residual φ motion is considered. Cutting at $0.17 \leq \beta^* \leq 0.28$ we selected $\sim 650$ million $K_S$-tagged events ($K_{\text{crash}}$ events in the following), which are used as a starting sample for the $K_S \rightarrow e^+e^-$ search.

$K_S \rightarrow e^+e^-$ events are selected by requiring the presence of two tracks of opposite charge with their point of closest approach to the origin inside a cylinder 4 cm in radius and 10 cm in length along the beam line. The track momenta and polar angles must satisfy the fiducial cuts $120 \leq p \leq 350\,$MeV and $30^\circ \leq \theta \leq 150^\circ$. The tracks must also reach the EMC without spiralling, and have an associated cluster. In Fig. 2, the two-track invariant mass evaluated in electron hypothesis ($M_{ee}$) is shown for both MC signal and background samples. A preselection cut requiring $M_{ee} > 420\,$MeV has been applied, which rejects most of $K_S \rightarrow \pi^+\pi^-$ events, for which $M_{ee} \sim 400\,$MeV. The residual background has two main components: $K_S \rightarrow \pi^+\pi^-$ events, populating the low $M_{ee}$ region, and $\phi \rightarrow \pi^+\pi^-\pi^0$ events, spreading over the whole spectrum. The $K_S \rightarrow \pi^+\pi^-$ events have such a wrong reconstructed $M_{ee}$ because of track resolution or one pion decaying into a muon. The $\phi \rightarrow \pi^+\pi^-\pi^0$ events enter the preselection because of a machine background cluster, accidentally satisfying the $K_{\text{crash}}$ algorithm. After preselection we are left with $\sim 5 \times 10^5$ events. To have a better separation between signal and background, a $\chi^2$-like variable is defined, collecting informations from the clusters associated to the candidate electron tracks. Using the MC signal events we built likelihood functions based on: the sum and the difference of $\delta t$ for the two tracks, where $\delta t = t_{cl} - L/3c$ is evaluated in electron hypothesis; the ratio $E/p$ between the cluster energy and the track momentum, for both charges; the
cluster depth, evaluated respect to the track, for both charges. In Fig. 2, the scatter plot of $\chi^2$ versus $M_{ee}$ is shown, for MC signal and background sources. The $\chi^2$ spectrum for background is concentrated at higher values respect to signal, since both $K_S \rightarrow \pi^+\pi^-$ and $\phi \rightarrow \pi^+\pi^-\pi^0$ events have pions in the final state.

A signal box to select the $K_S \rightarrow e^+e^-$ events can be conveniently defined in the $M_{ee} - \chi^2$ plane (see Fig. 2); nevertheless we investigated some more independent requirements in order to reduce the background contamination as much as possible before applying the $M_{ee} - \chi^2$ selection.

Charged pions from $K_S \rightarrow \pi^+\pi^-$ decay have a momentum in the $K_S$ rest frame $p_\pi^* \sim 206$ MeV. The distribution of track momenta in the $K_S$ rest frame, evaluated in the pion mass hypothesis, is shown in Fig. 2, for MC background and MC signal. For most of $K_S \rightarrow \pi^+\pi^-$ decays, at least one pion has well reconstructed momentum, so that the requirements

$$\min(p_\pi^*(1), p_\pi^*(2)) \geq 220 \text{ MeV} \quad , \quad p_\pi^*(1) + p_\pi^*(2) \geq 478 \text{ MeV}$$

rejects $\sim 99.9\%$ of these events, while retaining $\sim 92\%$ of the signal.

To reject $\phi \rightarrow \pi^+\pi^-\pi^0$ events we have applied a cut on the missing momentum, defined as:

$$P_{\text{miss}} = \left| \vec{P}_\phi - \vec{P}_L - \vec{P}_S \right|$$

(2)

where $\vec{P}_{L,S}$ are the neutral kaon momenta, and $\vec{P}_\phi$ is the $\phi$ momentum. The distribution of $P_{\text{miss}}$ is shown in Fig. 2, for MC background and for MC signal events. We require

$$P_{\text{miss}} \leq 40 \text{ MeV} ,$$

(3)

which rejects almost completely the $3\pi$ background source which is distributed at high missing momentum.

A signal box is defined in the $M_{ee} - \chi^2$ plane as shown Fig. 2. The $\chi^2$ cut for the signal box definition has been chosen to remove all MC background events: $\chi^2 < 70$. The cut on $M_{ee}$ is practically set by the $p_\pi^*$ cut, which rules out all signal events with a radiated photon with energy greater than 20 MeV, corresponding to an invariant mass window: $477 < M_{ee} \leq 510$ MeV. The signal box selection on data gives $N_{\text{obs}} = 0$. The upper limit at 90% CL on the expected number of signal events is $UL(\mu_S) = 2.3$.

4 Results

The total selection efficiency on $K_S \rightarrow e^+e^-$ events is evaluated by MC, using the following parametrization:

$$\epsilon_{\text{sig}} = \epsilon(K_{\text{crash}}) \times \epsilon(\text{sel}|K_{\text{crash}}) ,$$

(4)

where $\epsilon(K_{\text{crash}})$ is the tagging efficiency, and $\epsilon(\text{sel}|K_{\text{crash}})$ is the signal selection efficiency on the sample of tagged events. The efficiency evaluation includes contribution from radiative corrections. The number of $K_S \rightarrow \pi^+\pi^-$ events $N_{\pi^+\pi^-}$ counted on the same sample of $K_S$ tagged events is used as normalization, with a similar expression for the efficiency. The upper limit on BR($K_S \rightarrow e^+e^-$) is evaluated as follows:

$$UL(BR(K_S \rightarrow e^+e^-)) = UL(\mu_s) \times R_{\text{tag}} \times \frac{\epsilon_{\pi^+\pi^-}(\text{sel}|K_{\text{crash}})}{\epsilon_{\text{sig}}(\text{sel}|K_{\text{crash}})} \times \frac{BR(K_S \rightarrow \pi^+\pi^-)}{N_{\pi^+\pi^-}} ,$$

(5)

where $R_{\text{tag}}$ is the tagging efficiency ratio, corresponding to a small correction due to the $K_{\text{crash}}$ algorithm dependence on $K_S$ decay mode, and it is equal to 0.9634(1). Using $\epsilon_{\text{sig}}(\text{sel}|K_{\text{crash}}) = 0.465(4)$, $\epsilon_{\pi^+\pi^-}(\text{sel}|K_{\text{crash}}) = 0.6102(5)$ and $N_{\pi^+\pi^-} = 217, 422, 768$, we obtain

$$UL(BR(K_S \rightarrow e^+e^-(\gamma))) = 9.3 \times 10^{-9} , \text{ at 90\% CL} .$$

(6)
Our measurement improves by a factor of $\sim 15$ on the CPLEAR result $^2$, for the first time including radiative corrections in the evaluation of the upper limit.

References

CP ASYMMETRIES FROM NON-UNITARY LEPTONIC MIXING

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A low-energy non-unitary leptonic mixing matrix is a generic effect of some theories of new physics accounting for neutrino masses. We show how the extra CP-phases of a general non-unitary matrix allow for sizeable CP-asymmetries in the $\nu_\mu \to \nu_\tau$ channel. This CP-asymmetries turns out to be an excellent probe of such new physics.

1 Introduction

Non-zero neutrino masses constitutes one of the main evidence for physics beyond the Standard Model of particle physics (SM). In the complete theory accounting for them and encompassing the SM the complete mixing matrices should be unitary, as mandated by probability conservation. However, the effective $3 \times 3$ submatrices describing the mixing of the light fermionic fields need not to be unitary, since these known fields in the theory may mix with other degrees of freedom. We will assume that the full theory is indeed unitary, whereas, low-energy non-unitarity may result from BSM physics contributing to neutrino propagation, when the physical measurements are described solely in terms of SM fields\textsuperscript{1}. Specifically, the tree-level exchange of heavy fermions (scalars) will (not) induce low-energy non-unitary contributions through dimension six effective operators\textsuperscript{2}.

In Ref.\textsuperscript{1} the so-called MUV (minimal unitarity violation) was developed and the absolute values of the elements of the matrix $N$ were determined. However, no information on the phases of the mixing matrix is available, neither on the standard phases nor on the new non-unitary ones, as present oscillation data correspond mainly to disappearance experiments. Here we will explore the future sensitivity to the new CP-odd phases of the leptonic mixing matrix associated to non-unitary. In particular, it will be shown that CP-asymmetries in the $\nu_\mu \to \nu_\tau$ channel are an excellent probe of such new physics. Notice that CP-odd effects in that channel are negligible in the standard unitary case, in which the golden channel for CP-violation is $\nu_e \to \nu_\mu$.

2 Formalism

We start from the parametrization of the general non-unitary matrix $N$, which relates flavor and mass fields, as the product of an hermitian and a unitary matrix, defined by

$$\nu_\alpha = N_{\alpha \xi} \nu_\xi \equiv [(1 + \eta)U]_{\alpha \xi} \nu_\xi ,$$

(1)
with $\eta^3 = \eta$. Since $NN^\dagger = (1 + \eta)^2 \approx 1 + 2\eta$, the bounds derived in Ref.\cite{1} for the modulus of the elements of $NN^\dagger$ can also be translated into constraints on the elements of $\eta$. It follows that

$$|\eta| = \begin{pmatrix} < 5.5 \cdot 10^{-3} & < 3.5 \cdot 10^{-5} & < 8.0 \cdot 10^{-3} \\ < 3.5 \cdot 10^{-5} & < 5.0 \cdot 10^{-3} & < 5.1 \cdot 10^{-3} \\ < 8.0 \cdot 10^{-3} & < 5.1 \cdot 10^{-3} & < 5.0 \cdot 10^{-3} \end{pmatrix}, \quad (2)$$

at the 90% confidence level. The bounds have been updated with the latest experimental bound on $\tau \to \mu \gamma$\cite{3}. Eq. (2) shows that the matrix $N$ is constrained to be unitary at the per cent level accuracy or better. Therefore, the unitary matrix $U$ in Eq. (1) can be identified with the usual unitary mixing matrix $U_{PMNS}$, within the same accuracy. The flavor eigenstates can then be conveniently expressed as

$$|\nu_\alpha\rangle = \frac{(1 + \eta^*)_{\alpha\beta}U_{\beta i}^*}{|1 + 2\eta_{\alpha\alpha} + (\eta^2)_{\alpha\alpha}|^{1/2}} |\nu_i\rangle,$$

being the usual oscillation amplitude of the unitary analysis.

In order to study the possible new CP-violation signals arising from the new phases in $\eta$, one has to consider appearance channels, $\alpha \neq \beta$. The best sensitivities to such phases will be achieved in a regime where the first term in Eq. (4) is suppressed. This is possible at short enough baselines, where the standard appearance amplitudes are very small while $A^{SM}_{\alpha\alpha}(L) \simeq 1$. In that case the total amplitude is well approached by

$$< \nu_\beta | \nu_\alpha(L) > = A^{SM}_{\alpha\beta}(L) \left( 1 - \eta_{\alpha\alpha} - \eta_{\beta\beta} \right) + \sum_\gamma (\eta^*_{\alpha\gamma}A^{SM}_{\gamma\beta}(L) + \eta_{\beta\gamma}A^{SM}_{\alpha\gamma}(L)),$$

with

$$A^{SM}_{\alpha\beta}(L) \equiv < \nu_\beta^{SM} | \nu_\alpha^{SM}(L) > \quad (5)$$

being the usual oscillation amplitude of the unitary analysis.

For instance, within the above-described approximation, the oscillation probability for two families would read:

$$P_{\alpha\beta} = \sin^2(2\theta)\sin^2\left(\frac{\Delta L}{2}\right) - 4|\eta_{\alpha\beta}| \sin \delta_{\alpha\beta} \sin(2\theta) \sin(\Delta L) + 4|\eta_{\alpha\beta}|^2,$$

where $\Delta = \Delta m^2/2E$ and $\eta_{\alpha\beta} = |\eta_{\alpha\beta}|e^{-i\delta_{\alpha\beta}}$. The first term in the above equation is the standard oscillation probability. The third term is associated to the zero-distance effect stemming from the non-orthogonality of the flavor eigenstates\cite{1}. Finally, the second term is the CP-violating interference between the other two which we are interested in. Notice that, even in two families, there is CP-violation due to the non-unitarity.

\textsuperscript{a}The handy superscript $SM$ is an abuse of language, to describe the flavor eigenstates of the standard unitary analysis.
3 Sensitivity to the new CP-odd phases: the $\nu_\mu \to \nu_\tau$ channel

This channel is the best option to study the CP-violation coming from non-unitarity. This is because present constraints on $\eta_{e\mu}$ are too strong to allow a signal in the $\nu_e \to \nu_\mu$ one (see Eq. (2)), and $\nu_e \to \nu_\tau$ has extra suppressions by small standard parameters such as $\sin \theta_{13}$ or $\Delta_{12}^4$. In the numerical computation of $P_{\mu\tau}$, the only approximation performed is to neglect all $\eta$ elements but $\eta_{\mu\tau}$. They should be indeed subdominant (see Eq. (6)). This approximation has been checked numerically.

Eq. (7) suggests us that the best sensitivities to CP-violation will be achieved at short baselines and high energies, where the standard term is suppressed by $\sin^2 (\frac{\Delta L}{2})$. Therefore, we will study a Neutrino Factory beam resulting from the decay of 50 GeV muons, to be detected at a 130 Km baseline, which matches for example the CERN-Frejus distance. For this set-up, $\sin(\frac{\Delta m^2_{31} L}{2}) \simeq 1.7 \cdot 10^{-2}$ and $\sin(\frac{\Delta m^2_{21} L}{2}) \simeq 6 \cdot 10^{-4}$, where $\Delta_{jk} \equiv (m_j^2 - m_k^2)/2E$. Then, if the $\eta_{\mu\tau}$ value is close to their experimental limit in Eq. (2), all terms in the oscillation probability given in Eq. (7) can be of similar order for the channel $\nu_\mu \to \nu_\tau$. On the experimental side, we will assume $2 \cdot 10^{20}$ useful decays per year and five years running with each polarity, consider a 5 Kt Opera-like detector and, finally, sensitivities and backgrounds a factor 5 larger than those used for the $\nu_e \to \nu_\tau$ channel in Ref. 6.

Since this channel is not suppressed by small standard parameters such as $\sin \theta_{13}$ or $\Delta_{12}$, the two family approximation in Eq. (7), with $\theta = \theta_{23}$ and $\Delta = \Delta_{31}$, is very accurate to understand the results. This equation indicates that the CP-odd interference term is only suppressed linearly in $|\eta_{\mu\tau}|$. This can indeed be observed in the result of the complete numerical computation, Fig. 1, which shows the sensitivities to $|\eta_{\mu\tau}|$ and $\delta_{\mu\tau}$ obtained. The left panel represents two fits to two different input values of $|\eta_{\mu\tau}|$ and $\delta_{\mu\tau}$ (depicted by stars). The dashed lines correspond to fits done assuming the wrong hierarchy, that is the opposite sign for $\Delta_{31}$ to that with which the number of events were generated. As expected from Eq. (7), a change of sign for the mass difference can be traded by a change of sign for $\delta_{\mu\tau}$. Nevertheless, this does not spoil the potential for the discovery of CP violation, since a non-trivial value for $|\delta_{\mu\tau}|$ is enough.

\textbf{Footnote:} The complete expanded expression for $P_{\mu\tau}$ can be found in Ref. 4.
to indicate CP violation. Furthermore, the sinusoidal dependence implies as well a degeneracy
between $\delta_{\mu\tau} \to 180^\circ - \delta_{\mu\tau}$, as reflected in the figure.

The right panel in Fig. 1 depicts the $3\sigma$ sensitivities to $|\eta_{\mu\tau}|$ (solid line) and $\delta_{\mu\tau}$ (dotted
line), while the present bound from $\tau \to \mu\gamma$ is also shown (dashed line). The poorest sensitivity
to $|\eta_{\mu\tau}|$, around $10^{-3}$, is found in the vicinity of $\delta_{\mu\tau} = 0$ and $\delta_{\mu\tau} = 180^\circ$, where the CP-
odd interference term vanishes and the bound is placed through the subleading $|\eta_{\mu\tau}|^2$ term. The
latter is also present at zero distance and its effects were already considered in Ref. 1, obtaining a
bound of similar magnitude. The sensitivity to $|\eta_{\mu\tau}|$ peaks around $|\eta_{\mu\tau}| \approx 4 \cdot 10^{-4}$ for $\delta_{\mu\tau} \simeq \pm 90^\circ$, 
where $\sin \delta_{\mu\tau}$ is maximum. That is, for non-trivial values of $\delta_{\mu\tau}$ not only CP-violation could be
discovered, but values of $|\eta_{\mu\tau}|$ an order of magnitude smaller could be probed.

4 Conclusions

The flavour mixing matrix present in leptonic weak currents may be generically non-unitary,
as a result of new physics responsible for neutrino masses. Even if the effects are expected to
be extremely tiny in the simplest models of neutrino masses, it is important to analyze the
low-energy data without assuming a unitary leptonic mixing matrix, since it is a generic window
of new physics.

A non-unitary matrix has more moduli and phases than a unitary one. The last ones may
lead to new signals of CP-violation. In particular, we have shown that an asymmetry between
the strength of $\nu_\mu \to \nu_\tau$ oscillations versus that for $\bar{\nu}_\mu \to \bar{\nu}_\tau$ results to be a beautiful and excellent
probe of new physics, considering short-baselines ($\sim 100$ km.) and a Neutrino Factory beam of
energy $\mathcal{O}(20 \text{GeV})$. Non-trivial values of the new phases can also allow to probe the absolute
value of the moduli down to $10^{-4}$, an improvement of an order of magnitude over previous
analyzes of future facilities.

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References

DETERMINATION OF THE $B_s$ LIFETIME USING HADRONIC DECAYS

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We present a measurement of the $B_s^0$ meson lifetime using fully and partially reconstructed hadronic decays $B_s^0 \rightarrow D_s^- \pi^+(X)$ followed by $D_s^- \rightarrow \phi \pi^-$. The data sample was recorded with the CDF II detector at the Fermilab Tevatron and corresponds to an integrated luminosity of $1.3 \text{ fb}^{-1}$ from $p\bar{p}$ collisions at $\sqrt{s} = 1.96 \text{ TeV}$.

1 Introduction

For hadrons containing the heavy $b$ quark, one might assume that the light quark is only a spectator to the decay of the $b$, and the lifetimes of all the $b$ hadrons are the same regardless of the light quark’s flavor. However, the spectator quarks do participate in the decay, and the hierarchy $\tau(B_c) < \tau(\Lambda_b) \approx \tau(B_d) < \tau(B_u)$ has been both theoretically predicted and experimentally observed. The ratio of lifetimes is often quoted so some systematic uncertainties cancel. Recent theoretical calculations predict $\frac{\tau(B_u)}{\tau(B_d)} = 1.06 \pm 0.02$, $\frac{\tau(B_s)}{\tau(B_d)} = 1.00 \pm 0.01$, and $\frac{\tau(\Lambda_b)}{\tau(B_d)} = 0.86 \pm 0.05$. The world averages for the corresponding experimental numbers are $1.071 \pm 0.009$, $0.939 \pm 0.021$, and $0.921 \pm 0.036$, respectively. The experimental uncertainties are smaller than the theoretical uncertainties for all but the $B_s$ ratio. As $\tau(B_d)$ is already well measured, further improvements must come from reducing the uncertainty on $\tau(B_s)$.

In this Proceeding we present the most precise flavor-specific measurement of the $B_s$ lifetime to date. The data come from $p\bar{p}$ collisions at $\sqrt{s} = 1.96 \text{ TeV}$ produced at the Fermilab Tevatron. This analysis is based on an integrated luminosity of $\sim 1.3 \text{ fb}^{-1}$ collected by the CDF II detector between February 2002 and November 2006. After trigger and analysis selection criteria have been applied, the sample yields more than 1100 fully reconstructed $B_s \rightarrow D_s^- \pi^+$ candidates with $D_s \rightarrow \phi \pi^-$ and $\phi \rightarrow (K^+K^-)$. In addition, the sample reconstructed as $B_s \rightarrow D_s^- \pi^+$ includes a similar number of partially reconstructed $B_s$ candidates, for example, $B_s \rightarrow D_s^- \rho^+$ followed by $\rho^+ \rightarrow \pi^+\pi^0$ where the $\pi^0$ is not reconstructed, that can contribute to the lifetime measurement and double the number of events available for analysis. The increase comes with an uncertainty due to missing particles or incorrect mass assumptions, but the uncertainty can be properly accounted for and folded into the likelihood formulation.

*aCharge conjugation is implied throughout this Proceeding.
2 Analysis Method

A data sample rich in hadronic $B$ decays was selected with a three-level trigger system that searches for tracks displaced from the primary vertex. The trigger level requirements preferentially select longer lived particles, shaping the proper time distribution. Thus the exponential distribution of lifetimes no longer extends down to $ct = 0$. Instead there is a visible “trigger turn-on” in the distribution, which is visible in Figure 2.

The lifetime of the $B_s$ meson is determined from two sequential fits. The first fit is a maximum likelihood fit to the invariant mass distribution of candidates reconstructed as $D_s^- (\phi \pi^-) \pi^+$ with $m(D_s^- (\phi \pi^-) \pi^+) \in [4.85, 6.45]$ GeV/$c^2$. This fit determines the relative fractions of various decay modes and backgrounds in the data sample. Using the fractions determined in the mass fit as inputs, the second fit is to the proper decay time distributions for the $B_s$.

The lifetime of the $B_s$ meson is determined from an unbinned likelihood fit to the $B_s$ candidates with invariant masses in $[5.00, 5.45]$ GeV/$c^2$. We use separate probability distribution functions (PDFs) for the fully reconstructed (FR) modes, partially reconstructed (PR) modes, and the backgrounds. For the lifetime fit the variable of interest is the proper decay time, defined as $ct = \frac{L_{xy} m_{rBD}^2}{p_{TB}}$ where $L_{xy}$ is the decay length projected along the transverse momentum of the $B_s$, $p_{TB}$. Notice that the reconstructed mass is used instead of the world average $B_s$ mass. A salient feature of this analysis is the treatment of partially reconstructed $B_s$ mesons as signal events that contribute to the lifetime measurement. Since in the partially reconstructed cases $L_{xy}$, $m_{rBD}^2$, and $p_{TB}$ are extracted from candidates that are missing tracks after reconstruction or have the wrong mass assignment for a single daughter particle, a multiplicative correction factor $K$ to the decay length is needed. $K$ is defined as $K = \frac{1}{\cos \theta_{PR}} \frac{p_{T(PR)}}{p_{T(true)}} \frac{m_{BD(true)}}{m_{BD(PR)}}$ where $\theta_{PR}$ is the angle in the $x - y$ plane between the true momentum of the $B_s$ and the momentum of the partially reconstructed $B_s$.

There are separate PDFs for the three categories of proper time distributions. How each component is treated depends on its decay structure and whether it can provide information about the $B_s$ lifetime. The FR and PR PDFs depend on $\tau_B$, the $B_s$ lifetime, while the background PDFs have fixed shapes.

- **Fully Reconstructed**: The only fully reconstructed mode is the $D_s \pi$. The root functional form of the FR PDF (given in Equation 1) is an exponential with decay constant $\tau(B_s)$ convoluted with a Gaussian resolution function with width $\sigma$. A multiplicative “efficiency curve” of the form given in Equation 2 accounts for the trigger and analysis selection criteria. The shape parameters $(\sigma, \beta_i, N_i$ and $\tau_i)$ of the PDF come from a fit to a simulated $B_s$ sample where the lifetime used for generation is known. All the parameters for the PDF except for $\tau(B_s)$ are then fixed in the final fit to data. Note that although $\sigma$ is intended to be a detector resolution, it floats along with the efficiency curve parameters during the fits to the Monte Carlo. During this process it becomes correlated with the other parameters describing the overall PDF shape and loses some of its physical meaning.

$$P_{FR}(ct) = \left[ \frac{1}{c\tau} e^{-\frac{ct}{\tau}} \otimes \frac{1}{\sqrt{2\pi}\sigma} e^{-\frac{(ct - \beta_i)^2}{2\sigma^2}} \right] \cdot \text{eff}(ct) \quad (1)$$

$$\text{eff}(ct) = \begin{cases} 0 & \text{if } ct \leq \beta_i \\ \sum_{i=1}^{3} N_i \cdot (ct - \beta_i)^2 \cdot e^{-\frac{(ct - \beta_i)^2}{\tau_i}} & \text{if } ct > \beta_i \end{cases} \quad (2)$$

- **Partially Reconstructed**: As they also come from $B_s$ mesons, the $D_s K$, $D_s \rho$, $D_s^* \pi$, and other partially reconstructed modes can contribute to the $B_s$ lifetime measurement. However, a multiplicative correction factor $K$ to the decay length is needed. The PR PDF
given in Equation 3 is similar to the FR PDF with an additional convolution with the $K$ factor distribution for each mode. There are separate efficiency curve parameters for each mode, again determined from fits to simulation.

$$P_{FR}(ct) = \left[ \frac{1}{cT} e^{\frac{-ct}{cT}} \otimes e^{\frac{-1}{\sqrt{2\pi}K\sigma}} \right] \cdot \frac{-(K \cdot ct - ct')^2}{2K^2\sigma^2} \otimes K \cdot p(K) \cdot \text{eff}(ct) \quad (3)$$

- **Background:** The backgrounds can either come from single-$B$ modes (e.g., $B^0/B^- \rightarrow D^-\pi^+$, $B^0 \rightarrow D_sX$, and $\Lambda_b \rightarrow \Lambda_cX$), or they can be combinatorial in nature. There are two combinatorial background proxies available: the $B_s$ upper sideband taken from the $m_B$ interval $[5.7, 6.4]$ GeV/$c^2$ and the $D_s$ sidebands. Both proxies contain a mixture of real $D_s$+track and fake $D_s$+track events. The background PDFs come from fits to simulated samples (for $b$-hadron backgrounds) or from the combinatorial background proxies. All the background shape parameters are fixed in the final lifetime fit.

### 3 Fit Results

The fit procedure was tested extensively on three control samples: $B^0 \rightarrow D^-\pi^+$ with $D^- \rightarrow K^+\pi^-\pi^-$, $B^0 \rightarrow D^+\pi^+\pi^-$ with $D^+ \rightarrow D^0\pi^-$ and $D^0 \rightarrow K^+\pi^-$, and $B^+ \rightarrow D^0\pi^+$ with $D^0 \rightarrow K^+\pi^-$ before performing a blinded $B_s$ fit. Good agreement with the PDG values of the $B^0$ and $B^+$ lifetimes was found.

The projection of the $B_s$ mass fit is shown in Figure 1. The lifetime of $cT_{B_s} = 455.0 \pm 12.2\text{(stat.)}\mu s$ is obtained from the full fit. The $ct$ projection of the fit result is plotted in Figure 2.

We consider several sources of systematic uncertainty: combinatorial background fraction, modeling simulated backgrounds from single-$B$ decays, effect of reweighting the Monte Carlo to match the data’s $p_T$ and trigger distributions, impact parameter correlation, and detector alignment. The largest contribution to the systematic uncertainty comes from the uncertainty on the total amount of combinatorial background and the relative amount of promptly-produced real-$D_s$ background (the component with the lowest mean lifetime). The invariant mass shapes
Figure 2: ct projection of lifetime fit results for events reconstructed as $B_s \rightarrow D_s^- (\phi \pi^-)\pi^+$

for the real-$D_s$ and fake-$D_s$ combinatorial backgrounds are very similar, and several mass fit configurations that are equally valid from first principles yield dissimilar background fractions. The variations in the fractions returned from these mass fits are used to evaluate the systematic uncertainty.

4 Conclusions

A fit for the $B_s$ lifetime in $\sim 1.3$ fb$^{-1}$ of data reconstructed as $B_s \rightarrow D_s^- \pi^+$ is performed. The fit utilizes both fully and partially reconstructed modes. We measure

$$c\tau(B_s) = 455.0 \pm 12.2 \text{(stat.)} \pm 7.4 \text{(syst.)}\mu m$$

$$= 1.518 \pm 0.041 \text{(stat.)} \pm 0.025 \text{(syst.)} \text{ps}$$

The ratio of this single result and the world average $B^0$ lifetime yields $\tau(B_s)/\tau(B_d) = 0.99 \pm 0.03$. This agrees well with the theoretical prediction of $\tau(B_s)/\tau(B_d) = 1.00 \pm 0.01$. More information about this analysis can be found on its public webpage.  

References

Searching for Higgs Decaying to
$H \rightarrow WW^* \rightarrow \mu + \tau_{had}$ and $H \rightarrow WW^* \rightarrow ee$ at DO

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A search for the Higgs boson in $H \rightarrow WW^* \rightarrow ee$ and $H \rightarrow WW \rightarrow \mu \tau_{had}$ decays in pp collisions at a center-of-mass energy of $\sqrt{s} = 1.96$ TeV is presented. The data have been collected by the Run II DO detector. In order to maximize the sensitivity multivariate techniques such as artificial neural networks (NN), matrix element methods and likelihoods are used. No excess above the Standard Model background is observed and limits on the production cross section times branching ratio $\sigma \times BR(H \rightarrow WW^* \rightarrow ee)$ for Higgs masses between 115 and 200 GeV are set.

1 Introduction

Two searches for the Higgs boson decaying to the $WW^*$ final state are presented. The dileptonic decay mode with two electrons in the final state and $H \rightarrow WW \rightarrow \mu + \tau_{had}^{1-prong}$ leading to final states with one muon, a jet originating from a hadronically decaying tau and missing transverse momentum have been studied. The $\mu + \tau_{had}$ analysis has been performed using an integrated luminosity of about $\sim 1$ fb$^{-1}$ of RunII data recorded between 2001 and 2006, known as RunIIa. The $ee$ analysis is performed using data from June 2006 until August 2007, known as RunIIb, corresponding to an integrated luminosity of 1.2 fb$^{-1}$. These decay modes in combination with other dileptonic decay modes provide the best sensitivity to a Standard Model (SM) Higgs boson search at the Tevatron at a mass of $m_H \sim 160$ GeV/c$^2$. In order to maximize the signal to background separation multivariate techniques are used. If combined with searches exploiting the $WH$ and $ZH$ associated production, these decay modes increase the sensitivity for the Higgs boson searches.
2 Event Selection

The signal is characterized by two leptons, missing transverse momentum \( (p_T) \) and little jet activity. For the \( \mu + \tau_{\text{had}} \) selection one isolated muon with \( p_T^\mu > 12 \text{ GeV} \) and one isolated tau with \( p_T^\tau > 10 \text{ GeV} \) are required. A hadronically decaying tau lepton is characterized by a narrow isolated jet with low track multiplicity. Three tau types are distinguished:

- \( \tau \)-type I: A single track with a calorimeter cluster without any electromagnetic subclusters (1-prong, \( \pi \)-like).
- \( \tau \)-type II: A single track with a calorimeter cluster and electromagnetic subclusters (1-prong, \( \rho \)-like).
- \( \tau \)-type III: Two or three tracks with an invariant mass below 1.1 or 1.7 GeV, respectively (3-prong).

Due to the large background contamination \( \tau \)-type III is neglected in the analysis. For the dielectron analysis the leading electron is required to have \( p_T^e > 20 \text{ GeV} \) and the trailing electron to fulfill \( p_T^\mu > 15 \text{ GeV} \). Subsequently most of the QCD background is removed by selection requirements on the missing transverse energy \( E_T \) and the scaled missing transverse energy \( E_T^{\text{Scaled}} \), defined as
\[
E_T^{\text{Scaled}} = \frac{E_T}{\sqrt{\sum_{\text{jets}} E_T^2 || E_T}}
\]
which is the \( E_T \) divided by the \( E_T \) resolution.

This quantity is particularly sensitive to events where the missing energy could be a result of mismeasurements of jet energies in the transverse plane. A requirement on the minimal transverse mass, \( M_T^{\text{min}} = \sqrt{2 \cdot E_T \cdot p_T \cdot (1 - \cos(\Delta \phi))} \), between one of the leptons and the \( E_T \) reduces further the various background processes. Most of the \( Z/\gamma^* \rightarrow \ell \ell \) events are rejected by requiring the sum of the momentum of \( p_T^\mu + p_T^\tau + E_T \) to be within given lower and upper boundaries and by requiring the invariant mass to be less than the \( Z \) peak. The \( t\bar{t} \) contribution is reduced by requiring low values \( H_T \) which is defined as the scalar sum of the transverse momenta of all jets in the event. A large fraction of remaining back-to-back \( Z/\gamma^* \rightarrow \ell \ell \) is reduced by rejecting events with a wide opening angle between the leptons. Since the signal kinematics change as a function of the Higgs mass the selection is applied in a mass dependent way.

3 \( W + \text{jet/\gamma} \) and Multijet background estimation

The background contribution from QCD multijet production where jets are misidentified as leptons is estimated from the data by using like-sign lepton events of each analysis which were selected by inverting calorimeter isolation criteria. The samples are normalized to data as function of the lepton \( p_T^i \) \( (i = 1, 2) \) in a region of phase space which is dominated by multijet production. In the \( \mu + \tau_{\text{had}} \) analysis the shape of the \( W + \text{jet/\gamma} \) background is taken from MC, the normalization however is estimated using data.

4 Multivariate Techniques

At the final stage of the selection as described in Section 2 the remaining background is dominated by electroweak \( W + \text{jets/\gamma} \) and diboson production. To improve the background reduction further multivariate techniques have been used, in the \( \mu\tau_{\text{had}} \)-analysis a likelihood approach and for the \( ee \) final state artificial neural networks. For the likelihood approach two different likelihoods sensitive to different event properties are used. One likelihood is based on input distributions associated with the selected tau and the second one on kinematical properties of the particular event. All likelihoods are constructed according to formula (1). A non-negligible
fraction of the tau-candidates are electrons misreconstructed as taus which is taken into account by further constructing both classes of likelihoods for that particular events.

\[
\mathcal{L} = \frac{P_{Sig}(x_1, x_2, \ldots)}{P_{Sig}(x_1, x_2, \ldots) + P_{Bkgd}(x_1, x_2, \ldots)} \approx \frac{\prod_i P_{Sig}^i}{\prod_i P_{Sig}^i + \prod_i P_{Bkgd}^i} = \frac{\prod_i P_{Sig}^i / P_{Bkgd}^i}{\prod_i P_{Sig}^i / P_{Bkgd}^i + 1}
\]

(1)

Where \(P_{Sig}^i\) represents the signal and \(P_{Bkgd}^i\) the background value for a given bin \(i\). The value of the input distributions for bin \(i\) are given by the variables \(x_i\). \(P_{Sig}^i \equiv P_{Sig}(x_i)\) and \(P_{Bkgd}^i \equiv P_{Bkgd}(x_i)\) represents the probability density functions for the topological variables. These likelihoods are constructed for each Higgs boson mass point and both tau-types. The resulting likelihood distribution for \(m_H = 160\) GeV and tau-type I is displayed in Fig. 1. Using both likelihood classes a further selection requirement is applied. These selections have been optimized for each sample, tau type and Higgs mass. For the dielectron analysis the separation of signal from background is done using an artificial neural network. A separate NN is trained for each Higgs boson mass tested. A list of input variables has been derived based on the separation power of the various distributions. Those variables can be divided into three classes, object kinematics, event kinematics and angular variables. An additional input variable is a discriminant constructed using the matrix element (ME) method. Leading-order parton states for either signal or \(WW\) background are integrated over, with each state weighted according to its probability to produce the observed measurement.

![Figure 1: Distribution of (a) the kinematical likelihood distribution after all selection requirements (b) the neural network output after all selection requirements for data (points with error bars), background simulation (histograms, complemented with the QCD expectation) and signal expectation for \(m_H = 160\) GeV (empty histogram).](image)

The various signal contributions \(H \rightarrow WW \rightarrow ee\) and \(H \rightarrow WW \rightarrow \mu\tau_{had}\) are given by the solid red line.

5 Results and Conclusion

Limits on the cross section for Higgs boson production times the branching fraction into the discussed final states are derived at the 95% Confidence Level (CL). Whereas the number of expect signal events, expected background events and observed data is used in the \(\mu + \tau_{had}\) analysis, for the dielectron analysis the shape of the NN output distributions are taken into account as well.

There are various sources of systematic uncertainties included in the calculation of the expected and observed limits: Lepton identification and reconstruction efficiencies (0.3–13%), trigger efficiencies (5%), jet energy scale calibration in signal and background events(< 2%),
track momentum calibration (4%), detector modeling (1% for signal, 5-10% for background), PDF uncertainties (4%), modeling of multijet background (30%) and theoretical cross section (di-boson 7%, tt 16%). The uncertainty on the modeling of the electroweak $W + \text{jet}/\gamma$ production has been estimated to be 2.5-17.5%. The systematic uncertainty on the luminosity is mainly a combination of the PDF uncertainty, uncertainty on the NNLO Z cross section (4%) and data/MC normalization.

The total uncertainty on the background is approximately 20% and for the signal efficiency is 10%. The effects of these uncertainties on the NN output distribution shapes were also studied and included as additional systematic uncertainties. The expected and observed limits as function of the Higgs mass are shown in Fig. 2. After all selection requirements both the neural network output distribution and the likelihood distributions agree with data within their uncertainties with the expected backgrounds. Thus limits are set on the production cross section times branching ratio $\sigma \times BR(H \rightarrow WW^*)$. We calculate the limits for each channel using a modified frequentist approach CLs method with a log-likelihood ratio (LLR) test statistic. To minimize the degrading effects of systematics on the search sensitivity, the individual background contributions are fitted to the data observation by maximizing a profile likelihood function for each hypothesis$^5$.

![Graph](image)

Figure 2: (a) Expected and observed limit for the $\mu + \tau_{\text{had}}$ analysis (b) Expected and observed limit for the dilepton analysis

A search has been performed for the $H \rightarrow WW \rightarrow \ell\ell$ decay signature from the gluon-gluon fusion production of the Standard Model Higgs boson in leptonic channels with either muons and taus or two electrons, using data corresponding to an integrated luminosity altogether of $\approx 2.2$ fb$^{-1}$. No evidence for the Higgs particle is observed and no region of the SM Higgs can be excluded.

References

4. Search for $H \rightarrow WW \rightarrow \text{dileptons}$, 300-325 pb$^{-1}$, hep-ex/0508054.
SEARCH FOR WMAP-COMPATIBLE SIMPLE $SO(10)$ SUSY GUTs

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Unification of GUT-scale $t - b - \tau$ Yukawa couplings is a significant feature of simple $SO(10)$ SUSY GUTs. Here we present the results of a search that used the Markov Chain Monte Carlo technique to investigate regions of Yukawa unification and WMAP-compatible dark matter relic density in $SO(10)$-like MSSM parameter spaces. We mention the possible LHC signatures of Yukawa unified scenarios and discuss the consequences for dark matter.

1 Introduction

Grand Unification is regarded as an inspirational ingredient of models that claim to explain the fundamental laws of nature. A highly motivated scenario in this context originates from grand unification via the $SO(10)$ gauge group $^1$. Simple supersymmetric implementations of $SO(10)$ GUTs unify all matter fields in each generation within a 16-dimensional irreducible representation and two Higgs doublets of the MSSM within a 10-dimensional irreducible representation. Such a formalism automatically includes heavy right-handed neutrino states and the resulting structure of the neutrino sector implies a successful theory of baryogenesis via intermediate scale leptogenesis. Moreover the $SO(10)$ models are left-right symmetric and this enables them to provide a solution to the strong CP problem and to naturally induce R-parity conservation.

Besides gauge coupling unification, $SO(10)$ SUSY GUTs additionally require the unification of 3rd generation Yukawa couplings at the GUT scale ($M_{GUT}$). This is explicitly seen from the expression of the superpotential above $M_{GUT}$, which takes the form $\hat{f} \equiv f\bar{\psi}_{16}\phi_{16}\phi_{10} + \cdots$. An exact unification occurs at tree level while several percent corrections arise at the loop level. As a result we can assume that any sign of Yukawa unification from observations could be a hint to the existence of $SO(10)$ SUSY GUTs.

Our aim in the study is to investigate the characteristics that arise in a SUSY model when GUT scale Yukawa unification is imposed and to determine the experimental signatures that would distinguish such Yukawa-unified models from the others. In this context we assume a theoretical framework where nature is explained by an $SO(10)$ symmetry above $M_{GUT}$. Then at $M_{GUT}$, $SO(10)$ breaks to MSSM plus some heavy right-handed neutrino states. At the weak scale, the content of the theory is equivalent to that of the MSSM.

The GUT scale soft SUSY breaking parameters are constrained by the requirement of the $SO(10)$ symmetry. Unified representations would favor common SSB masses "$m_{16}$" for the scalars and "$m_{10}$" for the Higgses, but in order to achieve REWSB, SSB Higgs masses should be split, satisfying $m_{H_d} > m_{H_u}$. Here we examine two different methods to generate the necessary Higgs splitting: The first approach defines the Higgs masses as $m_{H_{u,d}}^2 = m_{10}^2 \pm 2M_D^2$. Here splitting is parametrized by $M_D$, which is the magnitude of the D-terms in the scalar potential.
of the extra $U(1)$ group that is a by-product of the $SO(10)$ breaking. The parameters of this GUT scale Higgs input (GSH) scenario are

$$m_{16}, m_{10}, M^2_D, m_{1/2}, A_0, \tan \beta, sgn(\mu)$$  \hspace{1cm} (1)

The second approach was put forward in order to generate Yukawa unified solutions with low $\mu$ parameter and low $m_A$. Such solutions were found to exist by Blaszek, Dermisek and Raby at a study where they assumed perfect Yukawa unification at $M_{\text{GUT}}$ and made a fit to the weak scale observables$^2$. In order to seek similar solutions, we start with GSH parameters at GUT scale, but additionally provide $\mu$ and $m_A$ as inputs. We run $m_{H_u,d}$ down, and at $Q = M_{\text{SUSY}}$ we compute what $m_{H_u}, m_{H_d}$ should have been in order to give our input $\mu$ and $m_A$, and run back up using these new boundary conditions. This weak scale Higgs (WSH) scenario has the parameters

$$m_{16}, m_{10}, M^2_D, m_{1/2}, \tan \beta, m_A, \mu$$  \hspace{1cm} (2)

We take the GSH and WSH scenarios and search in their parameter spaces for regions having a good Yukawa unification where Yukawa unification is parametrized by $R = \frac{\max(f_u,f_d)}{\min(f_u,f_d)}$. Additionally we seek sub-regions that are consistent with the WMAP measurements of dark matter relic density $\Omega h^2 \leq 0.1$. There have been previous searches using the GSH input based on random scans and they were able to achieve less than few percent of Yukawa unification for $\mu > 0$\textsuperscript{4}. However the dark matter relic densities for these solutions were always much higher than the WAMP upper bound. Here we implement the Markov Chain Monte Carlo (MCMC) technique which enables a much more efficient scanning of multi-dimensional parameter spaces. The following two sections summarize the MCMC technique, the characteristics of the regions found by utilizing it, and dark matter-related consequences of the $SO(10)$ SUSY GUTs.

2 The MCMC Search and the Yukawa-unified Solutions

A Markov Chain is a discrete time, random process where given the present state, the future state only depends on the present state, but not on the past states\textsuperscript{5}. The MCMC samples from a given parameter space as follows: It takes a starting point, and it generates a candidate point $x^c$ from the starting point $x^t$ using a proposal density $Q(x^t,x^c)$. The candidate point is accepted to be the next state $x^{t+1}$ if the ratio $p = \frac{P(x^c)Q(x^c|x^t)}{P(x^t)Q(x^t|x^c)}$ (where $P(x)$ is the probability calculated for the point $x$) is greater than a uniform random number $a = U(0,1)$. If the candidate is not accepted, the present point $x^t$ is retained and a new candidate point is generated. By repeating this procedure continuously the Markov Chain eventually converges at a target distribution around a point with the highest probability.

Our MCMCs were directed to approach regions with $R \sim 1.0$ and $0.094 \leq \Omega h^2 \leq 0.136$. Here we used a Gaussian distribution for the proposal density $Q$, and approximated the likelihood of a state to $e^{-\chi^2/2}$. We also chose multiple starting points ($\sim 10$) in order to search a wider range of the parameter space. We used ISAJET 7.75 for sparticle mass computations and micrOMEGAs 2.0.7 for DM relic density calculations. The MCMCs successfully located some regions with good $R$ and $\Omega h^2$. Here we will mostly emphasize the results from the GSH scenario.

Figure 1 shows the compatible regions in $m_{16} - m_{10}$ and $m_{16} - m_{16}/A_0$ planes. The light blue dots have $R \leq 1.1$, the dark blue dots have $R \leq 1.05$, the orange dots have $R \leq 1.1$ plus $\Omega h^2 \leq 0.136$ and the red dots have $R \leq 1.05$ plus $\Omega h^2 \leq 0.136$. We see that Yukawa unification occurs only at the regions where the input parameters are strongly correlated, having $m_{10} \simeq 1.2 m_{16}$ and $A_0 \simeq -(2 - 2.1) m_{16}$. A good DM relic density is achieved only at the constrained regions that have $m_{16} \sim 3 - 4$ TeV. Further search showed that $m_{1/2}$ takes the lowest possible values for a given $m_{16}$, generally giving $\sim 100$ GeV, and decreases steadily with increasing $m_{16}$.

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These highly confined parameter regions lead to strongly constrained mass spectra, and hence to significant LHC signatures. We see that Yukawa-unified solutions are distinguished by their heavy 1st/2nd generation scalars ($>2\,\text{TeV}$), lighter 3rd generation scalars ($\sim\text{TeV}$) and light gauginos (few hundred GeV). All Higgses except $h^0$ are about $1-3\,\text{TeV}$. Figure 2 shows the distribution of selected points on $m_{\tilde{t}_1}$ vs $m_{\tilde{g}}$ plane (left), and on $m_{h^0}$ vs $m_{\tilde{\chi}_2^0} - m_{\tilde{\chi}_1^0}$ plane (right) for the GSH scenario. The requirement of $\Omega h^2 < 0.136$ favors a gluino mass range around $350-450\,\text{GeV}$, which means we would expect a large amount of gluino pair production at the LHC with cross sections about $\sim 100\,\text{pb}$. The gluinos decay via 3-body channels such as $\tilde{g} \to \tilde{\chi}_1^0 b\bar{b}$, $\tilde{\chi}_1^0 b\bar{b}$, $\tilde{\chi}_1^+ t\bar{b}/b\tilde{t}$, since 2-body channels are closed due to the high squark masses. On the other hand favored $\tilde{\chi}_2^0 \simeq \tilde{\chi}_1^+$ mass range is $100-150\,\text{GeV}$, which leads to gaugino pair production cross sections about $\sim 10\,\text{pb}$, while $m_{\tilde{\chi}_2^0} \sim 50-75\,\text{GeV}$. The preferred mass difference $m_{\tilde{\chi}_2^0} - m_{\tilde{\chi}_1^0}$ is $52-65\,\text{GeV}$ which is smaller than $m_{Z,h^0}$, therefore $\tilde{\chi}_2^0$ decays are dominated again by 3-body channels as $\tilde{\chi}_2^0 \to b\tilde{g}\tilde{\chi}_1^0, g\tilde{g}\tilde{\chi}_1^0, \tilde{t}\tilde{\chi}_1^0$.

As a result we expect the $SO(10)$ models to manifest themselves as multi b-jet plus low missing $E_T$ final states at the LHC. Additionally it would be possible to investigate the OS/SF dilepton channels where the dilepton invariant mass is bound by $m_{\tilde{\chi}_2^0} - m_{\tilde{\chi}_1^0}$ and try to reconstruct the $\tilde{g} \to b\bar{b}\tilde{\chi}_2^0 \to bbl\tilde{\chi}_1^0$ cascades.
3 Consequences for Dark Matter

The majority of solutions shown in Figures 1 and 2 have excess DM relic density, while only a small portion of them gives $\Omega h^2 < 0.136$. To investigate the mechanism that provides the efficient annihilation in the DM-allowed solutions, we check the behavior of solutions in $m_{h^0}$ and $m_A$ mass resonances. Figure 3 shows the distribution of Yukawa-unified points on the $m_A - 2m_{\tilde{\chi}^0_1}$ vs. $m_{h^0} - 2m_{\tilde{\chi}^0_1}$ plane for GSH (left) and WSH (right) scenarios. In the GSH plot all DM-allowed solutions are on the $m_{h^0} \simeq 2m_{\tilde{\chi}^0_1}$ line, which shows that the relic density is reduced by annihilation via a light Higgs resonance. On the other hand $m_A > 2m_{\tilde{\chi}^0_1}$, so there are no A resonance solutions. Turning to the WSH scenario we see that annihilation via both $h^0$ and A resonances are at work. Actually the majority of solutions are generated by the latter due to the relatively small A masses allowed within the WSH scenario. However all of these solutions have $B_{s \to \mu\mu}$ branching ratios higher than the latest reported CDF upper limit $5.8 \times 10^{-8}$. So these A resonance points are ruled out, leaving us with only the $h^0$ resonance solutions.

One could also devise alternative methods for reducing the excess DM relic density. One way could be to assume that $\tilde{\chi}^0_1$ is not the LSP, but can decay to other candidates such as gravitino or axino via the mode $\tilde{\chi}^0_1 \to \gamma G/\tilde{a}$. Lifetime of $\tilde{\chi}^0_1$ would be long enough to let it escape the detectors. The resulting relic density would be $\Omega_{\tilde{\chi}^0_1} = \frac{m_{\tilde{G}/\tilde{a}}}{m_{\tilde{\chi}^0_1}} \Omega_{\tilde{\chi}^0_1}$ since the $\tilde{G}/\tilde{a}$s inherit the thermally produced neutralino relic number density. $\tilde{G}$ LSP can only reduce the relic density a few times, which is not satisfactory for our case, but axinos with $m_\tilde{a} \leq 1$ MeV would allow for a mixed cold/warm DM solution which can reduce the relic density below the WMAP bound.

Another method to reconcile $\Omega h^2$ is to relax some universalities in the GUT scale SSB terms. For example, increasing the $U(1)$ gaugino mass term $M_1$ (while keeping $M_{2,3} = m_{1/2}$) brings $m_{\tilde{\chi}^0_1}$ close to $m_{\chi^+_1}$, hence making $\tilde{\chi}^0_1$ more wino-like and inducing bino-wino coannihilation. A further possibility is to lower the 1st/2nd generation masses $m_{16}(16)$ (while keeping $m_{16}(3) = m_{16}$), which enables neutralinos to annihilate via light $\tilde{q}_R$ exchange and leads to neutralino-squark coannihilation.

4 Conclusions

By performing scans on the parameter space of simple $SO(10)$ SUSY GUT scenarios using the MCMC technique, we showed that solutions with both 5−10% Yukawa unification and WMAP-compatible $\Omega h^2$ can exist around $m_{16} \sim 3−4$ TeV. These regions defined by strictly constrained relations among the GUT scale inputs generate special sparticle mass relations that lead to distinguishable signatures at the LHC. With multi-TeV scalars, 350 − 450 GeV gluinos and...
50 – 150 GeV light gauginos, we expect dominant gluino production followed by 3-body cascade decays which will end up in b-rich multijet final states, occasionally including OS/SF lepton pairs from $\tilde{\chi}_0^0$ decays. Moreover the possibility to lower $\Omega h^2$ by assuming a LSP or introducing SSB non-universalities marks wider parameter space regions as compatible. So we can conclude that $SO(10)$ SUSY GUTs provide motivated scenarios with robust signatures relevant to be tested soon at the turn-on of the LHC.

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TOWARDS AN ESTIMATION OF THE MUON MULTIPLICITY IN ANTARES

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Antares is an underwater neutrino telescope under completion in the Mediterranean Sea. The telescope consists of a 3D network of photomultipliers able to detect the Cherenkov light emitted by relativistic charged particles. From January to December 2007 the Antares detector was composed of 5 lines and since December 7th 2007, it comprises 10 lines. The 12 line detector is being completed in 2008. In this paper we present a study with Monte-Carlo and data to estimate the muon multiplicity of atmospheric bundles, which are the most intense source of background to the signal of up-going muons.

1 Introduction

Antares is a high energy neutrino telescope located at 40 km off Toulon. The main goal of Antares is to detect neutrinos as a new probe to study the sky at energies greater than 1 TeV. Neutrinos are indeed neither deflected by the magnetic fields, nor absorbed by the interstellar dust. This enables them to travel on cosmological distances through the Universe. They can also emerge from regions of very high matter density, bringing us information on the core of astrophysical objects where catastrophic phenomena can happen. This work describes the characterization of atmospheric muons bundles which are the most intense background for neutrino detection.

2 Detection method

High energy cosmic muon neutrinos can interact with the surrounding detection medium and produce a muon. In water this muon travels faster than light and emit Cherenkov light in a cone whose angle is defined by:

$$\cos \theta = \frac{1}{\beta_n} \Rightarrow \theta \approx 42^\circ$$

Antares detects this dim light with a three-dimensional network of photomultipliers housed in glass spheres (called Optical Modules - OM).

Antares is installed at 2475 m depth. The final detector is constituted of 12 vertical lines, measuring 450 m of which 350 m are instrumented. Lines lay 60 m apart from each other. Each line is connected to a junction box through an interlink cable and the junction box is connected to shore via a 40 km long electro-optic cable for power and data transmission. The lines are composed of 25 floors of 3 optical modules each pointing downwards at 45° from the vertical.

The Antares PMT network is optimized to look downwards, i.e. to detect upward-going muons induced by neutrinos that have interacted in the Earth. Contrarily to these neutrinos, muons originating from atmospheric showers cannot travel through the Earth.
Because of the weak cross section and the low emission flux of the sources, few cosmic neutrinos are expected to be detected. This small signal has to be disentangled from two kinds of physical backgrounds:

- **Up-going atmospheric neutrinos** created by interaction of primary cosmic rays in the atmosphere which can interact near the detector, giving an up-going muon.
- **Down-going muons** that are sometimes mis-reconstructed as up-going ones. Due to the intensity of the atmospheric muon flux, this remains as the most intense source of background.

In this analysis, we will focus on the muon multiplicity, defined as the number of muons crossing the detector simultaneously, as a possible variable to reject this background. In our simulation, a “muon” is further considered as such if it produces at least 6 hits (minimum requisite for the reconstruction) in the detector.

### 3. Goal of the multiplicity study

When a primary cosmic ray (mostly proton) collides with a nitrogen or an oxygen nucleus, it produces an atmospheric cascade of secondary particles (called "shower" or "bundle" of particles), as illustrated in figure 1. During its development, the shower produces charged mesons which decay into muons and neutrinos.

Produced muons propagate to the sea level. Only the most energetic ones (E > 500 GeV) can go through more than two kilometres of water and reach the detector. As shown in figure 2, the down-going flux is dominated by the atmospheric muons by six orders of magnitude compared to the muons from the atmospheric neutrinos, for muon energies above 1 TeV at the detector.

**Figure 1:** Bundle of secondary particles produced by a cosmic ray interaction.

**Figure 2:** Muon flux versus zenith angle. \( \cos \theta = 1 \) represents the down-going muons and \( \cos \theta = -1 \) the up-going muons.

The reconstruction algorithm used in the Antares data analysis is optimised for a single (up-going) muon passing through the detector. However, 90% of the down-going events comprise several muons (see muon multiplicity in figure 3). These down-going multi-muons can produce some tricky effects in the reconstruction and in the understanding of the detector. This study is an attempt to separate single muons from multi-muons using appropriate variables with the aim of studying multi muon effects in the detector. Muon multiplicity depends on the primary energy and the thickness of water that the muons traverses (i.e. the incidence angle). Multi-muon events, when reconstructed as a single muon, are expected to produce bad quality fits. The study presented here is based on a variable relying on the time residuals after reconstruction.
4 Method

As illustrated in figure 4, the muon propagates in a straight line at the speed of light (in vacuum). The arrival time of a direct Cherenkov photon at the optical module is determined by

\[ t_{\text{direct}} = t_0 + \frac{1}{c} \left( \frac{k}{\tan \theta_C} \right) + \frac{1}{v_q} \left( \frac{k}{\sin \theta_C} \right) \]  

(1)

where \( v_q \) is the group speed of light in water, \( t_0, x, y, z \) define an arbitrary point on the muon trajectory, \( l \) and \( k \) are defined in figure 4 and \( \theta_C \) is the Cherenkov angle. The first term in equation (1) accounts for the muon travel time, from \( t_0 \) to the moment when it emits a detectable Cherenkov photon, while the last term represents the time needed by this photon to reach in a straight line the photomultiplier. Photons can however diffuse in the water before hitting a photomultiplier, this diffusion inducing some delay with respect to a direct hit. Moreover at high energy (above \( \approx 500 \text{ GeV} \)) the muon energy loss begins to be dominated by catastrophic losses producing secondary electromagnetic showers. These particles produce additional Cherenkov light leading to further delays. These delays create larger time residuals, defined as the difference between the arrival time on the photomultiplier and the computed hit time of a direct photon emitted by the reconstructed muon.

If a bundle of muons distant from each other by a few meters is reconstructed as a single muon, the time residual distribution will be broader than the one of isolated single muons. Therefore we introduce the variable Small Residual Fraction (SRF) which is the ratio of the number of hits with time residuals between \(-5 \text{ ns} \) and \(5 \text{ ns} \) over the number of hits with residuals between \(-50 \text{ ns} \) and \(100 \text{ ns} \).

5 Analysis

The simulation was done using the CORSIKA\(^3\) code for a 5-line detector.

The time residual distribution is shown in figure 5. As expected, the peak is narrower for single muons than for multi muons. The mean value of SRF is \(0.39\) (figure 6). For single muons it is \(0.46\) and is smaller for multi-muons. The single muon content increases with a cut on SRF, as shown in figure 7. For example with SRF \(> 0.5\), we have \(36.2\%\) of single muons contaminated by \(4.45\%\) of multi-muons. The mean multiplicity decreases from \(2.8\) to \(1.3\) and single muons represent \(53\%\) of the muons in the events. Fig. 8 shows a comparison of the SRF distribution between the Monte-Carlo sample and one month of real data (June 2007 - 18 days effective live time). By comparing with an other Monte-Carlo\(^4\), we find that one reason of disagreement is that the multiplicity of Corsika is lower than the observed one in data.
Figure 5: Time Residual distribution for different ranges of multiplicities

Figure 6: Small Residual Fraction distribution for all multiplicities and for ranges of multiplicities

Figure 7: Efficiency, purity and composition of lower SRF cut.

Figure 8: Small Residual Fraction distribution for data (red dotted line) and simulation (blue line)

6 Conclusion

In Antares, the neutrino signal is polluted by a fraction of the muons coming from atmospheric bundles. In this study we have shown that using a simple discriminating variable based on time residuals, the single muon sample can be substantially enriched. We increase the purity of the single muon sample and decrease the mean multiplicity. The separation of single from multiple muons allows a better understanding of reconstruction and event selection.

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Coincidence analysis in ANTARES: Potassium-40 and muons

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A new calibration technique using natural background light of sea water has been recently developed for the ANTARES experiment. The method relies on correlated coincidences produced in triplets of optical modules by Cherenkov light of $\beta$ particles originated from $^{40}\text{K}$ decays. A simple but powerful approach to atmospheric muon flux studies is currently being developed based on similar ideas of coincidence analysis. This article presents the two methods in certain detail and explains their role in the ANTARES experiment.

1 Introduction

ANTARES is a large water Cherenkov detector operating in Mediterranean sea 40 km offshore Toulon (France) at the depth of 2470 m. An ANTARES storey includes a triplet of optical modules oriented at 45° downwards and outwards of the vertical axis. Twenty-five storeys, chained together with a step of 14.5 m, form a detector line. Several calibration systems and techniques ensure a sub-nanosecond precision of the Cherenkov pulse measurements, which is required to achieve the high angular resolution of the neutrino telescope. Important roles are played by in situ measurements using LED Beacons and Potassium-40.

Potassium-40 is a $\beta$-radioactive isotope naturally present in sea water. The energy freed in $^{40}\text{K}$ decays amounts to 1.3 MeV, that well exceeds Cherenkov threshold for electrons in water (0.26 MeV) and is sufficient to produce up to 150 Cherenkov photons. If the decay occurs in the vicinity of a detector storey, a coincident signal may be seen by two of the three optical modules constituting the triplet (local coincidence). This effect and its use in the ANTARES experiment are explained in section 2. In the case of two modules located on different storeys the probability of a genuine coincidence from $^{40}\text{K}$ is negligibly small. Instead, a signal originated from atmospheric muons could become dominating, if was not overwhelmed by the random background from $^{40}\text{K}$ and bioluminescence. In order to reduce the random background we
require a local coincidence (within 20 ns time window) at each of the two storeys rather than just a hit. Section 3 presents such an analysis for the case of adjacent storeys.

For the purpose of detailed $^{40}$K measurements a dedicated type of calibration runs is defined in ANTARES (K40 runs). During a K40 run (typically 20 minutes) all local coincidences detected in any OM triplet of ANTARES are selected by a dedicated data filter algorithm, so-called K40 trigger, and saved on disk for later processing. Importantly, this type of runs is also perfectly suitable for the adjacent floor coincidence studies.

2 Calibration with Potassium-40

Due to the difference in positions of the optical modules (the distance between OM centers is 1.0 m) the signals are detected by the two OMs with a noticeable delay. When averaging over many events (locations of decayed nucleus) this results in a time spread of 4.0 ns (RMS), that is experimentally observed (see Fig. 1). This value is in excellent accordance with Monte Carlo simulations. As one can see from Fig. 1, the coincidence peak is comparable in amplitude with the pedestal of random coincidences, which are composed of truly uncorrelated background photons (from $^{40}$K or bioluminescent emission). Ideally, if the optical modules are all identical, perfectly calibrated, and arranged in triplets symmetrically, the coincidence peak must be aligned with zero. Experimental measurements (see Fig. 1) show a small spread of the offsets (0.68 ns RMS), that confirms the high accuracy of timing calibration in ANTARES and gives a measure of disagreement between the $^{40}$K data and currently used calibration set. The present results refer to so-called Dark Room calibration, which is performed on shore before line deployment. An improvement is foreseen with the in-situ LED Beacon calibration, which is currently in progress.

The rate of correlated coincidences can be defined as the integral under the coincidence peak (excluding pedestal) normalized to the effective duration of observation period, and properly corrected for dead time of the electronics and data acquisition. We rely on a Gaussian fit to compute the rate. The observed average value amounts to 14 Hz (see Fig. 2, left), that is in good agreement with MC simulations. It has been checked experimentally that the coincidence rate is not affected by variations of the bioluminescent background. This important observation suggests that the $^{40}$K measurements can be used as a robust calibration tool. In addition, it also provides a confirmation of single-photon character of the bioluminescent emission. Indeed, if a bioluminescent process produced bunches of photons, correlated in nanosecond scale, an increase of the correlated coincidence rate would be observed for high-background runs.

It can be shown by Monte Carlo simulation that the rate of correlated coincidences is proportional to the detection efficiencies of both the optical modules involved. Thus, for a triplet
one can write three equations as follows:

\[
\begin{align*}
    r_{12} &= k s_{12} s_{2} \\
    r_{23} &= k s_{23} s_{3} \\
    r_{31} &= k s_{31} s_{1}
\end{align*}
\]  

(1)

where \( s_{12,3} \) is the sensitivity of OM1(2,3) in arbitrary units, \( r_{12,23,31} \) is the rate of correlated coincidences between OM1 and OM2 (OM2 and OM3, OM3 and OM1), and \( k \) is a normalization factor. We will assume \( k \) constant, i.e. not varying from module to module or with time. For a triplet of optical modules one can unambiguously solve system (1) and extract the three sensitivity coefficients \( s_{12,3} \). Note that the absolute normalization of the OM sensitivities requires precise knowledge of the OM angular acceptance (alternatively, OM angular acceptance can be constrained by the \(^{40}\)K measurements).

So far, OM sensitivity measurements have been done with \(^{40}\)K for ten detector lines of ANTARES, helping to locate mistuned and degraded modules. A significant drop of sensitivity was observed for some optical modules during first months of operation. In average a decrease of coincidence rates by 2.5\% per month has been observed, that corresponds to about 1.2\% decrease in OM sensitivity per month (see Fig. 2, right). These measurements agree qualitatively with observations of counting rates, which are continuously monitored in ANTARES (but can be affected by bioluminescence). The present method also allowed to measure the effect of tuning of hardware thresholds, recently performed in ANTARES in order to compensate the sensitivity drop. It was found that the adjustment of pulse-height discrimination thresholds allowed to effectively recover the initial detector sensitivity (see Fig. 2, right). This supports the hypothesis that the sensitivity drop occurs due to a continuous decrease of gain of the photomultiplier tubes. Further measurements with \(^{40}\)K are being performed to better investigate the observed phenomena and maintain high accuracy of the OM sensitivity calibration.

3 Low energy atmospheric muons

An experimental plot of adjacent floor coincidences (A2 coincidences) exhibits a prominent peak shifted toward positive values (see Fig. 3). This is the most basic signal of (downward-going) atmospheric muons detected so far in ANTARES. A Monte Carlo simulation performed using MUPAGE\(^5\) agrees with the observations. The rate of correlated A2 coincidences (A2 rate), which can be defined as the integral under the peak, is directly linked to the atmospheric muon flux. One should note however that a good knowledge of absolute efficiency of the optical modules, as well as their angular response, is needed to accurately convert the A2 rate into the flux. Interestingly, some important information on the angular response curve can be extracted from
the shape of the A2 plot. A detailed study trying to exploit this possibility is currently ongoing. Importantly, the analysis can be repeated for each detector floor separately, thus providing access to the depth-intensity relation of the muon flux. A preliminary analysis (see Fig. 3, right) has demonstrated the effect of muon absorption with depth. A dedicated effort is currently ongoing to develop the necessary corrections for the different sensitivities of the optical modules (and therefore storeys), and further develop this technique.

Clearly this analysis does not allow to determine the properties of each event (e.g. muon zenith angle). However, certain collective properties of the muon flux might well be accessible. Moreover, any systematic errors and inefficiencies of likelihood maximization and event selection procedures are also avoided in this case. In addition, the energy threshold of the coincidence study is much lower, thanks to the use of highly-efficient K40 trigger and very low requirements imposed on the number of signal hits by the analysis itself. Thus the coincidence study can not only provide a redundancy to muon flux measurements but also deliver a lot of useful and complementary information concerning the detector response and atmospheric muon flux.

4 Outlook

A new calibration technique grew up from an academic study of local coincidences in an ANTARES storey. The method represents a unique tool for measurements of optical module sensitivities in situ, and a valuable tool for time calibration in ANTARES. Thanks to this method the effect of OM efficiency drop was precisely measured, and a preliminary time calibration of ANTARES was confirmed in an independent way. Presently the method is used on a regular basis to control the sensitivities of the optical modules. A new approach to atmospheric muon flux studies with ANTARES was suggested, which also relies on a simple coincidence analysis. The method allows to measure the muon flux and, more importantly, its depth dependence. Some preliminary results were presented in this paper, further analysis is ongoing.

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SNEUTRINO COLD DARK MATTER IN EXTENDED MSSM MODELS

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A thorough analysis of sneutrinos as dark matter candidates is performed, in different classes of supersymmetric models, as is typically done for the neutralino dark matter. First in the Minimal Supersymmetric Standard Model, sneutrinos are marginally compatible with existing experimental bounds, including direct detection, provided they compose a subdominant component of dark matter. Then supersymmetric models with the inclusion of right-handed fields and lepton-number violating terms are presented. These models are perfectly viable: they predict sneutrinos which are compatible with the current direct detection sensitivities.

1 Sneutrino in the Minimal Standard Supersymmetric Model

We wish to reconsider in a consistent way sneutrino as a cold relic from the early Universe and study its phenomenology relevant both for Cosmology and for relic–particle detection in low–energy supersymmetric extensions of the Standard Model, which does not (necessarily) invoke mSUGRA relations. We first review the phenomenology of the sneutrino as Cold Dark Matter (CDM) candidate in the case of the Minimal Supersymmetric extension of the Standard Model (MSSM). This model, not very appealing for sneutrino CDM and actually already almost excluded by direct detection searches sets the basis for the extended models described in the next section. In the MSSM, sneutrinos are the scalar partners of the left–handed neutrinos and are described by the usual superpotential and soft breaking terms, leading to the mass-term

\[ m_1 = m_L^2 + \frac{1}{2} m_Z^2 \cos 2\beta, \]

where \( m_L \) is the soft–mass for the left–handed SU(2) doublet \( \tilde{L} \), \( \beta \) is defined as usual from the relation \( \tan \beta = v_2/v_1 \), where \( v_2 \) is the vacuum expectation value of the neutral component of the \( H^2 \) Higgs field and \( m_Z \) is the \( Z \)–boson mass. First of all, we have calculated the sneutrino relic abundance, by taking into account all the relevant annihilation channels and co–annihilation processes which may arise when the sleptons are close in mass to the sneutrinos, as described in Ref. 1. In this minimal MSSM models, the three neutrinos are degenerate in mass: they therefore must be considered jointly in the calculation of the relevant
Figure 1: Left: Sneutrino relic abundance $\Omega h^2$ as a function of the sneutrino mass $m_1$. The horizontal solid and dotted lines delimit the WMAP interval for CDM. Right: Sneutrino–nucleon scattering cross section $\xi \sigma_{\text{nucleon}}^{(\text{scalar})}$ vs. the sneutrino mass $m_1$. The dashed–dotted curve shows the DAMA/NaI region, compatible with the annual modulation effect observed by the experiment. The vertical line denotes the lower bound on the sneutrino mass coming from the invisible $Z$–width. The solid (dashed) curves refer to models with (without) gaugino universality.

processes. An example of sneutrino relic abundance $\Omega h^2$ for the minimal MSSM is plotted in Fig. 1 as a function of the sneutrino mass. The sneutrino relic abundance is typically very small $^2$, much lower than the cosmological range for cold dark matter derived by the WMAP analysis$^3$, which is $0.092 \leq \Omega_{\text{CDM}} h^2 \leq 0.124$. For all the mass range from the experimental lower bound of about $m_Z/2$ up to 600–700 GeV sneutrinos as the LSP are cosmologically acceptable but they are typically underabundant. This means that sneutrinos in the minimal version of MSSM are not good dark matter candidates, except for masses in a narrow range which we determine to be 600–700 GeV. Dark matter direct search, which relies on the possibility to detect the recoil energy of a nucleus due to the elastic scattering of the dark matter particle off the nucleus of a low–background detector, is known to be a strong experimental constraint for sneutrino dark matter. The dependence of the direct detection rate on the DM particle rests into the particle mass and the scattering cross section. For sneutrinos, see Ref.$^1$ for details and references, coherent scattering arises due to $Z$ and Higgs exchange diagrams in the $t$–channel, therefore the relevant cross section on nucleus is given by $\sigma_N = \sigma_{N,Z} + \sigma_{N,H}$. Comparisons with experimental results are more easily and typically performed by using the cross section on a single nucleon $\xi \sigma_{\text{nucleon}}^{(\text{scalar})}$. We have to consider that, whenever the dark matter particle is subdominant in the Universe, also its local density $\rho_0$ in the Galaxy is very likely reduced with respect to the total dark matter density. This means that the dominant component of dark matter is not the sneutrino, but still sneutrinos form a small amount of dark matter and may be eventually detectable. In this case we rescale the local sneutrino abundance by means of the usual factor $\xi = \min(1, \Omega h^2/0.092)$. When compared with the DAMA/NaI annual modulation region$^4$ in Fig. 1 we see that direct detection is indeed a strong constraint on sneutrino dark matter in the minimal MSSM$^2$, but some very specific configurations are still viable and could explain the annual modulation effect.

2 Non minimal supersymmetric models

The models which we will be considering are natural and direct extensions of the MSSM which incorporate at the same time the new physics required to explain two basic problems of astro–particle physics: the origin of neutrino masses and the nature of dark matter. The first class of models (LR models), enlarge the neutrino/sneutrino sector by the inclusion of sterile right–handed superfields $\tilde{N}^i$. The relevant terms in the superpotential and in the soft breaking potential are:

$$W = \epsilon_{ij}(\mu H_1^I \tilde{H}_i^J - Y_{lI} \tilde{l}_i^J \tilde{l}_j^I + Y_{\nu} \tilde{H}_2^I \tilde{N}_i^J)$$

$$V_{\text{soft}} = (M_N^I)^i_j \tilde{N}_i^I \tilde{N}_j^I + (M_N^I)^{i_J} \tilde{N}_i^I \tilde{N}_j^J - [\epsilon_{ij}(A_{\nu}^{ij} H_1^I \tilde{R}_j^J + A_{\nu}^{ij} H_2^I \tilde{N}_j^J) + \text{h.c.}]$$

(1)
where $Y_{\nu}^{IJ}$ is a matrix, which we choose real and diagonal, from which the Dirac mass of neutrinos are obtained $m_D^{IJ} = v_2 Y_{\nu}^{IJ}$, as we do for the matrices $M_N^{ij}$, $\Lambda_\nu^{IJ}$, $M_\nu^0$ and $M_\nu^{ij}$. The parameter which mixes the left– and right–handed sneutrino fields may naturally be of the order of the other entries of the matrix, and induce a sizeable mixing of the lightest sneutrino in terms of left–handed and right–handed fields. The lightest mass eigenstate is therefore $\tilde{\nu}_1 = -\sin \theta \, \tilde{\nu}_L + \cos \theta \, \tilde{N}$, where $\theta$ is the LR mixing angle. Sizeable mixings reduce the coupling to the Z–boson, which couples only to left–handed fields, and therefore have relevant impact on all the sneutrino phenomenology: the lightest sneutrino may be lighter than $m_Z/2$ and the $\tilde{\nu}_1$ annihilation and scattering cross sections which involve Z exchange are reduced, see Ref \cite{1} for details. A supersymmetric model which can accommodate a Majorana mass–term for neutrinos and explain the observed neutrino mass pattern, may be built by adding to the minimal MSSM right–handed fields $\tilde{N}^I$ and allowing for lepton number violating ($L$) terms (Majorana models). The most general form of the superpotential \cite{5,6} and of the soft breaking potential \cite{7}, which accomplishes this conditions is:

$$
W = \epsilon_{ij}(\mu \bar{H}_1^1 \bar{H}_2^2 - Y_{\nu}^{IJ} \bar{H}_1^1 \bar{L}_I^j \bar{R}^j + Y_{\nu}^{IJ} \bar{H}_2^2 \bar{L}_I^j \bar{N}^j) + \frac{1}{2} M^{IJ} \tilde{N}^I \tilde{N}^J \\
V_{\text{soft}} = (M_N^{2})^{IJ} \bar{L}_I^i \bar{L}_I^j + (M_\nu^0)^{IJ} \tilde{N}^I \tilde{N}^J - \frac{1}{2} (m_B^{2})^{IJ} \tilde{N}^I \tilde{N}^J + \epsilon_{ij}(\Lambda_\nu^{IJ} \bar{H}_1^1 \bar{L}_I^j \bar{R}^j + \Lambda_\nu^{IJ} \bar{H}_2^2 \bar{L}_I^j \bar{N}^j) + h.c.
$$

(2)

where we again use the same assumptions of diagonality in flavour space for all the matrices as we already did before. For the $L$ parameters we therefore assume: $M^{IJ} = M \delta^{IJ}$, in order to reduce the number of free parameters. The Dirac mass of the neutrinos is obtained as: $m_D^I = v_2 Y_{\nu}^{IJ}$, while $M^I$ represent a Majorana mass–term for neutrinos. The Dirac–mass parameter is derived by the condition that the neutrino mass is determined by the see–saw mechanism: $m_\nu^I = m_D^I / M^2$. Sneutrinos now are a superpositions of two complex fields: the left–handed field $\nu_L$ and the right–handed field $\tilde{N}$. Since we introduced $L$ terms, it is convenient to work in a basis of CP eigenstates, where the $\tilde{\nu} – Z$ coupling is off diagonal. The non–diagonal nature of the $Z$–coupling leads to important consequences: first of all annihilation processes through Z channel become co–annihilation processes, thus reducing the annihilation cross sections; moreover the elastic scattering off nucleon becomes an inelastic scattering via $t$–channel Z exchange and under certain kinematics condition may be suppressed, leading to a lower value of the direct detection rate of the sneutrino dark matter. The lightest state, which is our dark matter candidate, may now exhibit the non–diagonal nature of the $Z$–coupling with respect of the CP eigenstates, $\tilde{\nu}_- - \tilde{\nu}_+ + \tilde{\nu}_0$, which are also mass eigenstates, and a mixing with the right–handed field $\tilde{N}$: $\tilde{\nu}_i = Z_{i1} \tilde{\nu}_+ + Z_{i2} \tilde{N}_+ + Z_{i3} \tilde{\nu}_- + Z_{i4} \tilde{N}_-$ with $i = 1, 2, 3, 4$.

In LR models sneutrinos may represent the dominant dark matter component for a wide mass range. The most relevant new feature is that for the full supersymmetric scan, the mass range allowed by the cosmological constraints is enlarged up to 800 GeV, and all the mass interval above the Z–pole may lead to strongly subdominant sneutrinos. From Fig. 2 we can conclude that after all experimental and theoretical constraints are imposed, sneutrino dark matter is perfectly viable, both as a dominant and as a subdominant component, for the whole mass range $15 \, \text{GeV} \lesssim m_1 \lesssim 800 \, \text{GeV}$. The lower limit of 15 GeV represents therefore a cosmological bound on the sneutrino mass in LR models, under the assumption of $R$–parity conservation. The sneutrino–nucleon cross section is shown in the right top in Fig. 2. Only points which are accepted by the cosmological constraint are shown. We see that the presence of the mixing with the right–handed $\tilde{N}$ fields opens up the possibility to have viable sneutrino cold dark matter. A fraction of the configurations are excluded by direct detection, but now, contrary to the minimal MSSM case, a large portion of the supersymmetric parameter space is compatible with the direct detection bound, both for cosmologically dominant (red crosses) and subdominant (blue points) sneutrinos. The occurrence of sneutrinos which are not in conflict.
with direct detection limits and, at the same time, are the dominant dark matter component, is a very interesting feature of this class of models. The relic abundance of the Majorana models at low–mass scale is shown on the left bottom in Fig. 2. It is remarkable that in the whole mass range from 5 GeV to 1 TeV sneutrinos can explain the required amount of CDM in the Universe. For the same mass range, sneutrinos may as well be a subdominant component. Direct detection is shown on the right bottom in Fig. 2. We see that three different populations arise: configurations on the upper right are clearly excluded by direct detection searches. Most of them refer to subdominant sneutrinos. Configurations on the lower right part of the plot are allowed but well below current direct detection sensitivity. Configurations on the center and left part of the plot all fall inside the current sensitivity region: a large fraction are cosmologically dominant and could explain the annual modulation effect observed by the DAMA/NaI experiment.

Acknowledgments
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References
EFFECTS OF LIGHTEST NEUTRINO MASS IN LEPTOGENESIS

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The effects of the lightest neutrino mass in “flavoured” leptogenesis are investigated in the case when the CP-violation necessary for the generation of the baryon asymmetry of the Universe is due exclusively to the Dirac and/or Majorana phases in the neutrino mixing matrix $U$. The type I see-saw scenario with three heavy right-handed Majorana neutrinos having hierarchical spectrum is considered. The “orthogonal” parametrisation of the matrix of neutrino Yukawa couplings, which involves a complex orthogonal matrix $R$, is employed. Results for light neutrino mass spectrum with normal and inverted ordering (hierarchy) are obtained. It is shown, in particular, that if the matrix $R$ is real and CP-conserving and the lightest neutrino mass $m_3$ in the case of inverted hierarchical spectrum lies the interval $5 \times 10^{-4} \text{ eV} \lesssim m_3 \lesssim 7 \times 10^{-3} \text{ eV}$, the predicted baryon asymmetry can be larger by a factor of $\sim 100$ than the asymmetry corresponding to negligible $m_3 \equiv 0$. As consequence, we can have successful thermal leptogenesis for $5 \times 10^{-6} \text{ eV} \lesssim m_3 \lesssim 5 \times 10^{-2} \text{ eV}$ even if $R$ is real and the only source of CP-violation in leptogenesis is the Majorana and/or Dirac phase(s) in the neutrino mixing matrix.

1 Introduction

We investigate the effects of the lightest neutrino mass in thermal leptogenesis\textsuperscript{1, 2} where lepton flavor effects\textsuperscript{3–8} play an important role in the generation of the observed baryon asymmetry of the Universe and the CP-violation required for the baryogenesis mechanism to work is due exclusively to the Dirac and/or Majorana CP-violating phases in the Pontecorvo-Maki-Nakagawa-Sakata (PMNS)\textsuperscript{9} neutrino mixing matrix. A detailed analysis of these effects has been performed in reference\textsuperscript{10}.

The minimal scheme in which leptogenesis can be implemented is the non-supersymmetric version of the type I see-saw\textsuperscript{11} model with two or three heavy right-handed (RH) Majorana neutrinos. Taking into account the lepton flavour effects in leptogenesis it was shown\textsuperscript{12} (see also\textsuperscript{4, 13, 14}) that if the heavy Majorana neutrinos have a hierarchical spectrum, the observed baryon asymmetry $Y_B$ can be produced even if the only source of CP-violation is the Majorana and/or Dirac phase(s) in the PMNS matrix $U_{PMNS} \equiv U$. In this case the predicted value of


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the baryon asymmetry depends explicitly (i.e. directly) on \( U \) and on the CP-violating phases in \( U \). The results quoted above were demonstrated to hold both for normal hierarchical (NH) and inverted hierarchical (IH) spectrum of masses of the light Majorana neutrinos. In both these cases they were obtained for negligible lightest neutrino mass and CP-conserving elements of the orthogonal matrix \( R \), present in the “orthogonal” parametrisation of the matrix of neutrino Yukawa couplings. The CP-invariance constraints imply that the matrix \( R \) could conserve the CP-symmetry if its elements are real or purely imaginary. We remark that for a CP-conserving matrix \( R \) and at temperatures \( T \sim M_1 \gtrsim 10^{12} \text{ GeV} \), the lepton flavours are indistinguishible (one flavour approximation) and the total CP asymmetry is always zero. In this case no baryon asymmetry is produced. One can prove that, for NH spectrum and negligible lightest neutrino mass \( m_1 \) successful thermal leptogenesis can be realised for a real matrix \( R \). In contrast, in the case of IH spectrum and negligible lightest neutrino mass \( m_3 \), the requisite baryon asymmetry was found to be produced for CP-conserving matrix \( R \) only if certain elements of \( R \) are purely imaginary: for real \( R \) the baryon asymmetry \( Y_B \) is strongly suppressed and leptogenesis cannot be successful for \( M_1 \lesssim 10^{12} \text{ GeV} \) (i.e. in the regime in which the lepton flavour effects are significant).

In the present work we have analysed the effects of the lightest neutrino mass on “flavoured” (thermal) leptogenesis. We considered the case when the CP-violation necessary for the generation of the observed baryon asymmetry of the Universe is due exclusively to the Dirac and/or Majorana CP-violating phases in the PMNS matrix \( U \). Our study is performed within the simplest type I see-saw scenario with three heavy RH Majorana neutrinos \( N_j, j = 1, 2, 3 \). The latter are assumed to have a hierarchical mass spectrum, \( M_1 \ll M_{2,3} \). As a consequence, the generated baryon asymmetry \( Y_B \) depends linearly on the mass of \( N_1 \), \( M_1 \), and on the elements \( R_{1j} \) of the matrix \( R \), \( j = 1, 2, 3 \), present in the neutrino Yukawa couplings of \( N_1 \). Throughout the present study we employ the “orthogonal” parameterisation of the matrix of neutrino Yukawa couplings. As was already mentioned previously, this parameterisation involves an orthogonal matrix \( R \), \( R^T R = RR^T = 1 \). Although, in general, the matrix \( R \) can be complex, i.e. CP-violating, in the present work we are primarily interested in the possibility that \( R \) conserves the CP-symmetry. We consider the two types of light neutrino mass spectrum allowed by the data: i) with normal ordering \( (\Delta m^2_\text{A} > 0) \), \( m_1 < m_2 < m_3 \); and ii) with inverted ordering \( (\Delta m^2_\text{A} < 0) \), \( m_3 < m_1 < m_2 \). The case of inverted hierarchical (IH) spectrum and real (and CP-conserving) matrix \( R \) is investigated in detail.

Our analysis is performed for negligible renormalisation group (RG) running of \( m_j \) and of the parameters in the PMNS matrix \( U \) from \( M_Z \) to \( M_1 \). This possibility is realised for sufficiently small values of the lightest neutrino mass \( \min(m_j) \) e.g., for \( \min(m_j) \lesssim 0.10 \text{ eV} \). The latter condition is fulfilled for the NH and IH neutrino mass spectra, as well as for spectrum with partial hierarchy. Under the indicated condition \( m_j \), and correspondingly \( \Delta m^2_\text{A} \) and \( \Delta m^2_\text{S} \), and \( U \) can be taken at the scale \( \sim M_Z \), at which the neutrino mixing parameters are measured.

2 Light Neutrino Mass Spectrum with Inverted Ordering and Real \( R_{1j} \)

The case of inverted hierarchical (IH) neutrino mass spectrum, \( m_3 \ll m_1 < m_2 \), \( m_{1,2} \equiv \sqrt{|\Delta m^2_\text{S}|} \), is of particular interest since, as was already mentioned, for real \( R_{1j} \), \( j = 1, 2, 3 \), IH spectrum and negligible lightest neutrino mass \( m_3 \approx 0 \), it is impossible to generate the observed baryon asymmetry \( Y_B \approx 8.6 \times 10^{-11} \) in the regime of “flavoured” leptogenesis, i.e. for \( M_1 \lesssim 10^{12} \text{ GeV} \), if the only source of CP violation are the Majorana and/or Dirac phases in the

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\(^{15}\) The case in which CP-violation arises from the combined effect between the “low energy” Majorana and/or Dirac phases in \( U_{\text{PMNS}} \) and the “high energy” CP-violating phases in a complex orthogonal matrix \( R \), in thermal “flavoured” leptogenesis scenario, has recently been addressed.\(^{16}\)
PMNS matrix. It can be proved that for \( m_3 = 0 \) and \( R_{13} = 0 \), the resulting baryon asymmetry is always suppressed by the factor \( \Delta m^2_\odot / (2 \Delta m^2_A) \approx 1.6 \times 10^{-2} \). We analyse the generation of the baryon asymmetry \( Y_B \) for real \( R_{1j} \), \( j = 1, 2, 3 \), when \( m_3 \) is non-negligible. We assume that \( Y_B \) is produced in the two-flavour regime, \( 10^9 \text{ GeV} < M_1 \sim 10^{12} \text{ GeV} \).

In Fig. 1 we show the correlated values of \( M_1 \) and \( m_3 \) for which one can have successful leptogenesis in the case of neutrino mass spectrum with inverted ordering and CP-violation due to the Majorana and Dirac phases in \( U_{\text{PMNS}} \). The result was obtained by performing, for given \( m_3 \) from the interval \( 10^{-10} \text{ eV} \leq m_3 \leq 0.05 \text{ eV} \), a thorough scan of the relevant parameter space searching for possible enhancement or suppression of the baryon asymmetry with respect to that found for \( m_3 = 0 \). The real elements of the \( R \)-matrix of interest, \( R_{1j} \), \( j = 1, 2, 3 \), were allowed to vary in their full ranges determined by the condition of orthogonality of the matrix \( R \): \( R^2_{11} + R^2_{12} + R^2_{13} = 1 \). The Majorana phases \( \alpha_{21,31} \) were varied in the interval \([0, 2\pi] \). The calculations were performed for three values of the CHOOZ angle \( \theta_{13} \), corresponding to \( \sin \theta_{13} = 0; 0.1; 0.2 \). In the cases of \( \sin \theta_{13} \neq 0 \), the Dirac phase \( \delta \) was allowed to take values in the interval \([0, 2\pi] \). The heavy Majorana neutrino mass \( M_1 \) was varied in the interval \( 10^9 \text{ GeV} \leq M_1 \leq 10^{12} \text{ GeV} \). For given \( m_3 \), the minimal value of the mass \( M_1 \), for which the leptogenesis is successful, generating \( |Y_B| \approx 8.6 \times 10^{-11} \), was obtained for the values of the other parameters which maximise \( |Y_B| \). We have found that in the case of IH spectrum with non-negligible \( m_3 \), \( m_3 \ll \sqrt{|\Delta m^2_\Lambda|} \), the generated baryon asymmetry \( |Y_B| \) can be strongly enhanced in comparison with the asymmetry \( |Y_B| \) produced if \( m_3 \approx 0 \). The enhancement can be by a factor of \( \sim 100 \), or even by a larger factor. As a consequence, one can have successful leptogenesis for IH spectrum with \( m_3 \gtrsim 5 \times 10^{-6} \text{ eV} \) even if the elements \( R_{1j} \) of \( R \) are real and the requisite CP-violation is provided by the Majorana or Dirac phase(s) in the PMNS matrix. As a consequence, successful thermal leptogenesis is realised for \( 5 \times 10^{-6} \text{ eV} \lesssim m_3 \lesssim 5 \times 10^{-2} \text{ eV} \). The results of our analysis show that for Majorana CP-violation from \( U_{\text{PMNS}} \), successful
leptogenesis can be obtained for $M_1 \gtrsim 3.0 \times 10^{10} \text{ GeV}$. A somewhat larger values of $M_1$ are typically required if the CP-violation is due to the Dirac phase $\delta$: $M_1 \gtrsim 10^{11} \text{ GeV}$. The requirement of successful “flavoured” leptogenesis in the latter case leads to the following lower limits on $|\sin \theta_{13} \sin \delta|$, and thus on $\sin \theta_{13}$ and on the rephasing invariant $J_{CP}$ which controls the magnitude of CP-violation effects in neutrino oscillations: $|\sin \theta_{13} \sin \delta|, \sin \theta_{13} \gtrsim (0.04 - 0.09)$, $|J_{CP}| \gtrsim (0.009 - 0.020)$, where the precise value of the limit within the intervals given depends on the $\text{sgn}(R_{11} R_{13})$ (or $\text{sgn}(R_{12} R_{13}))$ and on $\sin^2 \theta_{23}$.

The results we have obtained for light neutrino mass spectrum with normal ordering, $m_1 < m_2 < m_3$, can vary significantly if one of the elements $R_{1j}$ is equal to zero. In particular, if $R_{11} \approx 0$, we did not find any significant enhancement of the baryon asymmetry $|Y_B|$, generated within “flavoured” leptogenesis scenario with real matrix $R$ and CP-violation provided by the neutrino mixing matrix $U_{PMNS}$, when the lightest neutrino mass was varied in the interval $10^{-10} \text{ eV} \leq m_1 \leq 0.05 \text{ eV}$. If, however, $R_{12} \approx 0$, the dependence of $|Y_B|$ on $m_1$ exhibits qualitatively the same features as the dependence of $|Y_B|$ on $m_3$ in the case of neutrino mass spectrum with inverted ordering (hierarchy), although $\text{max}(|Y_B|)$ is somewhat smaller than in the corresponding IH spectrum cases. As a consequence, it is possible to reproduce the observed value of $Y_B$ if the CP-violation is due to the Majorana phase(s) in $U_{PMNS}$ provided $M_1 \gtrsim 5.3 \times 10^{10} \text{ GeV}$.

The analysis we have performed shows that within the thermal “flavoured” leptogenesis scenario, the value of the lightest neutrino mass can have non negligible effects on the magnitude of the baryon asymmetry of the Universe in the cases of light neutrino mass spectrum with inverted and normal ordering (hierarchy). In particular, as regards the IH spectrum, one can have an enhancement of the baryon asymmetry by a factor of $\sim 100$ with respect to the value corresponding to $m_3 \approx 0$, thus allowing for the generation of a matter-antimatter asymmetry compatible with the experimental observation.

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References

SEARCH FOR NEUTRINOLESS DOUBLE BETA DECAY WITH NEMO-3 USING $^{150}$ND

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The NEMO-3 detector looks for neutrinoless double beta decay from 10 kg of enriched $\beta\beta$ emitters. Using 440 days of data, background measurements give activities from each contribution from Neodinium foil. During the same period of data acquisition, preliminary study gives the half-life of the $2\nu\beta\beta$ process: $T_{1/2}(2\nu) = (9.75 \pm 0.35(stat.) \pm 0.85(syst.)) \times 10^{18}$ years. Up to now, no evidence of lepton number violation has been observed. A limit on the half-life of the neutrinoless process has been obtained: $T_{1/2}(0\nu) > 8.0 \times 10^{21}$ y at 90% CL.

1 Neutrinoless double beta decay

The goal of the NEMO-3 (Neutrino Ettore Majorana Observatory) detector is to search for evidence of lepton number violation in neutrinoless double beta decay ($0\nu\beta\beta$):

$$ (A, Z) \rightarrow (A, Z + 2) + 2e^- \quad (1) $$

The decay may be explained by the exchange of a light massive Majorana neutrino. Also (V+A) coupling or Majoron emission could describe the $0\nu\beta\beta$ decay. These processes, forbidden in the framework of the standard model, are the only experimental opportunity to prove the Majorana nature of the neutrino.

The half-life of the $0\nu\beta\beta$ process by light Majorana neutrino exchange can be written as:

$$ (T_{1/2}^{0\nu})^{-1} = G^{0\nu}(Z, Q_{\beta\beta})|M_{GT}^{0\nu}|^2 - \left(\frac{g_V}{g_A}\right)^2 M_F^{0\nu} |\langle m_{\beta\beta}\rangle|^2 \quad (2) $$

Where $G^{0\nu}(Z, Q_{\beta\beta})$ is the phase space factor proportional to the $Q_{\beta\beta}^2$, $M_{GT}^{0\nu}$ are the nuclear matrix elements (NME) and $\langle m_{\beta\beta}\rangle$ is the effective mass of neutrino given by:

$$ < m_{\nu_{\beta\beta}} >= |\cos^2 \theta_{13} |m_1| \cos^2 \theta_{12} + |m_2| e^{2i\phi_1} \sin^2 \theta_{12} | + |m_3| e^{2i\phi_2} \delta \sin^2 \theta_{12}| \quad (3) $$
<table>
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</tr>
<tr>
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</tr>
<tr>
<td>$^{82}\text{Se} \rightarrow ^{82}\text{Kr}$</td>
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<td>932</td>
</tr>
<tr>
<td>$^{100}\text{Mo} \rightarrow ^{100}\text{Ru}$</td>
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<tr>
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<td>$^{150}\text{Nd} \rightarrow ^{150}\text{Sm}$</td>
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<td>37</td>
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<tr>
<td>$^{48}\text{Ca} \rightarrow ^{48}\text{Ti}$</td>
<td>4271±4</td>
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</tr>
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Table 1: $\beta\beta$ emitters in the NEMO-3 detector.

The measurement of the $0\nu\beta\beta$ process half-life, given the NME, determines the absolute Majorana neutrino scale.

2 Searching for golden events in the NEMO-3 detector

The NEMO-3 detector is designed to detect and identify neutrinoless double beta decays using 10 kg of enriched $\beta\beta$ emitters (Cf. table 1).

Isotopes are spread on thin foils at the center of a drift chamber formed by 6180 open Geiger cells to reconstruct charged particles tracks. An ultra pure calorimeter made of 1940 plastic scintillators encloses the drift chamber to measure energy and time of flight of particles coming from the source foils. NEMO-3 is able to identify $0\nu\beta\beta$ candidate events (golden events) with some dedicated criteria.

In figure 1-(b) a partial top view of the detector shows a $0\nu\beta\beta$-like event in NEMO-3 detector. There are two $\beta$ tracks with negative curvature with a common vertex on the source foil. The time of flight is compatible with an event coming from the foil ($\delta t_{TOF}\approx0$). The total energy deposit in the calorimeter is equal to 3222 keV ($Q_{\beta\beta}=3367$ keV for $^{150}\text{Nd}$ and energy resolution : 14 % at 1 MeV). This event from 37 g of $^{150}\text{Nd}$ is a typical golden event ($\Delta L=2$ process).

3 Analysis and results

The golden event showed in figure 1-(b) is simulated from possible background contribution. Each process which is able to produce two electrons with a total energy around $Q_{\beta\beta}$ is considered as a dangerous background for the neutrinoless double beta decay investigation. For example, a $\beta\beta$-like event can be generated by a $\beta-\gamma$ emitter, the second electron being produced by Compton or photoelectric effect.

Typically the most important background comes from natural radioactive chain ($^{238}\text{U}$ and $^{232}\text{Th}$). That brings to light the crucial importance of the extreme radio-purity one needs to achieve in order to reach a sensitivity to the half-life of up to $10^{24-25}$ years. The NEMO-3 detector is hosted in the Frejus underground laboratory under 4800 m.w.e. to reduce the cosmic ray flux. More the allowed process ($2\nu\beta\beta$) is the ultimate background which cannot be suppressed due to the energy resolution (Cf. figure 1-(a)).

3.1 Background studies

The so-called tracko-calo technology used in NEMO-3 is able to measure and control the background sources using dedicated channels and to localize possible contaminations. The channels of interest are defined by the decay scheme from natural radioactive isotopes (Cf figure 2).
(a) Theoretical spectrum of the total energy for $\beta\beta$ processes in $^{150}$Nd. The standard model allowed process ($2\nu\beta\beta$) is observed in NEMO-3. Due to the energy resolution, the continuous energy spectrum is a significant source of background for the $0\nu\beta\beta$ search.

(b) Top view of a typical $0\nu\beta\beta$ event in NEMO-3 from $^{150}$Nd.

Figure 1:

- $^{208}$Tl from the $^{232}$Th chain is a $\beta$ emitter with $Q_\beta$ at 4.99 MeV. The most efficient signature to study this background consists in the measurement of one electron and several photons from deexcitation from the daughter nucleus. The corresponding channel is defined by an electron in coincidence with at least $\gamma$ rays. The $^{208}$Tl activity in $^{150}$Nd foil using such techniques has been measured at 17 $\mu$Bq.

- $^{214}$Bi from $^{238}$U chain is a $\beta$ emitter at $Q_\beta=3.27$ MeV followed by the $\alpha$ ($T_{1/2}=164$ $\mu$s) from $^{214}$Po decay in $^{210}$Pb. Dedicated electronics have been developed to identify the delayed $\alpha$ particle in the tracking chamber. The main channel to measure the BiPo cascade contamination is a prompt $\beta$ decay in coincidence with a delayed $\alpha$. $^{214}$Bi contamination has been measured on the Nd foil at the level of 90 $\mu$Bq.

These results are compatible with independent measurements using HPGe detector.

Figure 2: Visualization of some channels studied to measure background contamination in NEMO-3 detector. a) Typical event in the $e\gamma\gamma$ channel used to identify the $^{208}$Tl contamination. b) Typical event in the $e\gamma\alpha$ channel used in case of $^{214}$Bi contamination.
3.2 **Double beta decay analysis**

The total energy spectra from two internal electron channel is split into two main parts. The [0-3] MeV region is used to measure the $2\nu\beta\beta$ process and the [3-3.6] MeV energy range is used to search for $0\nu\beta\beta$ signal. Contaminations of this channel are $^{214}$Bi and $^{208}$Tl (and other isotopes from natural radioactive chains). The contamination level has been studied to determine the background contribution in the $\beta\beta$ candidate signal from the Neodinium foil.

Preliminary results for a total time acquisition of 10560 hours give 756 $\beta\beta$-like events with signal over background ratio $S/B = 3.4$ in the [400-3600] keV full spectra. The half-life of the $2\nu\beta\beta$ process is:

$$T_{1/2}(2\nu) = (9.75 \pm 0.35(\text{stat.}) \pm 0.85(\text{syst.})) \times 10^{18} \text{ years}$$ (4)

![Figure 3: Total energy spectra in the two internal electron channel. Dots are real data. Blue/red histograms represent the expected contamination from external/internal background. The blue line is the $2\nu\beta\beta$ contribution. The green histogram represents the sum of all contributions mentioned previously.](image)

There is no expected background event in the [3-3.6] MeV energy range. No $0\nu\beta\beta$ candidate event have been observed in the 10560 hours data sample from $^{150}$Nd foil. The limit obtained is $T_{1/2}(0\nu) > 8.0 \times 10^{21} \text{ y}$ at 90% CL for the light majorana exchange hypothesis. Up to now there is no lepton number violation evidence using $^{150}$Nd data in the NEMO-3 detector.

**References**

We study the possibility of identifying dark matter properties from direct (XENON100) and indirect (GLAST) detection experiments. In the same way, we examine the perspectives given by the next generation of colliders (ILC). All this analysis is done following a model-independent approach. We have shown that the three detection techniques can act in a highly complementary way, whereas direct detection experiments will probe efficiently light WIMPs, given a positive detection (at the 10% level for $m_\chi \sim 50$ GeV), GLAST will be able to confirm and even increase the precision in the case of a NFW profile, for a WIMP-nucleon cross-section $\sigma_{\chi-p} < 10^{-8}$ pb. However, for heavier WIMP ($\sim 175$ GeV), the ILC will lead the reconstruction of the mass.

1 Direct detection

Dark Matter (DM) direct detection experiments measure the elastic collisions between WIMPs and target nuclei in a detector, as a function of the recoil energy $E_r$. The detection rate depends on the density $\rho_0 \simeq 0.3$ GeV cm$^{-3}$ and velocity distribution $f(v_\chi)$ of WIMPs near the Earth (a Maxwellian halo will be considered). The differential rate per unit detector mass and per unit of time can be written as:

$$\frac{dN}{dE_r} = \frac{\sigma_{\chi-N} \rho_0}{2m_r^2 m_\chi} F(E_r)^2 \int_{v_{min}(E_r)}^{\infty} \frac{f(v_\chi)}{v_\chi} dv_\chi,$$

where $\sigma_{\chi-N}$ is the WIMP-nucleus scattering cross section, $m_\chi$ the WIMP mass and $m_r$ is the WIMP-nucleon reduced mass. $F(E_r)$ is the nucleus form factor; we assume it has the Woods-Saxon form.

The XENON$^1$ experiment aims at the direct detection of DM via its elastic scattering off xenon nuclei. In this study we will consider the case of a 100 kg Xenon-like experiment and 3
years taking data. We consider 7 energy bins between 4 and 30 keV. For this experiment we took a zero background scenario; of course a more detailed analysis could take into account non-zero background simulating the detector and in particular the neutron spectrum. So, in that sense our results will be optimistic.

One option to discriminate between a DM signal and the background is to use the $\chi^2$ method. Let us call $N_{\text{sign}}$ the signal, $N_{\text{bkg}}$ the background and $N_{\text{tot}} = N_{\text{sign}} + N_{\text{bkg}}$ the total signal measured by the detector. For an energy range divided into $n$ bins the $\chi^2$ is defined as:

$$\chi^2 = \sum_{i=1}^{n} \frac{(N_{\text{tot}} - N_{\text{bkg}})^2}{\sigma_i^2}.$$  \hspace{1cm} (2)

We assume a Gaussian error $\sigma_i \equiv \sqrt{N_{\text{tot}} \cdot T}$ on the measurement, where $M$ is the detector mass and $T$ the exposure time.

2 Indirect detection

The spectrum of gamma–rays generated in dark matter annihilations and coming from the galactic center can be written as

$$\Phi_{\gamma}(E_{\gamma}) = \sum_i \frac{dN_i}{dE_{\gamma}} \cdot Br_i \cdot \langle \sigma v \rangle \cdot \frac{1}{8\pi m_{\chi}^2} \cdot \int_{\text{line of sight}} \rho^2 \, dl,$$  \hspace{1cm} (3)

where the discrete sum is over all dark matter annihilation channels, $dN_i/dE_{\gamma}$ is the differential gamma–ray yield, $\langle \sigma v \rangle$ is the annihilation cross-section averaged over the WIMPs’ relative velocity distribution and $Br_i$ is the branching ratio of annihilation into the $i$-th final state. It is possible to concentrate ourselves on a process which gives 100% annihilation into $WW$ pairs, as this choice will not influence significantly the result of the study\(^2\). The dark matter density $\rho$ is usually parametrized as

$$\rho(r) = \frac{\rho_0}{(r/R)^\alpha \left[1 + (r/R)^\alpha\right]^{(\beta-\gamma)/\alpha}}.$$  \hspace{1cm} (4)

We assume a NFW profile with $\alpha = \gamma = 1$, $\beta = 3$ and $R = 20$ kpc, producing a profile with a behavior $\rho(r) \propto r^{-\gamma}$ in the inner region of the galaxy.

The gamma-ray telescope GLAST\(^3\) will perform an all-sky survey covering an energy range $1 - 300$ GeV. We will consider a GLAST-like experience with an effective area and angular resolution on the order of $10^{4}$ cm\(^2\) and $0.1^\circ \times 0.1^\circ$ ($\Delta\Omega \approx 10^{-5}$ sr) respectively, who will be able to point and analyze the inner centre of our galaxy. We consider also 3 years of effective data acquisition experiment.

For this experiment, the background can be modeled by interpolating\(^2\) the gamma-ray spectrum measured by HESS\(^4\) (for $E_{\gamma} > 160$ GeV) and EGRET\(^5\) (for $E_{\gamma} < 10$ GeV) missions.

3 Colliders

Recently an approach was proposed by A. Birkedal et al\(^6\) which allows to perform a model-independent study of WIMP properties at lepton colliders. Since the known abundance of DM gives specific values for the DM annihilation cross section, one might hope this cross section can be translated into a rate for a measurable process at a collider.

The starting point is to relate total annihilation cross section to the cross section into $e^+e^-$ pairs

$$\kappa_e \equiv \sigma(\chi\chi \rightarrow e^+e^-)/\sigma(\chi\chi \rightarrow all).$$  \hspace{1cm} (5)
Figure 1: Comparison between a 100 kg XENON-like experiments (dotted line) with $\sigma_{\chi p} = 10^{-7}$ pb, GLAST (dashed line) in the case of an NFW halo profile with $\langle \sigma v \rangle = 3 \cdot 10^{-26}$ cm$^3$s$^{-1}$, and unpolarized ILC sensitivity (solid line) at 2$\sigma$ of confidence level, for different WIMP masses $m_\chi = 100$ and 175 GeV, and $\kappa_e = 0.3$.

Then we can use the detailed balancing equation to relate $\sigma(\chi \chi \rightarrow e^+ e^-)$ to $\sigma(e^+ e^- \rightarrow \chi \chi)$, for non-relativistic WIMPs. But this kind of process containing only WIMPs in the final state is not visible in a collider since they manifest themselves just as missing energy. However this process can be correlated to the radiative WIMP pair-production $\sigma(e^+ e^- \rightarrow \gamma \chi \chi)$ using the collinear factorization. This approach is valid for photons which are either soft or collinear with respect to the colliding beams. The accuracy of the approximation outside the previous region has been discussed with the conclusion that the approach works quite well.

So, starting from the total annihilation cross section $\sigma_{an}$ we can compute $\sigma(e^+ e^- \rightarrow \gamma \chi \chi)$

$$\frac{d\sigma(e^+ e^- \rightarrow \chi \chi \gamma)}{dx \, dcost} \approx \frac{\alpha \kappa_e \sigma_{an}}{16 \pi} \frac{1}{x} \frac{1}{\sin^2 \theta} 2^{2J_0} \left( \frac{2 S_\chi + 1}{2} \right) \left( 1 - \frac{4 m_\chi^2}{(1 - x) s} \right)^{1/2 + J_0} \cdot (6)$$

Here $x \equiv 2 E_\gamma / \sqrt{s}$, $\theta$ is the angle between the photon and the incoming beam, $S_\chi$ and $J_0$ are the spin of the WIMP and the dominant value of the angular momentum in the velocity expansion for $\langle \sigma v \rangle$.

We place ourselves in the framework of the ILC project with a center-of-mass energy of $\sqrt{s} = 500$ GeV and an integrated luminosity of 500 fb$^{-1}$. For this process with only a single photon detected, the main background in the standard model is radiative neutrino production.$^2$

4 Complementarity

Recently, several works have studied the determination of the WIMP mass for the case of direct$^7$ and indirect$^8$ detection experiments. Furthermore, Drees and Shan$^9$ showed that one can increase such a precision with a combined analysis of two experiments of direct detection.

In figure 1 we compare the precision levels for direct and indirect detection experiments, along with the corresponding results of the method we followed for the ILC for $\kappa_e = 0.3$ and two cases of WIMP masses $m_\chi = 100$ (left panel) and 175 GeV (right panel). All the results are plotted for a 2$\sigma$ confidence level. The green-dashed lines correspond to the results for a GLAST-like experiment assuming a NFW halo profile. The total annihilation cross-section has been taken to be $\langle \sigma v \rangle = 3 \cdot 10^{-26}$ cm$^3$s$^{-1}$. The red-plain line represents the result for an ILC-like collider with non-polarized beams. The blue-dotted line corresponds to a 100 kg XENON-like experiment, assuming a WIMP-nucleus scattering cross-section of $10^{-7}$ pb. All the parameter space points that lie within the marked regions cannot be discriminated by the corresponding experiments. It is pertinent to study the complementarity between the three experiences listed above firstly
because the mass reconstruction yields comparable results, hence a combination of these data can substantially improve the final result. Secondly, because we can probe different regions in the parameter space. For the case of a 100 GeV WIMP, a GLAST- or an ILC-like experiment alone can provide a limited precision for the WIMP mass (~60%). Combined measurements can dramatically increase the precision, reaching an accuracy of ~25%. If we additionally include direct detection measurement, we can reach a precision of the order of ~9%.

In table 1 we show the precision expected for several dark DM masses. A light WIMP mass (~50 GeV) can be reconstructed by both direct and indirect DM experiments with a high level of precision; however for the ILC the model independent procedure fails because of the relativistic nature of the WIMP. On the contrary, the ILC will be particularly efficient to measure a WIMP with a mass of about 175 GeV. Concerning a 500 GeV WIMP, only a loose lower bound could be extracted from direct and indirect experiments. In this case the ILC will not be kinematically able to produce so heavy WIMPs.

Acknowledgments

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References

We study the decay rates of the $\mu \rightarrow e \gamma$ and $\tau \rightarrow \ell \gamma$ transitions in the framework of the type III seesaw model, where fermionic triplets are exchanged to generate neutrino masses. We show that the observation of one of those decays in planned experiments would contradict bounds arising from present experimental limits on the $\mu \rightarrow eee$ and $\tau \rightarrow 3\ell$ decay rates, and therefore imply that there exist other sources of lepton flavour violation than those associated to triplet of fermions.

1 Introduction

The Standard Model (SM) has the unique property of conserving flavour in the leptonic sector. However, since the experimental discovery of neutrino oscillations, we know that lepton flavour is violated in the neutrino sector. Neutrinos mass naturally arises within the framework of the seesaw mechanism (via the exchange of heavy fields). In such models flavour violating rare leptonic decays such as $\mu \rightarrow e \gamma$ and $\tau \rightarrow \ell \gamma$ are expected to be relevant. These decays have already been studied in type I\textsuperscript{1} and type II\textsuperscript{2} seesaw models. In the following, we study these decays in the framework of the type III seesaw model \textsuperscript{3,4}, where heavy triplets of fermions are exchanged.

2 The type-III Seesaw model

The type-III seesaw model consists in adding $SU(2)_L$ triplets of fermions $\Sigma$, with zero hypercharge, to the SM. At least two triplets are needed to account for the observation of neutrino masses, but in fact only one is sufficient to generate non-vanishing $\ell_1 \rightarrow \ell_2 \gamma$ rate. In the following we will not specify the number of triplets. The heavy fermions are in the adjoint representation
of the $SU(2)_L$ group and have a gauge invariant Majorana mass term. The Lagrangian of its interactions reads:

$$L = \text{Tr} \left[ \Sigma^i D^i \Sigma \right] - \frac{1}{2} \text{Tr} \left[ \Sigma M \Sigma^c + \Sigma^c M^\dagger \Sigma \right] - \bar{\phi} \Sigma \sqrt{2} Y \Sigma L - \bar{L} \sqrt{2} Y \Sigma^c \phi,$$

where $L \equiv (l, \nu)^T$, $\phi \equiv (\phi^+, \phi^0)^T \equiv (\phi^+, (v + H + i\eta)/\sqrt{2})^T$, $\bar{\phi} = i\tau_2 \phi^*$, $\Sigma^c \equiv C \Sigma^T$ and with, for each fermionic triplet,

$$\Sigma^c = \begin{pmatrix} \Sigma^0/\sqrt{2} & \Sigma^+ \\ -\Sigma^0/\sqrt{2} & -\Sigma^- \end{pmatrix}, \quad \Sigma^c = \begin{pmatrix} \Sigma^0/\sqrt{2} & \Sigma^- \\ -\Sigma^0/\sqrt{2} & -\Sigma^+ \end{pmatrix}.$$

After electroweak symmetry breaking, the neutrino mass matrix is given by:

$$m_\nu = -\frac{v^2}{2} Y \Sigma M \Sigma^c Y \Sigma.$$

The new Yukawa couplings are the source of mixing between the light leptons and the heavy fermions, which, combined with the presence of the Majorana mass term, allow lepton flavour violating processes. Thus, the study of these processes will enable to derive some bounds on the new couplings: $Y_\Sigma$ and $M_\Sigma$.

### 3 Flavour changing radiative leptonic decays

We briefly describe the main steps of the calculation of the $\mu \to e\gamma$ rate. The $\tau$ decay rates will straightforwardly follow. The on-shell transition $\mu \to e\gamma$ is a magnetic transition, and it can be written, in the limit $m_e \to 0$, as:

$$T (\mu \to e\gamma) = A \times \overline{\nu}_e (p-q) \left[ i q^\mu \varepsilon^\lambda \sigma_{\lambda\nu} (1 + \gamma_5) \right] u_\mu (p),$$

where $\varepsilon$ is the polarization of the photon, $p_\mu$ the momentum of the incoming muon, $q_\mu$ the momentum of the outgoing photon and $\sigma_{\mu\nu} = \frac{i}{2} [\gamma_\mu, \gamma_\nu]$. The fourteen diagrams contributing to these decays are shown in Fig. 1. The details of the calculation can be found in the related paper\textsuperscript{5}. In the limit $M_\Sigma \gg M_W$, at $O((Y_\Sigma v)^2)$, the total amplitude is given by:

$$T (\mu \to e\gamma) = i \frac{G^S_M}{\sqrt{2}} \frac{e}{32 \pi^2} m_\mu \overline{\nu}_e (p-q) (1 + \gamma_5) i \sigma_{\lambda\nu} q^\lambda u_\mu (p) \times \left\{ \frac{13}{3} + C \right\} \epsilon_{\nu\mu} - \sum_i x_{\nu_i} (U_{PMNS})_{ei} (U^\dagger_{PMNS})_{i\mu},$$

where $x_{\nu_i}$ are the neutrino oscillation parameters.
where \( C = -6.56, \epsilon = \frac{v^2}{4}Y_\Sigma^\dagger M_\Sigma^2 Y_\Sigma \) and \( x_{\nu} = \frac{m_\nu^2}{M_{\nu q}} \). The first part of the amplitude correspond to the contribution of the fermionic triplet, while the second one is the usual contribution from neutrino mixing (suppressed by a GIM cancellation). The branching ratio then reads:

\[
Br(\mu \to e\gamma) = \frac{3}{32\pi} \frac{\alpha}{\pi} \left( \frac{13}{3} + C \right) \epsilon_{e\mu} - \sum_i x_{\nu_i} (U_{PMNS})_{ei} \left( U_{PMNS}^\dagger \right)_{\mu i} \right|^2 .
\]

(5)

\( \tau \to l\gamma \) decays can be obtained from Eq. (5) by replacing \( \mu \) by \( \tau \), \( e \) by \( l \) and by multiplying the obtained result by the product of the two ratios. The latter branching ratios have already been calculated in the type III seesaw model. It turns out that the bounds obtained from these decays exactly apply on the same parameters \( \epsilon \) as those obtained in Eqs. (6)-(8). This relation between the two types of decays implies that the flavour structure of the \( \mu \)-to-\( e \) fermionic line is the same in both processes. Regarding the couplings, there is only one way to combine the two Yukawa couplings and two inverse \( M_\Sigma \) mass matrices to induce a \( \mu \)-to-\( e \) transition along a fermionic line: \( \epsilon_{e\mu} \). This relation between the two types of decays implies that the ratios of these branching ratios are fixed:

\[
|\epsilon_{e\mu}| = \frac{\nu^2}{2} \left| Y_\Sigma^\dagger \frac{1}{M_\Sigma} Y_\Sigma |_{\mu e} \right| \leq 1.1 \cdot 10^{-4},
\]

(6)

\[
|\epsilon_{\mu\tau}| = \frac{\nu^2}{2} \left| Y_\Sigma^\dagger \frac{1}{M_\Sigma} Y_\Sigma |_{\tau\mu} \right| \leq 1.5 \cdot 10^{-2},
\]

(7)

\[
|\epsilon_{e\tau}| = \frac{\nu^2}{2} \left| Y_\Sigma^\dagger \frac{1}{M_\Sigma} Y_\Sigma |_{\tau e} \right| \leq 2.4 \cdot 10^{-2}.
\]

(8)

### 4 Comparison to \( \mu \to eee \) and \( \tau \to 3\ell' \) decays

The presence of heavy fermions not only allows for one-loop lepton flavour violating processes, but also for tree level decays such as \( \mu \to eee \) and \( \tau \to 3\ell' \). The latter branching ratios have already been calculated in the type III seesaw model. It turns out that the bounds obtained from these decays exactly apply on the same parameters \( \epsilon \) as those obtained in Eqs. (6)-(8). To understand this property let us study the example of \( \mu \to e\gamma \) and \( \mu \to eee \). In both cases one wants to link a muon to an electron with a same fermionic line. The only way to achieve this is to mix a muon and an electron with a fermionic triplet. This implies that the flavour structure of the \( \mu \)-to-\( e \) fermionic line is the same in both processes. Regarding the couplings, there is only one way to combine two Yukawa couplings and two inverse \( M_\Sigma \) mass matrices to induce a \( \mu \)-to-\( e \) transition along a fermionic line: \( \epsilon_{e\mu} \). This relation between the two types of decays implies that the ratios of these branching ratios are fixed:

\[
Br(\mu \to e\gamma) = 1.3 \cdot 10^{-3} \cdot Br(\mu \to eee),
\]

(9)

\[
Br(\tau \to \mu\gamma) = 1.3 \cdot 10^{-3} \cdot Br(\tau \to \mu\mu\mu) = 2.1 \cdot 10^{-3} \cdot Br(\tau^\pm \to e^- e^+ \mu^-),
\]

(10)

\[
Br(\tau \to e\gamma) = 1.3 \cdot 10^{-3} \cdot Br(\tau \to eee) = 2.1 \cdot 10^{-3} \cdot Br(\tau^- \to \mu^- \mu^+ e^-). \quad \text{(11)}
\]

Since the processes \( \ell \to 3\ell' \) occur at tree level in this model while the \( \ell \to \ell'\gamma \) ones are one-loop, small values of the ratios are expected. The \( \mu \)-to-\( eee \), \( \tau \)-to-\( eee \) and \( \tau \)-to-\( \mu\mu\mu \) decays lead to \( |\epsilon_{e\mu}| < 1.1 \cdot 10^{-6}, |\epsilon_{\mu\tau}| < 4.9 \cdot 10^{-4}, |\epsilon_{e\tau}| < 5.1 \cdot 10^{-4} \). Those bounds are better than the one obtained in Eqs. (6)-(8). This means that in this model the tree-level processes will provide the most competitive bounds on the \( \epsilon_{\alpha\beta} \) parameters, even if the the experimental limit on the branching ratios of the radiative decays improves by two order of magnitude. Using the experimental bounds \( Br(\mu \to eee) < 1 \cdot 10^{-12}, Br(\tau \to eee) < 3.6 \cdot 10^{-8} \) and \( Br(\tau \to \mu\mu\mu) < 3.2 \cdot 10^{-8} \), one derives predictions for the bounds on branching ratios of the radiative decays:

\[
Br(\mu \to e\gamma) < 10^{-15} \quad \text{(12)}
\]

\footnote{Note that these ratios are obtained in the limit where \( M_\Sigma \gg M_{W,Z,H} \). Not working in this limit, these ratios can vary up to one order of magnitude.}
\[ Br(\tau \rightarrow \mu \gamma) < 4 \cdot 10^{-11} \]
\[ Br(\tau \rightarrow e\gamma) < 5 \cdot 10^{-11} \]

(13)

(14)

to be compared with experimental bounds\(^6,7\): \( Br(\mu \rightarrow e\gamma) < 1.2 \cdot 10^{-11} \), \( Br(\tau \rightarrow \mu \gamma) < 4.5 \cdot 10^{-8} \), \( Br(\tau \rightarrow e\gamma) < 1.1 \cdot 10^{-7} \).

5 Conclusion

In our work we were lead to the conclusion that the observation of one leptonic radiative decay in the upcoming experiments will rule out the seesaw mechanism with only fermion triplets. Indeed this would contradict bounds arising from present experimental limits on the \( \mu \rightarrow eee \) and \( \tau \rightarrow 3l \) decay rates, and therefore imply that there exist other sources of lepton flavour violation than those associated to triplet of fermions.

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References


The single largest background to future $\bar{\nu}_e \rightarrow \bar{\nu}_e (\nu_e \rightarrow \nu_e)$ oscillation searches is neutral current (NC) $\pi^0$ production. MiniBooNE, which began taking antineutrino data in January 2006, has the world’s largest sample of reconstructed $\pi^0$'s produced by antineutrinos. These neutral pions are primarily produced through the $\Delta$ resonance but can also be created through “coherent production.” The latter process is the coherent sum of glancing scatters of antineutrinos off a neutron or proton, in which the nucleus is kept intact but a $\pi^0$ is created. A signature of this process is a $\pi^0$ which is highly forward-going. It is advantageous to study coherent production using antineutrinos rather than neutrinos because the ratio of coherent to resonant scattering is enhanced in antineutrino running. The first measurement of NC coherent $\pi^0$ production in the MiniBooNE antineutrino data is discussed here.

1 Neutral Current $\pi^0$ Production

At low energy, neutral current (NC) $\pi^0$'s are produced via two different mechanisms:

$$\bar{\nu}N \rightarrow \bar{\nu}\Delta \rightarrow \bar{\nu}\pi^0N \quad (\text{resonant})$$
(1)

$$\bar{\nu}A \rightarrow \bar{\nu}A\pi^0 \quad (\text{coherent})$$
(2)

In resonant $\pi^0$ production, a(n) (anti)neutrino interacts with the target, exciting the nucleon into a $\Delta^0$ or $\Delta^+$, which then decays to a nucleon plus $\pi^0$ final state. In coherent $\pi^0$ production, very little energy is exchanged between the (anti)neutrino and the target. The nucleus is left intact but a $\pi^0$ is created from the coherent sum of scattering from all the nucleons. A signature of this process is a $\pi^0$ which is highly forward-going.
1.1 Why Study NC $\pi^0$ Production?

NC $\pi^0$ events are the dominant background to $\bar{\nu}_\mu \rightarrow \bar{\nu}_e$ ($\nu_\mu \rightarrow \nu_e$) oscillation searches. A $\pi^0$ decays very promptly into two photons ($\tau_{\pi^0} \sim 8 \times 10^{-17}$s), and can mimic a $\bar{\nu}_e$ ($\nu_e$) interaction if only one photon track is resolvable in a detector.

In particular, coherent production is much more challenging to predict theoretically than resonant processes. Unfortunately, there are currently only two published measurements of the absolute rate of antineutrino NC $\pi^0$ production (Figure 1); the lowest energy measurement reported with 25% uncertainty at 2 GeV$^1$. There exist no experimental measurements below 2 GeV. Furthermore, current theoretical models on coherent $\pi^0$ production$^2,3,4$ can vary by up to an order of magnitude in their predictions at low energy, the region most relevant for (anti)neutrino oscillation experiments.

The analysis presented at the 43rd Rencontres de Moriond represents the first time we are experimentally probing this process in this energy region.

![Figure 1: World data on antineutrino NC coherent $\pi^0$ production at low energy.](image)

2 The MiniBooNE Experiment

The Mini Booster Neutrino Experiment (MiniBooNE)$^5$, an experiment at Fermilab designed to measure $\nu_\mu \rightarrow \nu_e$ oscillations, turns out to be very well-suited for $\pi^0$ physics. Its large, open-volume Cherenkov detector with full angular coverage provides excellent $\pi^0$ identification and containment. In fact, MiniBooNE has the world’s largest samples of NC $\pi^0$ events in interactions with $\sim$1 GeV neutrinos ($\sim$28k) and with $\sim$1 GeV antineutrinos ($\sim$1.7k). Additional POT has been collected in $\nu$ mode since the MiniBooNE oscillation results$^5$ with additional POT being collected in $\bar{\nu}$ mode currently.

3 Coherent $\pi^0$'s in Terms of $E_\pi(1 - \cos \theta_\pi)$ in $\bar{\nu}$ Mode

As mentioned before, coherent and resonant $\pi^0$ production are distinguishable by $\cos \theta_\pi$, which is the cosine of the lab angle of the outgoing $\pi^0$ with respect to the beam direction. It turns
out that it is even better to study coherent $\pi^0$’s in terms of the pion energy-weighted angular distribution since in coherent events, $E_\pi (1 - \cos \theta_\pi)$ has a more regular shape as a function of momentum, than $\cos \theta_\pi$ alone. Thus, we will fit for the coherent content as a function of the pion energy-weighted angular distribution. Furthermore, we need to fit this quantity simultaneously with the invariant mass. This is due to the fact that, in the energy-weighted angular distribution, the resonant and background spectra have similar shapes that are distinct from the forward-peaking coherent spectrum while the resonant and coherent spectra have similar shapes in the invariant mass that are distinct from the background spectrum.

This fit has in fact been done in neutrino mode\textsuperscript{6}. Preliminary fits to the antineutrino data were shown at this conference. This sample is important because in antineutrino scattering, there is a helicity suppression for most interactions, including resonant production of $\pi^0$’s, but not for coherent production. Thus, the ratio of coherent to resonant scattering, which is small, is expected to be enhanced in antineutrino running.

3.1 Preliminary Results

Preliminary statistics-only fits between MiniBooNE antineutrino data and MC \textsuperscript{a} are shown (Figures 2 and 3). The initial MC includes a rescaling of the Rein-Sehgal\textsuperscript{8} coherent cross section based on the measurement in Ref. 6. These studies clearly show evidence for NC coherent $\pi^0$ production, as was the case in neutrino mode. When coherent $\pi^0$ production is absent from the MC (Figure 4), one can see the poor agreement between data and MC in the lower end of $E_\pi (1 - \cos \theta_\pi)$, where most of the $\pi^0$’s are expected to be produced from coherent scattering. The agreement improves considerably when coherent $\pi^0$ production is included (Figure 3). Furthermore, the coherent fraction in both $\nu$ and $\bar{\nu}$ modes is roughly 1.5 times lower than the Rein-Sehgal\textsuperscript{8} prediction, which is the most widely used model in (anti)neutrino experiments.

![Figure 2: Preliminary statistics-only invariant mass fit.](image)

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\textsuperscript{a}RES = NUANCE channels 6,8,13, and 15. COH = NUANCE channel 96. BGD = NUANCE channels other than 6,8,13,15, and 96. MC TOT = RES + COH + BGD. See Ref. 7 for channel definitions.
Figure 3: Preliminary statistics-only pion energy-weighted angular distribution fit.

Figure 4: Preliminary statistics-only pion energy-weighted angular distribution fit with no coherent contribution in the total MC.

References

WORKSHOP SUMMARY - EXPERIMENTAL

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The experimental results presented at the workshop are summarised and reviewed.

1 Introduction

The 43rd Rencontres de Moriond Electroweak Sessions produced a cornucopia of new results from accelerator and non-accelerator experiments across an impressively wide range, from searches for the Higgs and exotic particles at the Tevatron and HERA to precision electroweak measurements at the flavour factories and fixed target kaon experiments, and new results on neutrino interactions and astroparticle observations. In addition, the prospects for physics at the LHC were discussed.

The talks were excellent, and should be consulted for the details of the results discussed below. References to talks given at the workshop will be indicated by the italicised author, for example Peach.

2 Higgs Searches

There were several talks (Haas, Ochando, Masubuchi, Yorita, Zivkovic) on the searches for the standard model and non-standard Model Higgs at the Tevatron. So far, between 1 and 2 fb$^{-1}$ of data have been analyzed by CDF and D0, with a total of 3 fb$^{-1}$ each on tape and ready to be analyzed. If the Tevatron continues to run well, there is every prospect that each experiment will have accumulated 6 - 8 fb$^{-1}$ by the end of 2009, and running in 2010 would extend this further.
Although there is not yet any real hint of a signal, there has been remarkable recent progress. As well as the excellent performance of the Tevatron, there have been two particular developments that have improved the prospects for making a discovery. Firstly, the analyses are becoming very sophisticated, increasing the sensitivity from a given data sample. Secondly, combined CDF and D0 analyses are emerging; while this does not improve the statistical significance when compared with combining the two results obtained separately, the consistent treatment of backgrounds, cuts and phenomenological uncertainties should reduce the systematic errors, again improving the significance of the result from a given dataset. The result of this combined analysis is shown in Figure 1 (Zivkovic), where the expected limits from the two experiments separately are also compared with the combined expected limit. The advantages of the combined approach can clearly be seen. There is an opportunity for the Tevatron to place limits on the existence of a Higgs at around 160 GeV/c², and perhaps even make a discovery in this region, although it is unlikely that the nature of the Higgs could be firmly established.

Searches for non-Standard Model Higgs continue, but in general the limits are more difficult to interpret because of the large range of parameters that also have to be specified within a given model, as shown in new results presented in Figure 2 (Haas).

Figure 1: The Standard Model Higgs limit from CDF and D0 combined.

Figure 2: The MSSM Higgs limits in the $H \rightarrow \tau \tau$ and $bH \rightarrow b\tau\tau$ channels.
3 QCD with Electroweak probes

An important development is the use of electroweak probes (W, Z...) to test QCD. These were reported on by J Han and L Han for the Tevatron, and by de Boer for the HERA experiments. The Tevatron analyses are particularly important, since they indicate new ways in which the parton distribution functions can be constrained at the LHC. As an example, Figure 3 from J Han shows the measured differential cross-sections for $Z/\gamma^*$ productions as a function of the boson rapidity for D0 and CDF, compared with predictions from MRST NNLO (D0) and CTEQ NLO (CDF) expectations. Likewise, the di-boson production rates can be compared with the theoretical expectations to give a sensitive test of the Standard Model calculus (see Figure 4 from L Han).

![Figure 3: Differential cross-sections for $Z/\gamma^*$ production as a function of boson rapidity. (left) D0, compared with the MRST-04 NNLO expectation. (right) CDF, compared with CTEQ6.1M NLO expectation.](image)

![Figure 4: Di-boson (WW,WZ, ZZ, W$\gamma$ and $-\gamma$) production rates at the Tevatron for CDF and D0, compared with theory. (1) First evidence of W$\gamma$ RAZ. (2) Z$\gamma$ cross-section measurement. (3) WW/WZ $\rightarrow l\nu jj$. (4) WWZ Triple Gauge Coupling. (5) First evidence of ZZ.](image)

HERA continues to produce precision tests of our understanding of the proton structure, and the particular features of the $ep$ collisions allow tests that cannot be obtained easily from either electron-positron or proton-(anti)proton colliders. As well as checking the chiral structure of the Standard Model, de Boer reported on the extraction of limits on the quark radius from
high $Q^2$ neutral current events, placing limits of less than 0.001 times the proton radius (see Figure 5).

Figure 5: Limits on the quark radius from H1 (left) and ZEUS (right).

4 Searches for SUSY and other new physics

There were excellent reviews of the state of the searches for Supersymmetry and other more exotic extensions to the Standard Model by Jaffre, Dube, Barbagli and Sauvan, both at the Tevatron and at HERA. So far, despite significant innovation in search strategies and sensitivity to new physics over a wide region of parameter space, no convincing signals have yet been found. While this is, at one level, disappointing it is not particularly surprising. The electroweak fits show that, at the current energy scale, the Standard Model gives an excellent description of the data. Only one of the 18 key measurable quantities ($A_{f_0}^{0,b}$) is between two and three standard deviations away from the Standard Model prediction; indeed, without this point, the electroweak fit is if anything too good. There are hints of new physics in, for example, the difference between the measured and expected values of the muon $g - 2$, but both the experiment and the theory are triumphs of human ingenuity, which is another way of saying that they are very difficult. Finally, it might also be argued that the fact that the best fit value for the Higgs mass is well below the LEP exclusion limit could be a hint that there is new physics somewhere in the domain. However, the main conclusion is that is either no new physics at the electroweak symmetry-breaking scale, or that there is a fortuitous conspiracy of the parameters to keep the new physics hidden at this scale. Whichever of these is the case, it is clear that the energy scale at which the new physics, if it exists, becomes manifest is much higher.

There are two further remarks that I would like to make.

1. It is likely that when (if?) the new physics is found, for example at the LHC, we will discover that there were hints of the new physics in several places, but that lack of understanding prevented it being recognised as such.

2. New model-independent search strategies, using a topological classification of the data and comparing the distribution of many quantities with the Standard Model expectations, should be further developed. Although these methods are not yet as sensitive as the cuts-based search strategies, they have the ability to identify anomalies in the data very quickly. Often, these anomalies will turn out to be due to errors in the Monte Carlo, and indeed this is a very good way to look for such errors. But eventually, it should be possible, once the
Monte Carlo model is sufficiently robust, to use these techniques to look for new physics. A promising example was shown by Sauvan (see Figure 6), which compares the H1 data with Standard Model expectations for 23 classes of events with isolated high-\(P_T\) leptons, photons, jets or missing energy and up to four other leptons, photons, jets or missing energy for both \(e^+p\) and \(e^-p\) data, revealing no serious discrepancies, and providing yet another test that could have failed but did not.

Figure 6: Comparison between H1 \(e^+p\) and \(e^-p\) data and the Standard Model Monte Carlo for 23 different topologies of leptons, photons, jets and missing energy.

5 Heavy Flavour

5.1 Top Physics

Until the LHC begins operating at reasonable luminosity, the Tevatron is the only place where the top can be studied directly.

Besançon discussed the measurement of the top production cross-sections, which can be compared with the expectations from the Standard Model using the measured value of the top mass. Alternatively, assuming that the Standard Model describes the processes, the measured cross-sections can be used to derive a value for the top mass, which can then be compared with the direct measurements. This has been done by both CDF and D0 in the process \(t\bar{t} \rightarrow b\bar{b}W^+W^-\) where the pair of \(W\)'s decay into various combinations of leptons and jets. A new cross-section measurement of \((7.42\pm0.53(stat)\pm0.46(syst)\pm0.45(theory, m_{top} = 175 GeV/c^2))pb\) was reported by D0 on 910pb\(^{-1}\) of data in the “lepton+jets” channel, where the b-jets were identified either by tagging or topology, in excellent agreement with expectations.

Schwienhorst reported on the single top production, where both CDF and D0 are making measurements in several channels, albeit with large (\(\approx 20\%)\) errors, and D0 reported a promising new approach comparing the \(s^-\) and \(t^-\)-channel production (see Figure 7). Chen reported on the direct measurements of the top mass, which is now measured to 1.1% by CDF and D0. An interesting new method was reported (see Figure 8) by Fedorko from CDF in the Young Scientists Forum, which uses a coupled channel analysis on the di-lepton and lepton+jets channels, and obtains the single most precise measurement of the top mass, based on 2fb\(^{-1}\) of data. A measure of the increasing sophistication of the analyses is shown in Figure 9 from Chen, which compares the evolution of the uncertainty in the top mass from CDF as a function of the integrated luminosity to that expected on the basis of the first 680pb\(^{-1}\); the uncertainty on the top mass is reducing faster than would be expected from the increase in statistics, and a precision of 1%
on the top mass should be achievable with the presently obtained integrated luminosity. It is equally remarkable that this precision is about a factor two better than the Run-II goal set out in the 1996 TDR. All of these developments are very encouraging for the LHC, which will be,

amongst other things, a top-factory.

5.2 $b$, charm and $\tau$ Physics

BELLE and BaBar continue to produce a wealth of new results on B-decays. BELLE has accumulated 770fb$^{-1}$ and BaBar has accumulated 510fb$^{-1}$ of data. Aushev reported on the time dependent CP violating analyses, where there is mostly good agreement between the various channels which can be used to measure $\beta/\phi_1$. However, a preliminary result from BELLE (see Figure 10) in the channel $B^- \rightarrow K_s\pi^0\pi^0$ shows a deviation from the standard model expectations of more than 2 $\sigma$. Liventsev reported on analyses of $B \rightarrow D^{*}\ell\nu$ decays using fully reconstructed B-tags, with the first measurement of the $B \rightarrow D_s^\pm \ell\nu; D_s^\pm \rightarrow D\pi$ decay chain. Mazur reviewed the semi-leptonic B and D decays, which are reaching remarkable precision, leading to an impressive precision on $V_{ub}$ (less than 10% error) and $V_{cb}$ (less than 2% error), as well as measuring the b-quark mass to better than 0.7%. H Kim reported on the status of leptonic B-decays, including lepton-flavour-violating channels, where there has been recent progress from BELLE, BaBar,
CDF and CLEO, with improvements in the upper limits for \( B^0 \rightarrow e\mu, e\tau \) and \( \mu\tau \) decays (see Figure 11). Simi reviewed the status of B-decays via Penguin loops, including the observation of the decay \( B \rightarrow \pi^+ K \). Limosani reported new results from BELLE on \( B \rightarrow X_s\gamma \), where heroic efforts have been made to measure the whole of the \( \gamma \) spectrum above about 1.7 GeV; these new measurements are consistent with earlier measurements, but significantly more precise. Sordini reported many new results on the angle \( \gamma/\phi_3 \) of the standard unitarity triangle, demonstrating the remarkable progress that has been made recently (see Figure 12). This, and other recent results, have been used by the CKMFitter group to update their analysis, and this was presented by Descotes-Genon (see Figure 13).

Meanwhile, B-physics at the Tevatron continues to make progress, whetting the appetite for the data from LHCb. Parua reviewed the status of the lifetime measurements of b-hadrons, including several new results. Di Giovanni discussed \( B_s \) decays and CP-violation, including the first measurement of \( 2\beta_s \). In the Standard Model, CP-violation is expected to be small in the \( B_s \) system \( (2\beta_s = (0.04 \pm 0.01)\text{ rad}) \), but the measurement by D0, consistent with the constraint from CDF, is more than 2 \( \sigma \) away from the Standard Model expectations.

CLEO-c (Park) and BES (Ho) are producing a plethora of results on charm and \( \tau \) physics, with results on branching fractions and new charm states being reviewed, and on precision measurements of the hadronic cross-sections in the charm region. Xu reported on several new
charm states seen by BES, as well as updates on many previously discovered states. BELLE and BaBar also produce copious numbers of charmed particles and $\tau$ laptons, as reviewed by Zupanc.

During the workshop, the UT-fit collaboration produced an analysis of all the data on the $b \rightarrow s$ transition, claiming evidence for new physics at greater than 3$\sigma$. They parametrise the new physics ($A_s^{NP}$) relative to the standard model ($A_s^{SM}$) as

$$C_{Bs} e^{2i\phi_{Bs}} = \left( A_s^{SM} e^{-2i\beta_s} + A_s^{NP} e^{2i(\phi_s^{NP} - \beta_s)} \right) / A_s^{SM} e^{-2i\beta_s}$$

so that the “no new physics” scenario has $C_{Bs} = 1$ and $\phi_{Bs} = 0$. The results of the analysis are shown in in Figure 15. While $C_{Bs}$ is consistent with the Standard Model, the phase $\phi_{Bs}$ seems to be inconsistent with the Standard Model expectations, at a level with greater than 3$\sigma$ significance. While this is intriguing, more evidence is probably required before a strong claim for new physics can be convincingly made. In particular, it implies that the amplitude for the new physics is comparable to the Standard Model amplitude, which (if true) makes it perhaps surprising that this has not already revealed itself elsewhere.
Figure 13: The updated fit to the available experimental data for the CKM parameters from the CKMFit collaboration.

Figure 14: Plot of $\phi_s = -2\beta_s$ versus $\Delta \Gamma$ from D0, together with the limits on $\phi_s$ from CDF, showing a clear discrepancy with the Standard Model expectations.

6 Kaon Physics

Although it is more than 60 years since the kaon was first observed in cosmic rays by Rochester and Butler, there is still much to be understood, and thanks to the heroic efforts of NA48, KTeV and KLOE, there has been a significant improvement in the measurement of some rather basic parameters. Recent measurements reported by Ruggiero (NA48) and Testa (KLOE) on the $K_{\ell 2}$ and $K_{\ell 3}$ decay modes improve significantly on the precision of earlier measurements, and lead to tighter constraints on the unitarity of the CKM matrix, particularly through better determination of $V_{us}$, and place new limits on lepton flavour violation. Glazov (KTeV) reported an update on the value of $\epsilon'/\epsilon$, with a new world average of $(16.4 \pm 1.4) \times 10^{-4}$. Perhaps more importantly, there is now good agreement (see Figure 16) between NA48 and KTeV. In addition, Archilli (KTeV) reported a new upper limit on the important branching ratio $K_S \rightarrow e^+e^-$ of $9.3 \times 10^{-9}$ at 90% confidence level, an order of magnitude smaller than previous measurements. Finally, Zimmerman reported on a new measurement by KTeV of the $\pi^0$ double Dalitz decay, the first such measurement in more than 47 years. The result of $(3.26 \pm 0.30) \times 10^{-5}$, based on 30.5k events, is fully consistent with the 1960 measurement based on only 112 events.
7 Neutrino Physics

Although not neutrino physics, it is very important for the oscillation experiments to have a good understanding of the hadron production differential cross-sections for the particular incident proton energy and target configuration. This is important for both the present generation of experiments (K2K, MINOS, MiniBooNE) and future experiments such as T2K and a possible neutrino factory or superbeam facility. The HARP experiment has taken data at many proton energies and with a variety of targets, and new results were reported by Tsenov. These data are also important for calculations of the atmospheric neutrino fluxes.

The low energy neutrino and antineutrino cross-sections, which are still relatively poorly known, are a significant source of systematic error. New results on the low-energy quasi-elastic and NC $\pi^0$ production were reported by Katori, and there are new experiments (SciBooNE and Minerva) under way to measure these low energy cross-sections.

7.1 Neutrino Oscillations

Pistillo reported on the status of the OPERA experiment at Gran Sasso, in the CNGS neutrino beam from CERN, where 38 events have already been observed. Habig gave an update to the MINOS results with $|\Delta m^2_{32}| = (2.38^{+0.20}_{-0.16}) \times 10^{-3}eV^2$ and $\sin^2 2\theta_{23} = 1.00_{-0.08}$, with a fit $\chi^2$
of 41.2 for 34 degrees of freedom (see Figure 17). Polly reported on the latest status of the MiniBooNE analysis, which now includes data from 250 MeV to 300 MeV. The excess of events below 500 MeV is still there, and at the lower energy, but is not compatible with neutrino oscillations. In an attempt to understand even better the detector and the data, the electron- and muon-like events seen in MiniBooNE from the NUMI beam line have been analysed, and preliminary results are in good agreement with the expectations.

Figure 17: The latest results from MINOS on $|\Delta m^2_{32}|$ and $\sin^2 2\theta_{23}$.

Finally, Dalnoki-Veress reported on the first published result from Borexino, which has detected the $^7$Be solar neutrinos at a rate consistent with the LMA MSW solution for neutrino oscillations, albeit with a relatively large systematic error. This is, nevertheless, an important measurement since it studies these low energy solar neutrinos in real time.

7.2 Neutrinoless double $\beta$ decay

Although the observation of neutrino oscillations establishes beyond doubt that at least two of the neutrino have a finite mass, it does not set the absolute mass-scale. Direct measurement of the electron neutrino mass is challenging, and the KATRIN experiment is rising to that challenge, but will be limited to an absolute neutrino mass scale of a few hundred meV. There are constraints from cosmology, since the neutrino mass has an influence on the way the early universe unfolds. However, in my view, the neutrino mass-scale should be an input to cosmology. Provided that neutrinos have a Majorana component, the rate of neutrinoless double $\beta$ decay is proportional to a weighted average neutrino mass, and thus, when combined with the two independent mass-differences squared, sets the absolute neutrino mass scale. Vignati (CUORICINO) and Broudin-Bay (NEMO-3) reported new results on the search for neutrinoless double $\beta$ decay. CUORICINO reported a new preliminary result based on 15.53 kg-years of data on $^{130}$Te yields a limit (90% confidence level) of $3.1 \times 10^{21}$ y, which translates to a limit on the the weighted average neutrino mass $m_{\beta\beta}$ of between 0.20 eV and 0.68 eV, where the uncertainty depends upon uncertainties in the nuclear matrix elements. NEMO-3 reported a new preliminary limit of $1.46 \times 10^{22}$ for $^{150}$Nd, almost an order of magnitude higher than the best previous limit, corresponding to a limit on $m_{\beta\beta}$ of between 3.7 eV and 5.1 eV.

8 Dark Matter

There were several talks on Dark Matter experiments. Perhaps the most imaginative development, in my opinion, is the COUPP (Chicagoland Observatory for Underground Particle
Physics) experiment in the MINOS near detector hall at Fermilab, reported by Szydagis. This uses the (very old) Heavy Liquid Bubble Chamber technique in a novel way, by deliberately desensitizing the superheated liquid so that minimum ionising particle leave no trace, and holding the liquid in a metastable phase for long enough to act as a WIMP detector. Its first result is an impressive limit compared with previous measurements (see Figure 18), incompatible with the DAMA-allowed region. Lang reported on a new result from the CRESST experiment, which is not yet competitive with CDMS, but is very promising since the analysis is based on only 67 kg-d of data taking. S-K Kim reported on the KIMS Dark Matter experiment, which uses CsI as the target and detector and so is more directly comparable with the DAMA NaI data; their results are incompatible with the DAMA claims of a signal for spin-independent interactions, and are incompatible with the DAMA limits for spin-dependent interactions for WIMP masses above 20 GeV/c². Santorelli reported on the status of the XENON experiment, with a recent spin-independent limit comparable to the latest CDMS limit (see Figure 19) Finally, Ahlen discussed the options for a large, affordable gaseous TPC which could be used as a directional Dark Matter detector.

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Figure 18: The first results from COUPP on the search for Dark Matter.

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Figure 19: Recent on spin-independent limits from the XENON Dark Matter experiment.

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*This workshop took place before the most recent result from DAMA, making stronger claims for a Dark Matter signal based on significantly more data; the experiments reported here are also inconsistent with this recent DAMA result (see arXiv:0804.2741v1 (astro-ph)).*
9 Astroparticle Physics

The past decade has seen a dramatic growth in the field of Astroparticle Physics (or Particle Astrophysics). Many of these experiments use particle physics techniques to explore astrophysical objects, using particles (other than light) as messengers. These experiments explore the Universe in a new way, and could reveal many new and unexpected phenomena. However, these experiments also have the potential to reveal new features in particle physics.

Glicenstein reported on the high energy gamma ray experiments (HESS, MAGIC, Whipple) and the prospects for GLAST (2008) and upgrades to MAGIC (2008) and HESS (2009). These experiments are sensitive in some models to Dark Matter through the annihilation processes in the galactic centre, and can yield interesting bounds on the scale of quantum gravity. Roucelle reported on the status of the AMANDA and ICE-CUBE experiments, looking for point sources of neutrinos, so far without success, and Pradier reported on the status of ANTARES, which is beginning to see candidate neutrino interactions.

The CAST experiment at CERN, recycling a prototype LHC dipole magnet, is looking for evidence for axions from the sun via the inverse Primakoff effect, and is beginning to achieve sensitivities that will constrain the models, as reported by Borghi.

There were two reports from the AUGER experiment, where Maris reported on the spectrum of ultra-high-energy cosmic rays (UHECR) and the evidence for the GZK cutoff, and on the evidence that their composition has several components. Bonino reported on the evidence for anisotropies and correlations with known astrophysical objects for these ultra-high-energy cosmic rays, although there was a comment by Tinyakov that the statistical significance may be less than claimed. Arqueros presented a review of methods of calculating the air fluorescence yield from UHECR, emphasizing the progress that had been made but also the fact that there were still some discrepancies and model dependencies despite general agreement, so that the absolute values have an uncertainty of about 10%.

10 The LHC machine and experiments

Of course, many of the unresolved issues discussed above - the existence and mass of the Higgs and any low-lying Supersymmetry, Dark Matter, flavour physics etc - should be greatly illuminated by the data from the LHC. Evans gave a report on the status of the LHC machine, where the final cool-down was well under way. At the time of the meeting, about half of the machine was cold or in the process of cooling down, although one of the sectors would be warmed up to complete the repair of the inner triplets. The schedule (see Figure 20) should see the machine cold and ready for beam in July, provided that there is no unforeseen reason to warm up one of the other sectors. Plamondon and Christiansen reported on the state of readiness of the ATLAS and CMS detectors, and demonstrated that both experiments are ready to extract the first physics once collisions are achieved, and Tsuno and Bellan showed that the Higgs and any exotic physics would be clearly seen.

The next two or three years are going to be very exciting. The Higgs (if it exists), or the mechanism that so closely mimics (it if it does not), will be revealed at last and its nature explored. We can also hope that some of the many new phenomena (some more exotic than others) that could begin to address some of the unexplained features of the Standard Model may be discovered, probably sooner rather than later. Of course, we may find little new, but I would be both surprised and disappointed if that were the case - the Standard Model, while an excellent description of physics at the electroweak scale, is clearly only an effective theory. We need the LHC to illuminate the road ahead, towards a more fundamental and more satisfying description of the universe around us.
11 Summary and Conclusion

The Rencontres de Moriond Electroweak Sessions provide an opportunity to review and discuss the state of the Standard Model description of the atto-world, and this year has seen the presentation of many new results. Most confirm that the Standard Model is an excellent description of the interactions up to the electroweak unification scale. There are a few results which challenge the Standard Model, but these are still at the level of “hints” of, rather than solid evidence for, new physics. There is, however, a feeling “in the air” that new physics is just around the corner - at the Tevatron, the flavour factories, fixed-target or non-accelerator experiments and, if not found there, at the LHC. The next few years should be very exciting, and should reveal the physics that underpins the Standard Model.

Acknowledgments

I would like to thank all of the speakers for their contributions, and for letting me use their results. I would also like to thank the organisers for their excellent organisation, and particularly Jean Tranh Tan Van, Lydia Iconomidou-Fayard and Jean-Marie Frère for making Moriond possible, and for preserving its essential spirit. The tranquility and staggeringly beautiful scenery of La Thuile and the fresh mountain air provide a stimulating environment in which to be immersed in physics. I enjoyed myself enormously.

References

### List of Participants

**RENCONTRES DE MORIOND 2008**  
**Electroweak Interactions and Unified Theories**

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