STRONG INTERACTIONS
AND GAUGE THEORIES
The Hadronic Session of the Twenty-first Rencontre de Moriond on
STRONG INTERACTIONS AND GAUGE THEORIES

was organized by

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FOREWORD

The XXIst Rencontre de Moriond was held in 1986 in Les Arcs, Savoie (France).

The first such meeting was 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 at the first meeting were mainly French physicists interested in electromagnetic interactions. In subsequent years, a session on high energy strong interactions was also added.

The main purpose of these meetings 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 meeting helps to develop better human relations as well as a more thorough and detailed discussion of the contribution.

This concern for research and experimentation of new channels of communication and dialogue which from the start animated the Moriond meetings, inspired us to organize a simultaneous meeting of biologists on Cell Differentiation and to create the Moriond Astrophysics meeting. Common meetings between biologists, astrophysicists and high energy physicists are organized to study the implications of the advances in one field into the others. I hope that these conferences and lively discussions may give birth in the future to new analytical methods or new mathematical languages.

At the XXIst Rencontre de Moriond, in 1986, four physics session and one biology session were organized:

* January 25 - February 1   "Massive Neutrinos in Particle Physics and Astrophysics"

* March 9 - 16   "Perspectives in Electroweak Interactions and Unified Theories"

"Accretion Processes in Astrophysics"

* March 15 - 16   "Future Colliders"

* March 16 - 22   "Strong Interactions and Gauge Theories"

"Cell Differentiation"
I thank the organizers of the XXIst Rencontre de Moriond:

- O. FACKLER, M. MUGGE, and F. VANNUCCI for the "Massive Neutrinos in Particles Physics and in Astrophysics" session,

- A. BOUQUET, J. ERNWEIN, P. FAYET, G. GOLDHABER, J. F. GRIVAZ, G. KANE, A. MOREL, L. OLIVER, and I. VIDEAU for the "Progress in Electroweak Interactions" session,

- J. AUDOUZE, C. CESARSKY, Ph. CRANE, Th. GAISSER, D. HEGYI, and J. W. TRURAN, for the Astrophysics meeting,

- L. LEDERMAN and H. SOPER for the "Future Colliders" session,

- A. CAPELLA, D. DENEGRI, P. FRANZINI, H. FRISCH, L. MONTANET, R. PESCHANSKI, B. PIETRZYK, and F. M. RENARD for the "Strong Interactions and Gauge Theories" session,

- G. BELLIARD, D. DECARIES, M. FELLOUS, D. RIQUIER and K. TRAN THANH VAN for the biology meeting,

and the conference secretaries J. BORATAV, S. BOUHET, K. CLERJEAUD, K. DOUPLITZKY, M. FURGOLLE, C. JOUANEN, O. LEBEY, F. LEFEVRE, N. LESBRE, C. PIETRZYK, LE VAN SUU, L. NORRY, A. RAMOS, E. TRAN THANH TAM and V. URBAN who have devoted much their time and energy to the success of this Rencontre.

I am also grateful to M. J. DUPUY, E. ROCCA-SERRA, D. TOURAILLE, C. IRIGOYEN who contributed through their hospitality and cooperation to the well-being of the participants enabling them to work in a relaxed atmosphere.

This Rencontre is sponsored by the "Centre National de la Recherche Scientifique" and by the Commissariat à l'Energie Atomique". The "Massive Neutrinos" workshop is also sponsored by "Lawrence Livermore Laboratory". I would like to express my thanks to their encouraging support.

I wish sincerely that a fruitful exchange and an efficient collaboration between the physicists, the astrophysicists and the biologists will arise from this Rencontre as in the previous one.

J. TRAN THANH VAN
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JET PHYSICS IN THE UA1 EXPERIMENT

UA1 Collaboration

Presented by R. Sobie
University of Victoria, Victoria, B.C., Canada

ABSTRACT

This paper reviews the results of the analysis of jets in the UA1 experiment. The inclusive jet cross-section, two-jet angular distribution, and three-jet cross-section are presented and compared with current QCD predictions.
1. INTRODUCTION

The analysis of jets in the UA1 experiment is based primarily on data taken during the 1983 and 1984 runs. The results come from an integrated luminosity of approximately 100 nb$^{-1}$ for the 1983 run at $\sqrt{s} = 546$ GeV and 300 nb$^{-1}$ for the 1984 run at $\sqrt{s} = 630$ GeV. In this paper we present the inclusive jet cross-section, the two-jet angular distribution, and the three-jet cross-section. All results are compared with QCD predictions.

In UA1, a jet is defined by a jet algorithm\textsuperscript{1}, which combines calorimeter hits within a cone of radius 1 unit in $\eta$ and $\phi$ space around the highest transverse energy hits ($\eta$ is the pseudorapidity and $\phi$ is the azimuthal angle in radians). The energy and momentum of each jet are computed by taking the respective scalar and vector sum over the associated calorimeter cells. A correction was applied to the measured energy and momentum of each jet, as a function of the pseudorapidity and azimuth of the jet, on the basis of Monte Carlo analysis, to account for the effect of uninstrumented material and containment losses.

2. INCLUSIVE JET CROSS-SECTION

The inclusive jet cross-section for the reaction

$$p\bar{p} \rightarrow \text{jet} + X$$

is shown in Fig. 1a for the two centre-of-mass energies $\sqrt{s} = 546$ and 630 GeV\textsuperscript{2}. The errors shown include statistical errors and an energy-dependent uncertainty. The overall systematic uncertainty in the cross-section is estimated to be $\pm 70\%$. The present data are in good agreement with similar results from the UA2 Collaboration\textsuperscript{3}.

![Graph](image-url)

Fig. 1 (a) The inclusive jet cross-section; and (b) the scaled jet cross-section
We can compare the inclusive jet cross-sections with QCD calculations evaluated at $\eta = 0$. The curves shown in Fig. 1a were calculated from a leading order in $\alpha_s$, two-parton scattering QCD calculation using the structure functions of Eichten et al.\(^5\). The scale is defined to be $Q^2 = p_T^2$ and the $\Lambda$ parameter in the strong coupling constant $\alpha_s$ is equal to 0.2 GeV. The QCD calculation gives a good description of the $p_T$ dependence of the data. The best fit to the data requires that the theoretical curves be multiplied by a factor of 1.5, which is within the systematic errors and accuracy of the calculation. In Fig. 1, all curves have been scaled upward by a factor of 1.5. Uncertainties in the $Q^2$ scale, the QCD scale parameter, the choice of structure function, and higher-order corrections do not allow a greater accuracy than a factor of 2 in the QCD calculations.

The increase in the cross-section between the two centre-of-mass energies can be accounted for in terms of $x_T$ scaling. Figure 1b shows the dimensionless quantity $p_T^2 \, E(d^3s/dp^3)$ plotted against $x_T = 2p_T/\sqrt{s}$ for the two different beam energies. On this plot the two sets of data overlap, demonstrating that the observed increase in the cross-section with $\sqrt{s}$ is entirely consistent with perfect scaling. A much larger lever arm in energy, for example, $\sqrt{s} = 2000$ GeV (broken curve), would be needed to be sensitive to non-scaling QCD effects, although such effects have been observed previously in ISR experiments\(^6\).

3. TWO-JET CROSS-SECTIONS

The two-jet sample is selected from all events with $\geq 2$ jets defined by the jet algorithm. For each event the corrected four-vectors are transformed into the rest frame of the two highest-$p_T$ jets.

The two-jet angular distributions as a function of $\cos \theta$ are shown in Fig. 2a. The data are plotted for two-jet masses $M_{2J} = 150-250$ GeV and are based on data from the 1983 run at $\sqrt{s} = 546$ GeV\(^6\). The data can be seen to exhibit the characteristic Rutherford angular dependence $(1 - \cos \theta)^{-2}$. The dashed curve is a first-order QCD scaling prediction for the angular distribution, while the solid curve incorporates various non-scaling effects (i.e. variation of the structure functions and $\alpha_s$ with $Q^2$).

In Fig. 2b, we replot the angular distribution as a function of the variable $x = (1 + \cos \theta)/(1 - \cos \theta)$. In the limit as $x$ becomes large, the scaling or Rutherford-like angular distribution becomes flat. As in Fig. 2a, the dashed curve is the scaling QCD prediction, while the solid curve is the non-scaling prediction. It is clear from both figures that the scaling prediction does not give good agreement, even when all the various subprocesses are included. Only when non-scaling effects are included is there good agreement with the data.

Fig. 2 The two-jet angular distribution plotted against $\cos \theta$ (a) and $x$ (b) from the data at $\sqrt{s} = 546$ GeV
The two-jet angular distributions from the 1984 run at $\sqrt{s} = 630$ GeV has also been measured. In Fig. 3 we plot the angular distribution for the two-jet masses $M_{21} = 200-240$ GeV and $M_{21} = 240-300$ GeV. The higher trigger threshold used during the 1984 run meant that only two-jet events above $M_{21} > 200$ GeV cover the full angular acceptance as compared to $M_{21} > 150$ GeV for the 1983 run. The solid curves shown in Fig. 3 are the (non-scaling) QCD predictions for the given mass bins. In both cases, the data are consistent with the QCD prediction.

4. THREE-JET CROSS-SECTION

The three-jet sample is selected from all events with $\geq 3$ jets defined by the jet algorithm. For each event the corrected four-vectors are transformed into the rest frame of the three highest-$p_T$ jets.

The three-jet cross-section can be expressed in terms of four independent dimensionless variables in the centre-of-mass system: $x_3, x_4, \theta_3,$ and $\phi$ (see Fig. 4). The variables $x_3$ and $x_4$ are the Dalitz-plot variables defined as $x_i = 2E_i/\Sigma E_i$, where $E_i$ is the energy of jet-$i$ ($i = 3, 4, 5$) and ordered such that $x_3 > x_4 > x_5$ ($x_3 + x_4 + x_5 = 2$). The variable $\theta_3$ is the angle of jet-3 with respect

---

**Fig. 3** The two-jet angular distributions for $M_{21} = 200-240$ GeV (a) and $M_{21} = 240-300$ GeV (b) from the data at $\sqrt{s} = 630$ GeV

**Fig. 4** The three-jet variables in the subprocess centre-of-mass system
to the beam, and $\psi$ is the angle between the plane made by jet-4 and jet-5 and the plane made by jet-3 and the incoming partons.

The dominant subprocesses in the three-jet cross-section are predicted to be $gg \rightarrow ggg$, $qg \rightarrow qgg$ and $q\bar{q} \rightarrow q\bar{g}$. The leading-order QCD prediction for these subprocess cross-sections is given by

$$d^4\sigma/dx_3 dx_4 d\cos \theta_3 d\psi = \left(\frac{\alpha_s}{\pi}\right) \left| x_3 x_4 x_5 (1-x_3)(1-x_4)(1-x_5)\right|^{-1},$$

(2)

where $x_{T1} = x_i \sin \theta$.

The three-jet cross-section can become extremely large in certain cases. The first case is when $x_{T1} \rightarrow 0$ or one of the jets becomes close to the beam, which is the preferred configuration for initial-state bremsstrahlung. The second case is when $x_3 \rightarrow 1$ and jet-4 and jet-5 become collinear, which is the preferred configuration for final-state bremsstrahlung. The cuts applied to the three-jet sample attempt to confine the region of phase space where all the jets are well resolved from the beam and from each other, and where the theoretical cross-sections are well-behaved and relatively slowly varying.

The Dalitz plot for $x_3$ versus $x_4$ is shown in Fig. 5. The solid points are the data from the 1983 run at $\sqrt{s} = 546$ GeV and the crosses are the preliminary results from the 1984 run at $\sqrt{s} = 630$ GeV. Both data sets are found to be consistent with each other. From the plot we observe that there is a deviation from phase space as $x_3 \rightarrow 1$, indicating that $x_4$ and $x_3$ are collinear, which is suggestive of final-state bremsstrahlung. The solid curves shown in Fig. 5 are the single-bremsstrahlung QCD prediction based on Eq. (2). The dashed curves are the phase-space predictions. Clearly the data are consistent with the QCD prediction and not the phase-space prediction.
In Fig. 6, the three-jet angular distribution is plotted for the variables $\cos \theta_3$ and $\psi$. As with the Dalitz plot, the 1983 data are the solid points and the 1984 data are the crosses. Again, both data sets are entirely consistent with each other. The $\cos \theta_3$ distribution has the familiar Rutherford dependence seen in the two-jet events. The $\psi$ distribution is peaked at $0^\circ$ or $180^\circ$, indicating that jet-4 and jet-5 prefer to lie in the same plane as jet-3 and the incoming partons. This can be interpreted as evidence for final-state bremsstrahlung. In addition, no asymmetry is observed in either the $\cos \theta_3$ or $\psi$ distributions. Both $\cos \theta_3$ and $\psi$ are defined so that $\cos \theta_3$ is positive when jet-3 points along the fast parton and $\psi$ is positive when jet-4 points along the fast parton. QCD predicts a slight asymmetry in the $\psi$ distribution due to $qg$ initial states\(^\text{4}\), that is unobservable within the errors of the data.

5. **COMPARISON OF TWO-JET WITH THREE-JET EVENTS**

In QCD the relative yield of three-jet and two-jet events is directly related to the value of $\alpha_s$. Integrating the differential cross-sections over the dimensionless variables gives

\[
\sigma_{2j} = \frac{C_{2j} \alpha_s^2}{8}
\]

\[
\sigma_{3j} = \frac{C_{3j} \alpha_s^3}{8}
\]

where $C_{2j}$ and $C_{3j}$ are constants. Although $C_{2j}$ and $C_{3j}$ depend strongly on the subprocess, the ratio $C_{3j}/C_{2j}$ is very similar for the various incoming parton combinations. Thus the three-jet to two-jet ratio can be simply written as

\[
\frac{\sigma_{3j}}{\sigma_{2j}} = \frac{C_{3j}}{C_{2j}} \alpha_s
\]

and is independent of the parton densities, which largely cancel in the ratio.
Figure 7 shows the three-jet to two-jet ratio plotted as a function of the subprocess centre-of-mass energy. The dots correspond to the published 1983 data at $\sqrt{s} = 546$ GeV and the crosses to preliminary 1984 data at $\sqrt{s} = 630$ GeV. The solid curve is a leading-order QCD prediction calculated on the assumption that the $Q^2$ scales for the three-jet and two-jet events are identical. The broken curve, which fits the data better, has a smaller $Q^2$ scale for the three-jet events than the two-jet events [$Q_{3I} = (2/3)Q_{2I}$].

The determination of $\alpha_s$ from the data is limited by the lack of higher-order corrections to the QCD predictions for the two-jet and three-jet cross-sections. This problem is closely related to the uncertainty in the $Q^2$ scale. If the $Q^2$ scales for the two-jet and three-jet events are not identical, then Eq. (5) needs a correction due to the non-cancellation of the $\alpha_s$ factors and structure functions due to scaling deviations. Since this uncertainty cannot be eliminated until the appropriate theoretical calculations have been done, we quote results for $\alpha_s(K_{3I}/K_{2I})$, where $K_{3I}/K_{2I}$ represents the effect of the higher-order corrections. Over the full angular acceptance, we find $\alpha_s(K_{3I}/K_{2I}) \approx 0.23$.

On the assumption that the $Q^2$ scale for the three-jet events is lower than that for the two-jet events, we can compare the three-jet sample with a subset of two-jet events with smaller scattering angles. Thus for $Q^2 \approx 4000$ GeV$^2$, UA1 obtains from the 1983 data $\alpha_s(K_{3I}/K_{2I}) = 0.16 \pm 0.02 \pm 0.03$. (6)

It is clear that this value depends strongly on the angular acceptance chosen and care must be taken when comparing results from different experiments.

6. LIMITS ON QUARK COMPOSITENESS

Recently there have been theoretical speculations that quarks and leptons have compositeness structures and that they are bound states of more fundamental constituents called preons. Following the model of Eichten et al.\textsuperscript{4,10}, the energy scale of compositeness is defined as $\Lambda_c$. In this model, values of $\Lambda_c$ different from infinity (i.e. consistent with ordinary QCD) produce an excess of events in the high-p$_T$ region of the inclusive jet cross-section. To determine the value of $\Lambda_c$, we have renormalized the QCD calculation to the low-p$_T$ region. In all instances, the $\Lambda_c = \infty$ solution gave a good fit to the data. Taking into account theoretical and experimental uncertainties, the lower limit becomes $\Lambda_c > 400$ GeV.

One can also extract a limit on quark compositeness from the two-jet angular distribution. In contrast to the inclusive jet cross-section, the two-jet angular distribution is independent of the structure functions when scaling violations are ignored. Hence, in the limit of infinite statistics, a higher limit for $\Lambda_c$ from the angular distribution is intuitively expected. The high-mass ($M_{2J} = 240$–300 GeV) two-jet angular distribution was used to determine the value of $\Lambda_c$. Values of $\Lambda_c$ different from infinity produce an excess of events at small $\chi$ (see Fig. 3). The limit is obtained by fitting the angular distribution over the full $\chi$.
Taking into account the effect of the structure functions and $Q^2$ scales, as well as experimental uncertainties, we calculate a preliminary limit of $\Lambda_c > 400 \text{ GeV}$ at the 95% confidence level. This limit is comparable to that obtained by fits to the inclusive jet cross-section.

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A STUDY OF MULTI-JET EVENTS
AT THE CERN pp COLLIDER

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ABSTRACT

The UA2 experiment has performed a study of events containing three hard jets in the final state. The angular distributions of the three jets show evidence for gluon bremsstrahlung, in good agreement with a QCD model to leading order in the strong coupling constant $\alpha_s$. The yield of three-jet events relative to that of two-jet events provides a measure of the strong coupling constant: $\alpha_s K_3/K_2 = 0.23 \pm 0.01 \pm 0.04$, where $K_2$ and $K_3$ represent the contributions arising from higher order corrections in $\alpha_s$ to the two- and three-jet exclusive cross-sections. A preliminary study has also been performed on events containing four jets in the final state. The observed features of these events agree well with a recent leading order QCD calculation, indicating a strong deviation from phase space behavior. These results are not consistent with the presence of a large contribution from multi-parton scattering, in which more than one hard scattering takes place among the nucleon constituents.
1. Introduction

In recent years the CERN $pp$ Collider has proven to be an excellent laboratory in which to examine hard collisions between quarks and gluons. These collisions have been extensively studied and their properties have been successfully compared with the predictions of perturbative quantum chromodynamics (QCD). This is made possible by the presence of hadron jets which can be isolated from the other collision products and identified with the outgoing partons. In particular, most features of two-jet final states are well described by the leading term in the perturbative expansion of the parton-parton cross-section in powers of the strong coupling constant $\alpha_s$. In this report we discuss the observation [1,2] and present a detailed study [3] of the properties of the three-jet and four-jet final states which are proportional to higher powers of $\alpha_s$.

We divide this study into two areas of investigation. First, we examine the three-jet events and compare several distributions with a QCD model and with phase space, in order to test the underlying dynamics. These comparisons confirm the qualitative features expected of QCD, in particular, a Rutherford-like angular dependence for the scattering angle of the leading jet and a bremsstrahlung spectrum for the energy and angle of the softest jet. This agreement encourages us to perform more quantitative tests of QCD. To do so, we measure the cross-section for the production of three-jet events relative to that for two-jet events. This ratio provides a measurement of the strong coupling constant $\alpha_s$.

Second, we present a preliminary study of four-jet events. The features of these events are compared with a QCD model and with phase space with the goal of testing perturbative QCD to order $\alpha_s^4$. The event shapes are found to agree well with the QCD model, and a measurement of $\alpha_s$, using the ratio of the four-jet cross-section to the three-jet cross-section, is in progress. In addition, several authors have predicted a second mechanism for producing four-jet events, in which more than one hard parton scattering takes place in a single collision [4]. A simple model for this so-called multi-parton mechanism has been used to search for the presence of such events. The preliminary conclusion is that the observed events are consistent with leading order QCD, and there is no evidence for the multi-parton process in the current data sample.
2. The Data

The data for the present study have been obtained using the UA2 central calorimeter [5] which covers the pseudo-rapidity region $-1 \leq \eta \leq +1$ over the whole azimuthal range. Its granularity (240 cells, each covering 15° in azimuth and 10° in polar angle) is well matched to the jet structure of hard collision final states at Collider energies. The trigger is obtained from a coincidence between a minimum bias signal, ensuring the occurrence of a non-diffractive $pp$ collision, and the requirement that $\Sigma E_T$, the total transverse energy measured in the 240 calorimeter cells, exceeds 60 GeV. The data were collected at $\sqrt{s} = 630$ GeV and correspond to a total integrated luminosity of 310 nb$^{-1}$ from 1984 and an additional luminosity of 440 nb$^{-1}$ from 1985.

In each event, energy clusters are constructed using a refined algorithm optimized for the study of multi-jet final states [3]. We have studied the angular resolving power associated with this refined algorithm by superimposing jets obtained from two different two-jet events. The fraction of such events for which the clustering algorithm finds two clusters is a measure of the resolving power. It is found to have a distribution with an approximately energy-independent cut-off at $30° \pm 10°$ (Figure 1).

Once the clusters have been defined, we proceed to select a sample of events which arise primarily from hard scattering processes. This topic has been studied previously by UA2 in an analysis of the emergence of the hard scattering contribution to two-jet final states [6]. This analysis was based on a simple two component model for hadron collisions. The soft component is characterized by a fairly uniform distribution of energy among a large number of low energy clusters, and dominates at low total transverse energy. The hard component consists of two high transverse energy jets accompanied by an underlying event of low transverse energy secondaries, and dominates at high total transverse energy. The results of this study are summarized in Figure 2, which shows the inclusive cross-section and the fraction of hard scattering events as a function of two transverse energy variables: the total transverse energy ($\Sigma E_T$) and the leading jet transverse energy ($E_T^1 = E_T^1 + E_T^2$). This figure suggests that the contamination from the soft component decreases much more rapidly if we require a large transverse energy in the leading jets ($E_T^1$).
In each event, the clusters are sorted in order of decreasing transverse energy and labeled 1, 2, 3, etc. In addition, we calculate the transverse momentum $P_T$ of the jet system as the vector sum of the transverse momenta of the leading clusters. The events are then required to satisfy the following conditions:

\[ E_T^1 + E_T^2 + ... + E_T^n > 70 \text{ GeV}, \]  
\[ E_T^i > 10 \text{ GeV}, \quad E_T^{n+1} < 10 \text{ GeV}, \]  
\[ |\eta_i| < 0.80, \quad i = 1...n. \]
\[ P_T < 20 \text{ GeV} \]

The first condition ensures that the configuration of the leading jets is unbiased by the trigger threshold, and reduces the contamination from the soft collisions discussed above. The second condition retains events in which the leading clusters are likely to be associated with the hard collision rather than with spectator fragments. The third condition defines a fiducial volume in which the jet energy measurements are reliable. The fourth condition eliminates events with large missing momentum. Using these criteria, we define the two, three, and four jet samples which will be used in the subsequent analysis.

3. The QCD Model

The present study relies on comparisons between experimental distributions associated with the multi-jet sample and the corresponding QCD predictions. It is therefore useful to describe the QCD model and its implementation in a Monte Carlo simulation.

To leading order in $\alpha_s$, the cross-section for producing $n$ final state partons with a total invariant mass $\sqrt{s}$ is expressed in terms of elementary subprocesses in which two incident partons, $i$ and $j$, carrying fractions $x_i$ and $x_j$ of their nucleon parent momentum, interact:

\[ \sigma_n^{LO} = \left( \frac{\alpha_s^2}{\pi^2 s} \right) \sum_{i,j} F_i(x_i) F_j(x_j) Q_{ij}^n \Phi_n \left( \frac{dx_i}{x_i} \right) \left( \frac{dx_j}{x_j} \right) \]

where $\Phi_n$ is the n-body phase-space factor and $x_i x_j s = \hat{s}$.

Explicit expressions for the contribution $Q_{n}^{ij}$ of each elementary subprocess are available in the literature for $n = 2$ [7], $n = 3$ [8], and $n = 4$ [9]. In Equation (2), they are weighted according to the structure functions $F_i(x)$ which describe the parton content of the incident nucleons.
As a result of the bremsstrahlung nature of the gluon radiation spectrum, $Q_n^{|\vec{q}|}$ diverges when the mass of any two partons, one from the final state and one from either the initial or final state, approaches zero. These divergences are cancelled by non-leading contributions to the topological cross-sections. Since we are working with leading order calculations, we will only consider final states in which the initial and final state partons are well separated in phase-space.

We implement this requirement in our model by defining cut-offs to ensure that the parton configuration in the final state is free of these divergences: we only consider partons having a transverse momentum in excess of 8 GeV and parton pairs having a separation in excess of 20°. With these cut-offs, the ratio $K_n$ between the topological cross-section $\sigma_n$ and its leading order approximation $\sigma_n^{LO}$ can, in principle, be computed, although the necessary calculations have not yet been performed. Lacking a precise calculation of $K_n$ we shall have to compare the multi-jet data to expressions proportional to $K_n \sigma_n$, which means that the ratio of cross-sections is sensitive to the quantity $\alpha_s K_3/K_2$, instead of simply $\alpha_s$.

The quark and antiquark structure functions, $Q(x)$ and $\bar{Q}(x)$, are taken from low $q^2$ neutrino data [10] evolved to the $q^2$-range of the present experiment. The gluon structure function is taken to be

$$G(x) = F(x)/\sqrt{K - 4/9 \{Q(x) + \bar{Q}(x)\}}$$

(3)

where $F(x)$ is the effective structure function obtained from the two-jet data of the UA2 experiment [11]. In this evaluation, we assume that the inclusive $K$-factor takes the value $K = 2$. In addition, both the structure functions and the strong coupling constant $\alpha_s$ are functions of $q^2$. We choose to identify $\sqrt{q^2}$ with the largest transverse momentum among the final state partons.

In order to convert the final state partons of Equation (2) into observable hadrons, we use a fragmentation method based on the Field-Feynman algorithm [12], modified to reproduce the cluster radius distribution observed in the present experiment [1]. We have checked that this method of fragmentation, when used in conjunction with the clustering algorithm described previously, produces extra clusters in less than 2% of the events for cluster transverse energies in excess of 15 GeV. This empirically demonstrates that in our model the radiation of hard partons can be separated from the soft fragmentation process, and ensures that the jets in the final state arise from the hard-scattering matrix element and not from the fragmentation model.
The acceptance of the UA2 detector and the details of the energy response of the central calorimeter are simulated in a Monte Carlo program which reproduces the experimental details of relevance \cite{1}. The underlying event, associated with spectator particles, is simulated by superimposing actual minimum bias events onto the jets produced by the hard collision.

4. The Three-Jet Sample and a Comparison with the QCD Model

In this section we study the event configurations for a three-jet sample selected from the 1984 data sample of 310 nb$^{-1}$ according to criteria (1). After boosting to the center of mass of the jet system, the leading jets are arranged in order of decreasing $p_T$. The three-jet system can be described in terms of six variables, three defining the orientation of the plane containing the three-jet system in space, and three defining the configuration of the jets within this plane. In order to emphasize those features peculiar to QCD, the distributions of several variables will be compared with the predictions of the model described in the previous section, as well as with a phase space model obtained by setting the $Q_{nji}$ in Equation (2) equal to one.

We first examine the distribution of $\cos \theta^*$, the angle of the scattering plane relative to the beam direction. This distribution, normalized to one at $\cos \theta^* = 0$ and corrected for the limited UA2 acceptance, is shown in Figure 3. The figure also contains the corrected distribution for two-jet events \cite{11}. The shapes are very similar, and are characteristic of vector gluon exchange. The curve in Figure 3 shows the parton level calculation of the cross-section for the $gg \rightarrow ggg$ sub-process, computed using Equation (2) and the transverse momentum and angular separation cut-offs described in the preceding section. The curves for the other sub-processes are not significantly different, and the overall agreement with QCD is very good.

We next discuss the internal configuration of the three-jet system, independent of its orientation. The qualitative features can be seen in a Dalitz plot, constructed by defining scaled variables using the two-body masses $m_{ij}$:

$$x_{ij} = \frac{m_{ij}^2}{s}$$

(4)
Figure 4 shows the distribution of $x_{12}$ versus $x_{23}$. If the events were distributed according to a phase space density, the population would be uniform across the plot (ignoring acceptance effects). Instead, there is an increased density in the region of small $x_{23}$ (corresponding to final state bremsstrahlung of a soft third jet) compared to that of large $x_{23}$ (corresponding to equal sharing of energy among the three jets). The absence of events near $x_{12} = 1$ and $x_{23} = 0$ is a result of the event selection criteria (1) and of the two-jet angular resolving power. The projections of the Dalitz plot are also shown. They are compared with the QCD model (solid histogram) and the phase space model (dashed curve). The QCD model agrees well with the data, and there is a large excess of events above the phase space curve in the region of small $x_{23}$.

In the preceding analysis we used selection criteria which allowed us to display the dominant features of the three-jet data sample. In particular, loose cuts were used on the jet transverse momenta and on the opening angle of jet pairs (Figure 1) to study the effect of bremsstrahlung. In this way, we have shown that the QCD model gives a good description of the shapes of the three-jet distributions. We now pursue a more quantitative comparison with the QCD model. To this end, we measure the quantity $a_s K_3/K_2$ by adjusting $a_s$ in the QCD model until the theoretical value of the ratio between the three-jet cross-section and the two-jet cross-section ($R_{QCD}$) is equal to its experimental value, $R_{exp}$.

There are several sources of systematic error associated with this measurement which can be reduced by applying stricter event selection criteria. We want to avoid those regions of phase space where either two of the partons are separated by a small angle, or one of the partons is soft. For this reason, we redefine the two-jet and three-jet data samples using two additional criteria:

i. the standard clustering algorithm is followed by an additional step in which secondary cluster pairs, having an opening angle smaller than 50° and a transverse energy in excess of 5 GeV, are merged into a single cluster (Figure 1).

ii. the 10 GeV threshold used in (1b) is replaced by a 15 GeV threshold to further reduce the contamination from spectator particles.

The new selection criteria retain 14635 two-jet and 2596 three-jet events. The experimental value $R_{exp}$ is simply calculated as the ratio between these two numbers. This procedure gives a value $R_{exp} = 0.177 \pm 0.004$. 
The selection criteria applied to the data samples are also applied to the Monte Carlo event samples. The value of \( R_{\text{QCD}} \) is then the ratio between two cross-sections, \( \sigma_3^{\text{QCD}} \) and \( \sigma_2^{\text{QCD}} \). Due to the limited acceptance and resolving power of the UA2 detector, some \( n \)-jet final states will be observed as \( m \)-jet events, \( m < n \). This requires us to include in \( \sigma_3^{\text{QCD}} \) both events from \( \sigma_3^{\text{LO}} \) and from \( \sigma_2^{\text{LO}} \) which have three and only three jets obeying the selection criteria. Similarly, \( \sigma_2^{\text{QCD}} \) contains events generated from \( \sigma_2^{\text{LO}} \) and \( \sigma_3^{\text{LO}} \) which have two and only two jets obeying the selection criteria. At the time of this analysis, the full QCD calculation of \( \sigma_4^{\text{LO}} \) was not yet complete. An approximate estimate of \( \sigma_4^{\text{LO}} \) was obtained by taking Equation (2) for \( n = 3 \), and adding the bremsstrahlung of a single gluon. The estimate was then normalized to the total number of observed four-jet events. The relative contributions to \( \sigma_3^{\text{QCD}} \) from \( \sigma_4^{\text{LO}} \) and to \( \sigma_2^{\text{QCD}} \) from \( \sigma_2^{\text{LO}} \) are respectively 21% and 16%. Higher order contributions, such as from \( \sigma_5^{\text{LO}} \) to \( \sigma_3^{\text{QCD}} \) and from \( \sigma_4^{\text{LO}} \) to \( \sigma_2^{\text{QCD}} \) are ignored.

By varying \( \alpha_s \) in the QCD model, we find that the value of the strong coupling constant that makes \( R_{\text{QCD}} \) equal to \( R_{\text{exp}} \) is

\[
\alpha_s \frac{K_3}{K_2} = 0.23 \pm 0.01. \tag{5a}
\]

The effects of the dependence of \( \alpha_s \) on \( q^2 \) are displayed in Figure 5, which shows the values of \( \alpha_s \frac{K_3}{K_2} \) in different bins of the multi-jet mass \( M \). The \( q^2 \)-range covered by the data sample is from \(~600 \) to \( 10,000 \text{ GeV}^2 \), with the average \( q^2 \sim 1700 \text{ GeV}^2 \). The curve in Figure 5, obtained by computing \( \alpha_s \) for an average \( q^2 \) calculated for each mass bin in the figure, shows the expected \( \pm 7\% \) variation of \( \alpha_s \) over this mass range. The data are in agreement with the curve, but show no significant deviation from a constant value.

Several sources of systematic uncertainties affect our evaluation of the strong coupling constant. We now consider each of these in turn and evaluate their contributions to the error on \( \alpha_s \frac{K_3}{K_2} \).

i. Fragmentation effects: In the high energy range of the present experiment, there is virtually no ambiguity in distinguishing between two-jet and three-jet configurations. Instead, the main consequence of using different fragmentation functions is to modify the fraction of the original parton energy which is collected in the corresponding calorimeter cluster. This affects the three-jet sample, which often contains a relatively soft third jet, more than the two-jet sample. Changes in the fragmentation model produce variations of \( \alpha_s \) which do not exceed \( (\Delta \alpha_s / \alpha_s) \approx \pm 10\% \).
ii. Underlying event: The spectator scattering contribution is simulated by superimposing a minimum bias event onto each generated hard collision. However, we know [1] that the value of \( \langle E_T \rangle \) is, on average, twice as high for hard collisions as for minimum bias events. Part of the difference can be accounted for by particles directly associated with the hard collision. Nevertheless, we have studied the effect of doubling the transverse energy carried by the underlying event. From this study we evaluate an uncertainty on \( \alpha_s \), \( \Delta \alpha_s/\alpha_s = \pm 10\% \).

iii. Structure functions: Uncertainties in the structure functions may affect the \( \alpha_s \) measurement to the extent that the quark to gluon ratio may be modified. Because of the higher colour charge of gluons, \( \alpha_s \) relative increase of the gluon content results in a larger value of \( R_{QCD} \) and consequently in a lower value of \( \alpha_s \). Changes in the quark to gluon ratio result in an uncertainty \( (\Delta \alpha_s/\alpha_s) = \pm 7\% \).

iv. Higher order contributions: The contribution to \( R_{QCD} \) from \( \alpha_s^{LO} \) is important, of the order of 21\%. As mentioned earlier, \( \alpha_s^{LO} \) was computed without the benefit of the full QCD calculation. We estimate an uncertainty of \( \pm 40\% \) on the absolute scale of \( \alpha_s^{LO} \) which appears in the \( \alpha_s \) measurement as an error \( (\Delta \alpha_s/\alpha_s) = \pm 8\% \).

v. Energy response: Uncertainties in the energy response of the UA2 central calorimeter partially cancel since they affect both the two-jet and the three-jet samples. They introduce an uncertainty on the \( \alpha_s \) measurement of \( (\Delta \alpha_s/\alpha_s) = \pm 3\% \).

vi. Event selection: The measurement of \( \alpha_s \) should not depend on the criteria used to select the jet samples, as long as these selection criteria are compatible with the cut-offs used in the QCD model. After varying the selection, we infer an upper limit \( (\Delta \alpha_s/\alpha_s) < 4\% \) on the uncertainty resulting from this source.

After this summary of the individual uncertainties, we can compute the total systematic error. Some effects are accounted for in several of the systematic uncertainties described above. We conservatively ignore this double counting and retain as a global systematic uncertainty the number obtained by adding all individual contributions, i) to vi), in quadrature. The result is

\[
\alpha_s K_s/K_2 = 0.23 \pm 0.01 \pm 0.04
\]
where the second error represents the systematic uncertainty.

Finally, we note that our choice of $q^2$ in the scaling violation factors is somewhat arbitrary. A different choice of $q^2$ will not affect $\alpha_s$, but only modifies the QCD scale $\Lambda$, as long as the two-jet and three-jet samples are treated equally. However, it is possible to choose $q^2$ definitions which alter the average $q^2$ of the two-jet sample with respect to that of the three-jet sample. We have considered a concrete example by replacing our choice, $\sqrt{q^2} = \text{Max}(p_T^2)$, by $\sqrt{q^2} = <p_T^2>$, thereby reducing the three-jet $<q^2>$ with respect to the two-jet $<q^2>$. The result is to increase only the value of $\alpha_s$ that appears in the calculation of $\alpha_s^{LO}$, typically by 8%, and hence to increase the value of $R_{QCD}$. As a consequence, the measured value of $\alpha_sK_3/K_2$ decreases by 25%.

5. A Study of Four-Jet Events

In this section, we describe an analysis of four-jet events selected from a combined data sample of 750 nb$^{-1}$. The selection is made according to criteria (1), with an additional step in the clustering algorithm in which secondary clusters with an opening angle smaller than 30° are merged into a single cluster. This reduces the sensitivity to fragmentation effects while still retaining good angular resolution, and results in a sample of 2219 events. A typical high $\Sigma E_T$ event from this sample is shown in Figure 6, with four clearly resolved jets in the final state. If we apply the stricter cuts, used in the previous section for the $\alpha_s$ measurement, we arrive at a much smaller sample of 255 events.

For this analysis, we again boost to the center of mass of the jet system and arrange the jets in order of decreasing $p_T^*$, The four jet system can be described in terms of nine variables: three describing the orientation of the jet system and six describing the internal configuration of the jets. This large number of variables make the description of this final state rather complex. The jets do not lie in a plane, nor can they be described by simple Dalitz plots. A natural set of variables are the six $x_{ij}$ defined in Equation (4), or equivalently the space angles between the jets in the center of mass, $\cos \omega_{ij}^*$. Only five of these variables are independent — the sixth degree of freedom corresponds to the total mass of the jet system.

As for the three-jet analysis, comparisons will be made between the data, a parton level QCD model based on recent calculations by Kunszt and Sterling [9], and a phase space model obtained by setting $Q_{n ij}$
in Equation (2) equal to one. At this preliminary stage of the analysis, we will only study the shapes of the distributions – questions of normalization will be avoided by setting the area of the distributions to one.

To examine the qualitative features of the events, we first look at the sphericity calculated from the jet momenta in the center of mass. The distribution is shown in Figure 7, and it is apparent that the observed event shape agrees well with the QCD model, and is significantly less spherical than phase space. A more detailed understanding of the features of four-jet events can be obtained by examining the space angles defined previously. Figure 8 shows the distribution of \( \cos \omega_{12} \). The data shows a significant enhancement above the phase space model in the region of small angular separation, indicating the presence of bremsstrahlung. This enhancement agrees very well with the QCD model. A similar, less dramatic, enhancement can also be observed in the distribution of \( \cos \omega_{14} \).

Several authors have predicted an additional source of four-jet events [4]. This is the multi-parton mechanism, in which multiple hard parton collisions take place in a single hadron collision. This mechanism, which is suppressed by extra powers of \( s \), becomes important for small values of the parton momentum fraction, where the density of partons in the nucleon becomes very large. The simplest form of this process, in which two independent pairs of partons interact, can be described as follows:

\[
\sigma_{\text{mp}} \sim \int G(x_1, x_2) G(x_3, x_4) \, d\hat{\sigma}_{12} \, d\hat{\sigma}_{34}
\]

where the \( G \) are double structure functions and the \( d\hat{\sigma} \) are the cross-sections for the two parton sub-processes. In order to search for evidence of multi-parton interactions, we have constructed a simple model in which the double structure functions take the form:

\[
G(x_1, x_2) \sim F(x_1) F(x_2)
\]

We simulate the effects of soft gluon radiation from the initial state partons by giving each jet-jet system a Gaussian \( P_T \) kick. The magnitude of this kick has been adjusted to agree with the measured two-jet data [11]. Using this model, one can then look for deviations from the leading order QCD prediction, choosing variables which should be sensitive to a multi-parton contribution.
The most characteristic feature of multi-parton events should be the appearance of pair-wise correlations among the jets. To search for these correlations, we choose to look at transverse variables in the lab frame, since they are relatively insensitive to center of mass motion. A simple variable of this kind is the difference between the azimuthal angle of the leading jet, and the azimuthal angles of the other jets in the event. This variable, called $\phi_{\text{lead}}$, is sensitive to the presence of a second jet opposite to the leading jet. The distribution of this variable, with three entries per event, is shown in Figure 9a. The data agree well with the QCD model, and show no sign of the narrow peak that is expected for the multi-parton process. An alternate variable is the $P_T$ unbalance in the event, defined by the expression:

$$P_T^{\text{unb}} = 2 \min \left( \sum_{i} P_T^{i} \right)$$

where $i$ is chosen to minimize the unbalanced $P_T$. This variable should take on small values for the multi-parton process since there will be a second jet balancing the $P_T$ of the leading jet. The observed distribution is shown in Figure 9b, and agrees well with the QCD model. Again, there is no sign of the enhancement expected at small $P_T$ unbalance from the multi-parton mechanism.

A slightly more sophisticated approach involves finding the pairing of the four jets which minimizes the $P_T$ sum for pairs of jets:

$$\min \left( \sum_{i} P_T^{i} \right)$$

Once the jets have been paired, it is possible to calculate the difference in azimuthal angle between the jets in each pair. The resulting distributions are shown in Figure 10, where $\phi_{\text{pair}}$ is always the difference for the higher $E_T$ pair. The pairing process itself introduces some correlations, but there is no evidence for the narrow peak expected in $\phi_{\text{pair}}$ from the multi-parton interaction. Once again, the observed events are consistent with the leading order QCD calculation.

Two preliminary conclusions emerge from this discussion. The first is that the observed four-jet distributions agree well with a leading order QCD calculation, and differ significantly from four-body phase space. The second is that there is no evidence for the additional contribution from the higher twist multi-parton processes. These studies of four-jet events are being extended to include a more quantitative test of the QCD model using the ratio of the four-jet to three-jet cross-sections, as well as an attempt to find multi-parton interactions at lower $\Sigma E_T$. 
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Figure 6

Figure 7

Figure 8
Figure 9

Figure 10
INCLUSIVE JET CROSS SECTIONS AND
PRODUCTION IN THE UA2 DETECTOR

The UA2 Collaboration

Bern - CERN - Copenhagen (NBI) - Heidelberg - Orsay (LAL) -
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ABSTRACT

We report on measurements of very large $p_T$ jets in the UA2 experiment at the CERN pp Collider for $\sqrt{s} = 546$ GeV and $\sqrt{s} = 630$ GeV including a preliminary analysis of the latest data taking period at the end of 1985. The emergence of two jet dominance has been studied using shape analysis techniques and the conclusions agree with previously published results. The inclusive jet production cross-sections exhibit a $p_T$ dependent increase with $\sqrt{s}$; this behaviour can be described by QCD calculations and by approximate $x_T$ scaling. No structure at high mass is seen in the 2-jet mass cross-section and no significant deviation of the data from QCD predictions is observed at very large $p_T$, allowing to place a lower limit $\Lambda_C > 370$ GeV at 95% C.L. on the characteristic scale $\Lambda_C$ of a hypothetical superstrong contact interaction.
1. Introduction

The CERN \( \bar{p}p \) Collider is an excellent laboratory for the study of short distance (hard) collisions between quarks and gluons. Hadronic jets resulting from these collisions were clearly identified using the UA2 calorimeter during the first Collider run in 1981\(^1\) and their features have been the subject of intensive studies by the UA2 \(^2\) and UA1 \(^3\) collaborations. All these results strongly support perturbative Quantum Chromo–Dynamics (QCD) as a predictive theory of strong interactions, where \( \bar{p}p \) interactions are interpreted as a hard scattering of partons which subsequently fragment into final state hadrons.

We would like to review here the latest results on inclusive jet physics in the UA2 experiment after the Collider run at \( \sqrt{s} = 630 \) GeV has taken place at the end of 1985, and compare these results with the values measured at lower \( \sqrt{s} \), at the Collider and at the ISR, and with QCD calculations. Special care will be devoted to a detailed investigation of the emergence of two jet dominance at Collider energies using a shape analysis of the global energy flow in the UA2 calorimeter.

2. Apparatus and data taking

The UA2 detector has been described in detail in other publications\(^4\); the data for the present study have been obtained using the high granularity central calorimeter which covers the pseudo–rapidity region \( |\eta| < 1 \) over the whole azimuthal range. It is composed of 240 cells, each segmented into an electromagnetic and two hadronic compartments and covering 10° in \( \theta \) and 15° in \( \phi \), and it is built in a tower structure pointing to the center of the interaction region. The response of the calorimeter has been studied for electrons, muons and hadron beams from 1 to 70 GeV and the resolution \( \sigma_E/E \) is 0.14/\( \sqrt{E} \) for e.m. showers and varies from 32\% at 1 GeV to 11\% at 70 GeV (decreasing as \( \approx E^{-1/4} \)) for single pions.

All triggers used are calorimetric triggers made by summing the transverse energy deposited in a single wedge of \( \Delta \phi = 120^\circ \) ("inclusive jet" trigger) or in two opposite azimuthal wedges of \( \Delta \phi_1 = 60^\circ \) and \( \Delta \phi_2 = 120^\circ \) ("two jet" trigger) or in the total acceptance of the calorimeter ("\( \Sigma E_T \)" trigger), with given thresholds. Each of them is in coincidence with a "minimum bias" trigger, obtained from scintillator arrays on both sides of the interaction region in the angular range \( 0.44^\circ < \theta < 2.84^\circ \), ensuring the
presence of a non-diffractive collision. This trigger is also used to monitor the luminosity integrated in the different run periods. The collected luminosity values are $\mathcal{L} = 120 \text{ nb}^{-1}$ in 1983 at $\sqrt{s} = 546 \text{ GeV}$, $\mathcal{L} = 310 \text{ nb}^{-1}$ in 1984 at $\sqrt{s} = 630 \text{ GeV}$, both with a systematic error of $\pm 8\%$. A preliminary value for 1985, also at $\sqrt{s} = 630 \text{ GeV}$ is $\mathcal{L} = 442 \text{ nb}^{-1}$ with approximately the same error.

3. Two jet dominance

Since the beginning of the Collider era it has been observed that with increasing $\Sigma E_T$ (the total transverse energy of all final state particles in a given rapidity interval) more and more events consist mostly of two hadron jets having approximately opposite transverse momenta. In a QCD picture this large $\Sigma E_T$ interaction is described as a short distance "hard" collision between two of the incident partons which subsequently fragment into hadron jets, whilst the remnant nucleons (spectators) independently produce a number of "soft" secondaries. In contrast with this large $\Sigma E_T$ regime the overwhelming majority of $\bar{p}p$ interactions are "soft" collisions that produce final state where $\Sigma E_T$ is more uniformly distributed among the collision products. As a function of $\sqrt{s}$ hard and soft collisions behave differently; the former obey approximate Bjorken scaling 5) while the latter evolve typically as $\ln s$. As a consequence the transition between the two regimes takes place in a region which depends upon $\sqrt{s}$ 6).

Fig. 1 shows the $\Sigma E_T$ cross-section as seen in UA2 where the transition between the soft and hard regimes is visible as a change of slope in the distribution and takes place in the region $50 \leq \Sigma E_T \leq 90 \text{ GeV}$.

In order to analyze in detail the features of the events over the whole range of $\Sigma E_T$ and to investigate the behaviour of the transition region we have applied a methodological scheme well known in $e^+e^-$ experiments 7) by using collective shape variables to describe the pattern of the energy flow inside the central calorimeter. For each event we build energy vectors pointing from the origin of the detector to the center of each cell having $E_T^{\text{cell}} > 500 \text{ MeV}$ and we select a suitable data sample according to two general criteria:

- given the limited rapidity coverage of the detector, the distortions arising from edge effects should be minimized: this is accomplished by limiting the longitudinal and transverse momenta of the event and requiring the direction of the maximum energy flow to be well-contained in the calorimeter.
• an adequate sampling of the available phase space is necessary in order to get reliable values of the shape variables: this translates into a minimal requirement on the number of energy vectors in each event.

We then boost the events into their center of mass system and we calculate various shape variables and jet measures. Fig. 2 shows the dependence of the mean value of two of them upon $\Sigma E_T$. The variables are defined as \[ S = \frac{(3/2)\min \sum (p_i \cdot \hat{e})^2}{\Sigma |p_i|^2} \]

\[ H_2 = \frac{1}{E} \sum |p_i| \cdot P_2(\cos \theta_{ij}) \]

where $\hat{e}$ is the "jet axis", $P_2(\cos \theta_{ij})$ the 2nd Legendre polynomial in the angle between each pair of energy vectors, $E$ is the total energy and the sums run over all the energy vectors $p_i$ in the event.

As can be seen, $<S>$ and $<H_2>$ start from the values expected for an almost spherical configuration at low $\Sigma E_T$ (high $S$, low $H_2$) and evolve towards values characteristic of collimated 2-jet structures (low $S$, high $H_2$) at high $\Sigma E_T$, passing through the region between 50 and 90 GeV where a transition between these two different patterns takes place.

It is possible to give a pictorial representation of the energy flow for these events using a technique which is also well known in $e^+e^-$ experiments. To this aim we first look for the Thrust axis $\hat{e}_1$:

\[ \text{Thrust} = T = \max \sum (p_i \cdot \hat{e}_1)/\Sigma |p_i| \]

which identifies the direction of the maximum energy flow, then we project the energy vectors in the plane orthogonal to $\hat{e}_1$ and we find the direction which maximizes the energy flow in this plane ("Major" axis $\hat{e}_2$). To treat the events on the same basis we can now apply an ordering algorithm superimposing for each event the T axis and the corresponding $\hat{e}_2$ axis. In fig. 3 is plotted the Energy Flow in the Thrust-Major plane together with differential distributions of $H_2$ for different slices of $\Sigma E_T$: a) 10-25 GeV, b) 40-60 GeV, c) 60-90 GeV, d) 160-210 GeV. The radial scale is different for each plot and is given in units of $\text{GeV}/10^5$ per event. The impressive evolution of the multi-lobe pattern with increasing $\Sigma E_T$ is clearly seen, with a corresponding change in the $dN/dH_2$ distribution. This evolution starts with configurations that reflect the almost random direction of the $\hat{e}_1$ axis expected from a Phase Space dominated final state and the ordering procedure followed (10 $< \Sigma E_T < 25$ GeV). Then, crossing the transition region, structures which are typical of collimated 2-jet events arise, shrinking more and more at the highest $\Sigma E_T$ values (160 $< \Sigma E_T < 210$ GeV).
4. Inclusive cross-sections

In order to reduce background contamination we reject events coming from the jet triggers if they do not have a good pattern of energy deposition in the calorimeter (see ref. 2 for details) or if they are associated with an early signal in the minimum bias counters. These criteria reduce this contamination to less than 2%, while the loss of good events is < 3%.

In order to find a correspondence between hadron jets and the energy deposited we apply a direct cluster algorithm (which is possible due to the high segmentation of the central calorimeter) joining cells which share a common side and have an energy > 400 MeV. Moreover we retain only clusters with |η| < 0.85 and, to partly account for final state gluon radiation, \(^9\) the jet momentum is modified by adding to it the vectors of all the clusters having \(E_T > 3\) GeV and separated by an angle \(\cos \omega > 0.2\). In the presence of such clusters the jet acquires a mass and its \(P_T\) differs from its \(E_T\).

The cross-sections for the reactions

\[
\begin{align*}
pp &\rightarrow \text{jet + anything} \\
pp &\rightarrow \text{jet + jet + anything}
\end{align*}
\]

have been evaluated from the observed number of jets using the integrated luminosity \(\mathcal{L}\) for normalization and a Monte Carlo simulation to account for detector acceptance. The results are summarized for \(\sqrt{s} = 546\) GeV and \(\sqrt{s} = 630\) GeV in fig. 4 for inclusive jet production \(d^2\sigma/dP_Td\eta\) at \(\eta = 0\) and in fig. 5 for two jet production. The errors include the statistical error and an energy dependent systematic uncertainty coming from the acceptance functions. An additional systematic error of 45%, arising from an energy independent uncertainty of the acceptance functions (35%), calorimeter calibrations (30%) and luminosity (8%), is not included in the plots.

We have also performed a systematic search for structures at high masses (> 100 GeV/c\(^2\)) in the jet-jet invariant mass cross-section for the 2-jet sample only. As illustrated in fig. 6 no evidence for any structure is observed. The results for a fit with a smooth exponential background plus a gaussian signal give 52±44 events at 162±6 GeV/c\(^2\), without any statistical significance.
5. Comparisons with QCD and lower $\sqrt{s}$

As can be seen in fig. 4 the preliminary analysis of the 1985 data agrees very well with previously published results at $\sqrt{s} = 630$ GeV (the errors on the 1985 data are almost identical to those of 1984 data) and a clear increase of the jet production cross-sections with $\sqrt{s}$ is visible in fig. 4 and 5. In order to examine this increase more closely the ratio of cross-sections at the two $\sqrt{s}$ has been evaluated and is shown in fig. 7 after suitable rebinning. Only statistical errors are given because the systematic errors approximately cancel in the ratios.

The inclusive jet cross-sections have also been measured in pp collisions at the ISR. These results are compared in fig. 8 to the present ones in terms of scaled invariant cross-section $p_T^*E_0/dp^3$ as a function of $x_T = 2p_T/\sqrt{s}$. Scale breaking effects can be parametrized in the form

$$E_0/dp^3 = p_T^{-n} \cdot f(x_T), \quad f(x_T) = A(1-x_T)^{m/x_T}$$

as shown in fig. 9. A global fit gives $n = 4.74 \pm 0.06, m = 6.54 \pm 0.15$.

All these measurements can be compared with QCD calculations: the curves shown in fig. 4, 5 and 7 have been obtained from a leading order (in $\alpha_s$) QCD calculation using $Q^2 = p_T^2$ as a scale ($\Lambda = 0.2$ GeV) and the structure functions of ref. 10. This prediction describes the cross-sections at both collider energies and their increase with $\sqrt{s}$, even if the $\pm 45\%$ systematic error on the measurements precludes an accurate evaluation of the effects coming from higher order $\alpha_s$ contributions ($K$ factors).

The inclusive jet cross-section data at $\sqrt{s} = 630$ GeV have been analyzed in terms of deviations from QCD including the effects of a hypothetical super-strong interaction binding preons inside quarks, following the parametrization of ref. 11 with the choice $g^2/4\pi = 1$ for the coupling constant. In this model a substructure of partons would manifest itself as a contact interaction visible at momentum transfers well below the characteristic energy scale $\Lambda_c$. Finite values of $\Lambda_c$ would produce an excess of events compared to ordinary QCD predictions ($\Lambda_c = \infty$) at large $p_T$ values. Hence the analysis is possible to the extent that the main uncertainties (systematic errors and ignorance of the $K-$factor) are approximately constant over the $p_T$ range, so that we can normalize the data in the low $p_T$ region and look for deviations in the high $p_T$ tail.
In fig. 4 and 5 the results are plotted for the pure QCD calculation ($\Lambda_c = \infty$) and for the best fit ($\Lambda_c = 460$ GeV) together with the expected behaviour for $\Lambda_c = 300$ GeV as an illustration of the sensitivity to $\Lambda_c$. When both theoretical and experimental uncertainties are taken into account we obtain a lower limit of $\Lambda_c = 370$ GeV at 95% C.L. which agrees with recently reported results from the UA1 Collaboration 12).

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UA2 PRELIMINARY
1983-1985 data
2-jet only

Fig. 6

Fig. 7

Fig. 8

Fig. 9
CHARGED PARTICLE MULTIPLEITIES IN $\mu^+\mu^-$ ANNIHILATION AT 29 GeV: MEASUREMENTS IN THE CENTRAL RAPIDITY REGIONS AND IN THE GLUON JET

HRS COLLABORATION
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Abstract

The charged particle multiplicity distribution for $\mu^+\mu^-$ annihilations at $\sqrt{s} = 29$ GeV has been measured using the High Resolution Spectrometer at PEP. Firstly, the multiplicity distributions for particles selected in different central rapidity spans are presented. All of these are well represented by the negative binomial distribution. As the rapidity span is narrowed, the distributions become broader and approach a constant value of the parameter $k$.

Secondly, the charged particle multiplicities of the quark and gluon jets in the three-fold symmetric $e^+e^- + q\bar{q}g$ events have been studied. A value of $\langle n \rangle = 6.7 \pm 1.1 \pm 1.0$ for gluon jets with an energy of $9.7 \pm 1.5$ GeV is measured. The ratio, $\langle n \rangle_q/\langle n \rangle_g$, is $1.29 \pm 0.21 \pm 0.20$, which is significantly lower than the value of $9/4$ naively expected from the ratio of the quark-to-quark color charges.

1. **Introduction and Data Analysis**

In this paper we report a study of the charged particle multiplicity in $e^+e^-$ annihilation at $\sqrt{s} = 29$ GeV. First we present the multiplicity distributions for particles selected in different central rapidity spans.\(^1\) Then we present the mean charged particle multiplicity of a gluon jet.\(^2\) The data sample which this analysis is based on corresponds to an integrated luminosity of 185 pb\(^{-1}\) and was obtained with the High Resolution Spectrometer (HRS)\(^1,2\) at PEP. The HRS is a solenoidal spectrometer that measures charged particles and electromagnetic energy over 90% of the solid angle.\(^1,2\)

The events were selected to have $n$ well reconstructed tracks where $5 < n < 40$. The sum of the charged particle momenta, plus the energy observed in the shower counter, was required to be greater than 12 GeV, and more than 7.5 GeV/c in the scalar sum of the charged particle momenta. To ensure good tracking efficiency, the thrust axis of the event was selected to be within 30° of the equatorial plane of the detector, and each track had to have an angle with respect to the $e^+e^-$ beam direction of more than 24° and had to register in more than one-half of the drift chamber layers traversed. With these selections, the reconstruction efficiency for isolated tracks was greater than 99%. For a typical annihilation event, with several close neighboring tracks, the reconstruction efficiency is 80% or better and varies slowly with dip angle; for the higher momenta, $p > 2$ GeV/c, this efficiency increases to 90%.

The true multiplicity distribution was determined from the observed data by means of a transposed matrix technique. The matrix was determined from a Monte Carlo simulation of the experiment which includes the effects of the experimental cuts as well as the tracking inefficiencies. The results were found to be stable for reasonable variations, both of the cuts and of the definition of a good track. Events with no charged particle in the selected rapidity span were included in the data sample.

2. **Rapidity Dependence**

The KNO multiplicity distributions for the two-jet data sample and for particles contained in selected rapidity ranges from $|Y| < 0.1$ to $|Y| < 2.5$, corresponding to $Y$ spans from 0.2 to 5.0 units, are shown in Fig. 1. Each successive data set has been displaced lower by a factor of ten for clarity. The two-jet data with no $Y$ selection are also shown with their own ordinate scale. These events are always even prongs so the normalization differs by a factor of two from that of the data with a rapidity selection. The distributions clearly widen as the rapidity span is restricted, and for $Y \leq 1$ events with $z$ values of 3 to 5 are seen. The data of Fig. 1 have been fitted to

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\(^1\)\(^2\)
the negative binomial distribution

\[ P(n, <n>, k) = \frac{k(k + 1) \ldots (k + n - 1)}{n!} \left( \frac{<n>/k}{1 + <n>/k} \right)^n \left( 1 + \frac{<n>}{k} \right)^{-k}. \]  

(1)

In the parametrization, the shape of the distribution is determined by the parameter \( k \) and the position of the maximum by \(<n>\). The variable \( k \) is related to the mean multiplicity \(<n>\) and the dispersion \( D \) by

\[ \frac{D^2}{<n>^2} = \frac{1}{<n>} + \frac{1}{k}. \]  

(2)

The histogram in Fig. 1 show the results of the fits to Eq. (1). It is remarkable how well the negative binomial distribution represents the data. In all cases, the value of \( k \) calculated from the measured value of \( D \) and \(<n>\) agrees to within a few percent that given by the fit of the data to Eq. (1).

The values of \( k \) obtained by fitting the data of Fig. 1 to the negative binomial are shown in Fig. 2(a), connected by the solid angle. The \( k \) value falls rapidly as the \( Y \) limit is restricted but levels off for the central rapidity region at values of \( k \) between 5 and 10. The lower data set, connected by the dashed line, corresponds to the full data sample without the two-jet selection. The overall variation of \( Y \) is similar for the two data sets. Fig. 2(b) shows the (two-jet) mean charged particle multiplicities (scale on left) and dispersion values (scale on right) for the different \( Y \) selections. The lines are drawn to guide the eye.

The variation of \( k \) with \( Y \), shown in Fig. 2(a), differs qualitatively from the observation of the UA5 group who studied
pp interactions at $\sqrt{s} = 540$ GeV and found a continuous broadening of the multiplicity distributions as the pseudorapidity range was narrowed: no tendency to a limiting shape was seen. However, the $\bar{p}p$ distributions are broader, characterized by $k$ values falling from $\sim 3.8$ to 1.5 as the pseudorapidity span is restricted.

Bialas and Hayot have fitted the TASSO data for the full rapidity range in a model with independent emission of clusters; they predict a limiting KNO distribution for the central rapidity region. This expectation is supported by our data and so may be taken to confirm the dominant role of energy and momentum conservation in narrowing the $e^+e^-$ KNO distribution for the full rapidity span. The narrow multiplicity distribution observed for the data in the outer rapidity spans also supports this conclusion.

Chao, Meng, and Pan also have fitted our data in various rapidity spans in a statistical model. Their model has no free parameters. Nevertheless the fits are very good.

Several authors have interpreted the hadronic data on cluster models. Giovannini and Van Hove have also considered the applications for $e^+e^-$ annihilation, and, in a cascade model, the cluster charged particle multiplicity, $<n>_c$, is controlled by a parameter $b = <n>/<(n + k)$ being given by $<n>_c = \frac{6}{b-1} \frac{1}{\ln(1-b)}$. The present results give $b$ values in the range 0.15 and 0.3 as the rapidity span is changed corresponding to $<n>_c$ values between 1.1 and 1.2. The 540 GeV $\bar{p}p$ data by contrast gives $<n>_c \sim 3$ when interpreted in this model.

3. Gluon Jet

The events passing the cuts described previously are mixtures of the two- and three-jet topologies. The low sphericity region ($0 < S < 0.25$) is dominated by the two-jet events, and contains 82% of the data sample; the higher sphericity region ($S > 0.25$) is strongly enriched in three-jet events. In order to select a clean sample of the latter, a jet-finding algorithm was applied to the hadronic data sample in the kinematic region with $S > 0.15$ and $A < 0.1$, where $A$ is the aplanarity. These selections in $A$ and $S$ ensured that the three jets are in a plane. The three-jet events were then selected by a cut on the normicity variable ($C_3$), which is the ratio of the triplicity to thrust. For normicity $> 1.05$, 89% of the $qqg$ partonic states are selected by the algorithm; the estimated contamination in this three-jet data sample from two-jet events is 16%.

The normalized particle flow of these three-jet events on the event plane is shown in Fig. 3(a). In this plot the jets are ordered according to the
angles between neighboring jets; jet 1 is defined as that jet opposite to the smallest angle and, similarly, jet 3 is opposite to the largest angle. The prediction of the Lund Monte Carlo model [9], shown by the line, agrees with the data. With our selections, according to the Monte Carlo study, about 56% of the third jets result from gluon fragmentation.

The multiplicity distribution of each jet was determined by the matrix method described previously. The corrected mean multiplicities for the individual jets as a function of their energies are shown in Fig. 4. The errors are dominated by systematics coming from the contamination from two-jet events, the estimated track-finding efficiency, and the uncertainty of the correction matrix. The solid line is the fit to half of the mean multiplicity for all $e^+e^-$ annihilation events plotted as a function of $\sqrt{s}/2$. Since $e^+e^-$ annihilations in the energy range shown in Fig. 4 are dominated by $e^+e^- + q\bar{q}$, this line shows the expected $\langle n \rangle$ values for quark fragmentation averaged over all flavors.

The multiplicities derived from the three-jet events agree with this curve; no evidence is seen for higher $\langle n \rangle$ values even though jet 3 is enriched in gluons.

The three-fold-symmetric three-jet events were selected by the three-jet finding algorithm with cuts: $S > 0.15$, $A < 0.1$, $C_3 > 1.10$, and the angle between the jet axes in the range $100^\circ$ and $140^\circ$. The sample of 276 events passing these cuts has a total mean multiplicity of $16.3 \pm 0.3 \pm 0.7$ and dispersion of $4.2 \pm 0.5 \pm 0.3$. According to the Monte Carlo simulation, the energy of each jet is 9.7 GeV with a width of $+1.5$ GeV and $-2.0$ GeV measured at the half maximum value. We note that the line on Fig. 4 gives $\langle n \rangle = 5.2$ for a quark jet of this energy.

The three jets are now ordered according to their charged particle multiplicities, with jet 1 defined as the jet with the lowest number of charged
particles and jet 3 with the highest multiplicity. Figure 3(b) shows the particle flow on the event plane with this ordering. The angle is now measured with respect to the axis of jet 1 in the direction of jet 2. The areas in the three peaks are then proportional to the mean particle multiplicities. The ratio of the mean multiplicity of jet 3 to that of the average of jet 1 and jet 2 is $2 < n_3 > / < n_1 + n_2 > = 1.71 \pm 0.06 \pm 0.05$. Because of the multiplicity ordering, we expect this ratio to be greater than one even if the multiplicity of the gluon jet is the same as that of the quark jet.

In order to compare this ratio with the values expected for different multiplicities of the gluon jet, we made a simple simulation of the experiment in which each 9.7 GeV jet with $< n > = 5.2$ particles was allowed to independently fragment with a Poisson distribution adjusted to fit the observed single jet multiplicity distributions. This model predicts $2 < n_3 > / < n_1 + n_2 > = 1.69$ for $< n >_g = < n >_q = 5.2$ and $2 < n_3 > / < n_1 + n_2 > = 2.31$ for $< n >_g = 9/4 < n >_q = 11.7$. The corrected multiplicity distribution for the jet 3 sample is shown in Fig. 5. The solid line, which agrees reasonably well with the data, shows the expectation from this simple Poisson model with $< n >_g = < n >_q = 5.2$, and the dashed line corresponds to that with $< n >_g = 9/4 < n >_q = 11.7$. If $< n >_g$ is allowed to vary, then a value of $< n >_g = 6.7 \pm 1.1 \pm 1.0$ at a jet energy of $9.7 \pm 2.0$ GeV is found to best reproduce the jet 3 multiplicity data and the measured ratio of $2 < n_3 > / < n_1 + n_2 >$. This value gives a ratio $< n >_g / < n >_q$ of $1.29 \pm 0.21 \pm 0.20$. We conclude, therefore, that in the energy range of our experiment, the mean multiplicity of the gluon jet, for the particles with $x = p / p_{\text{beam}} > 0.01$, is consistent with being equal to that of the quark jet, and significantly different from the value of $9/4$ coming from the color factor ratio. This result is consistent with the measurement by the MARK II collaboration.

There are several uncertainties that cloud these comparisons of quarks and gluon jet multiplicities. First, some of the charged particles in the quark jets result from the decay of the mesons containing the heavy quarks (c and b). The multiplicity measurements of jets resulting from heavy quark fragmentation gives an increase of $\sim 0.5$ particles per quark jet from this
effect. Second, the next-to-leading order QCD corrections to the ratio of quark to gluon multiplicities gives a reduction of about 10%.\textsuperscript{13} At high jet energies these calculations, therefore, predict a $<n_g>/<n_q>$ ratio of two. Third, the very low momentum particles are difficult to assign to a specific jet.

The simulation of $e^+e^-$ annihilation, with our experimental selections, based on the Lund string model, predicts a total mean multiplicity of 16.7 and $2 <n_3>/<n_1 + n_2> = 1.64$ in agreement with the data. In this approach, the gluon is considered to be a kink in the string joining the $q$ and $\bar{q}$, and so the $9/4$ color factor does not explicitly enter. At infinite energy, the dual string topology at the gluon kink leads to a rapidity density which is twice that at the quark ends of the string. The difference between these expectations and the agreement of our data with the Lund simulation reflects the kinematic constraints coming from the low energy of the experiment.

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STUDIES OF PARTON FRAGMENTATION AND BARYON PRODUCTION WITH THE TPC AT PEP

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ABSTRACT

We present here results from the TPC on the analysis of 3-jet events, on the width of gluon jets versus quark jets, and on proton and lambda production and correlations.

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1. INTRODUCTION

We report here the analysis of data taken with the Time Projection Chamber\(^1\) (TPC), which provides three-dimensional tracking and dE/dx particle identification, at the PEP electron-positron colliding beam ring. The center-of-mass energy was 29 GeV. The total number of multi-hadronic events in the data is about 29,000 from the 1982-83 run, and about 14,000 from the 1984-85 run. During this latter run the TPC had improved momentum resolution due to the use of a 13.25 kG superconducting coil, and the reduction of drift field distortions with a gated grid\(^2\).

The topics covered in this paper are:

1. Introduction
2. Tests of Parton Fragmentation Models Using 3-jet Events
3. Width of Gluon Jets vs. Quark Jets
4. Baryon Production
5. Conclusions

2. TESTS OF PARTON FRAGMENTATION MODELS USING 3-JET EVENTS

The three main types of models for the fragmentation of quarks and gluons into observed hadrons, shown schematically in Figure 1, all have similar predictions for 2-jet events. They differ, however, in their predictions for the distribution of particles between jet axes in 3-jet events.

In the Independent Fragmentation (IF) model\(^3\), each parton fragments independently with cylindrical symmetry in the overall center-of-mass (CM) system. Thus all regions between the jets are populated equally.

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**Figure 1** Schematic representation of a 3-jet event in the
(a) Independent Fragmentation Model
(b) String Fragmentation Model
(c) Cluster Fragmentation Model

**Figure 2** The ratio \(N_{31}/N_{12}\) of the population between jets, for the data and models.
In the String Fragmentation (SF) model\(^4\), it is assumed that strings are stretched between the partons along the directions of color flow. Each string fragments in its own rest frame with cylindrical symmetry. The hadrons are boosted to the CM system, thus depleting the region between the q and \(\bar{q}\) of hadrons (relative to the gq and \(\bar{q}g\) regions). Due to its boost origin, this relative depletion effect is enhanced by selecting heavy particles or those particles with a large momentum component out of the event plane (\(P_{\text{out}}\)).

In the Cluster Fragmentation (CF) model\(^5\), the primary partons initiate a quark-gluon shower described by leading-log QCD. After the parton shower evolution, color singlet clusters are formed from neighboring partons, which then decay into hadrons. Soft gluon interference produces an effective angular ordering of gluons, which forces partons in the forward direction along the jet axis. In 3-jet events, this interference effect gives less partons in the \(q\bar{q}\) region than in the gq and \(\bar{q}g\) regions\(^6\).

The 3-jet event selection is described in Reference\(^7\). The jets are labeled 1, 2 and 3 such that jet 1 is opposite the smallest angle between jets and jet 3 is opposite the largest angle. Monte Carlo studies indicate that jet 3 is the gluon jet in about 55% of the events. To search for a depletion of particles in the \(q\bar{q}\) region (i.e. between jets 1 and 2), relative to the gq region (i.e. between jets 3 and 1), one calculates \(N_{31}/N_{12}\), where \(N_{ij}\) is the number of hadrons between jets \(i\) and \(j\). For IF models we expect \(N_{31}/N_{12} = 1\) independent of the particle mass or \(P_{\text{out}}\), while for the SF and CF models we expect this ratio to be greater than 1 and to increase in magnitude as mass and \(P_{\text{out}}\) increase. Figure 2 shows \(N_{31}/N_{12}\) for the data and models. It is seen that the IF model fails to describe the data, while the SF and CF models do agree with the data. The failure of the IF model is fundamental, and cannot be explained by parameter tuning or by using different variants of the model.

3. WIDTH OF GLUON JETS vs. QUARK JETS

In the SF model (fig. 1b), the gluon has two strings stretched between it and the initial quarks. Each string fragments into hadrons, and thus one would expect that the average transverse momentum (in the event plane) of the hadrons in the gluon jet would be greater than in the quark jets. We have used our 3-jet sample to look for this effect, and find that the ratio of the average transverse momentum (in the event plane) of jet 3 (usually the gluon jet) to jet 2 (the \(\bar{q}\) jet) is \(1.08 \pm 0.02\) (preliminary). The SF model predicts 1.10 - 1.11 for this ratio, in agreement with the data. As expected, the IF model gives 0.99 - 1.01 for this ratio, indicating that the effect is not just kinematical. Thus the gluon jet appears to be "fatter" than the quark jets.
4. BARYON PRODUCTION

We can use baryon angular distributions and transverse momentum correlations to help distinguish between the baryon production models shown in Figure 3.

In the "fundamental diquark" model, it is assumed that diquarks are fundamental entities, and thus diquark-antidiquark pairs can occasionally be formed in the color field (instead of $q\bar{q}$ pairs, which lead to meson production). The combination of a diquark and a quark forms a baryon (fig. 3a).

In the "effective diquark" model, it is assumed that occasionally $q\bar{q}$ pairs of non-screening color will "pop" out of the color field. In order to screen the remaining field, another $q\bar{q}$ pair is produced (fig. 3b). These quarks form loosely bound diquarks, which then combine with another quark to form a baryon. In this model, however, it is possible for more than one $q\bar{q}$ pair to be created in the remaining color field, leading to the production of a meson "in between" the baryon and antibaryon (fig. 3c).

In the "cluster decay" model, parton showers create low-mass color singlet clusters, some of which may be heavy enough to decay into a baryon and antibaryon (fig. 3d).
The diquark and cluster models both reproduce the measured inclusive proton spectra quite well, but the model predictions differ for the angular dependence of proton-antiproton correlations. If we define the angle $\Theta^*$ as the angle between the proton momentum and the jet axis in the CM frame of the $pp$ pair, then in the cluster model, since the cluster decays spherically symmetric, the distribution in $\cos \Theta^*$ will be flat. In both diquark models, the proton and antiproton are more likely to be produced along the jet direction because the diquarks and antidiquarks are pulled apart by the tension in the color string. In this case the angular distributions would show an enhancement at $\cos \Theta^* = \pm 1$ (fig. 4a).

The $\cos \Theta^*$ distribution of $pp$ pairs measured by the TPC is shown in Figure 4b, together with the predictions of the diquark model and cluster decay model. Because of the limited range of the proton momentum (0.5 - 1.5 GeV/c), the predictions are lower at large $\cos \Theta^*$ than expected, but the qualitative difference between the two models is maintained. The diquark models (solid curve) are consistent with the data, while the cluster model (dashed curve) disagrees with the data and is excluded at greater than 95% CL. Thus baryon pairs are oriented primarily along the jet axis, and are not produced isotropically in the CM frame of the pair.

![Figure 4](image)

Figure 4 Distribution of $pp$ pairs in the angle $\Theta^*$ between the proton direction and the sphericity axis, measured in the $pp$ rest frame. a) Predictions of the LUND diquark model (solid line) and of the Webber cluster model (dashed line). b) Experimental points and model predictions, as in a), for momenta between 0.5 and 1.5 GeV/c.
Angular correlations between the $p$ and $\bar{p}$ momentum transverse to the jet axis can discriminate between the two variants of the diquark model\textsuperscript{12}. In the "fundamental" diquark model, the diquark and antidiquark will be produced with opposite transverse momenta. In the "effective" diquark model, when a meson is produced between the baryon and antibaryon, this correlation is largely destroyed. We define a correlation coefficient, $\alpha$, proportional to the dot product of the proton and antiproton transverse momentum vectors out of the event plane. Using the "effective" diquark model, one can calculate the value of $\alpha$ as a function of the probability of a $\text{BMB}$ configuration (i.e. a meson between the baryon and antibaryon). This is shown in Figure 5, along with the value of $\alpha$ measured with the TPC data. It is seen that the probability of a $\text{BMB}$ configuration is greater than 65% at 90% CL (preliminary). Thus the "effective" diquark model can account for a large part of baryon production.

We can also use baryon correlations to investigate whether baryon number conservation is local or global. First, all multihadron events having an antiproton in the rapidity range of 0.2 - 0.9 are selected. Then for these events we plot the number of protons minus the number of remaining antiprotons as a function of rapidity. The results are in Figure 6, and show an excess of protons at about the same rapidity that was selected for the antiproton, indicating that baryon number conservation is local. The solid curve in Figure 6 is the diquark model prediction, and the dashed curve is just the inclusive rapidity distribution for protons (which one would get if there were no baryon correlations).

![Figure 5](image1.png)

Figure 5: Correlation coefficient $\alpha$ of $p,\bar{p}$ momentum components. Shaded bands: data $\pm 1$ SD. Full line: model prediction as a function of the probability to find a $\text{BMB}$ configuration instead of $\text{BB}$.

![Figure 6](image2.png)

Figure 6: Baryon density ($N_p - N_{\bar{p}}$) after selecting out an antiproton in the rapidity range of 0.2-0.9. Solid curve is the diquark model prediction; dashed curve is the inclusive proton rapidity distribution.
The TPC also sees evidence for local baryon number conservation in events with lamdas and antilamdas, as the $\Lambda \bar{\Lambda}$ tend to be in the same jet rather than opposite jets, as the following preliminary results show:

<table>
<thead>
<tr>
<th></th>
<th>Same Jet</th>
<th>Opposite Jet</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\Lambda \bar{\Lambda}$</td>
<td>21 events</td>
<td>6 events</td>
</tr>
<tr>
<td>$\Lambda \Lambda$</td>
<td>0</td>
<td>2</td>
</tr>
<tr>
<td>$\bar{\Lambda} \bar{\Lambda}$</td>
<td>1</td>
<td>1</td>
</tr>
</tbody>
</table>

5. CONCLUSIONS

The String Fragmentation Model and the Cluster Fragmentation Model satisfactorily predict the relative depletion of hadrons in the region between the $q$ and $\bar{q}$, while the Independent Fragmentation Model fails to describe the data.

We have presented preliminary evidence that gluon jets are "fatter" than quark jets.

We have tested the predictions of various baryon production models using angular correlations between protons and antiprotons, and find that diquark models agree with the data, and cluster models do not. Transverse momentum correlations indicate that the "effective" diquark model can account for a large part of baryon production.

Baryon correlations ($p\bar{p}$ and $\Lambda \bar{\Lambda}$) indicate that baryon number is conserved locally.

We acknowledge the efforts of the PEP staff, and the engineers, programmers and technicians of the collaborating institutions who made this work possible. This work was supported by the United States Department of Energy, the National Science Foundation, the Joint Japan-United States Collaboration in High Energy Physics, and the Foundation for Fundamental Research on Matter in The Netherlands.
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EXPERIMENTAL STUDIES ON MULTIJET PRODUCTION IN $e^+e^-$ ANNihilation AT PETRA Energies

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Abstract.
Jet production of the multihadronic data in the whole PETRA energy range of 14.0 to 46.7 GeV has been studied in terms of a cluster algorithm. The relative cluster multiplicities of the data are presented and compared with model calculations, based on 2nd order QCD perturbation theory as well as on leading log QCD calculations including soft gluon interference. The 2nd order perturbative QCD model produces too few spherical and 4-jet like events at the higher energies. The observed discrepancy in the ratio of 4-jet to 3-jet production cannot be explained in the framework of 2nd order perturbative QCD calculations. This result has also consequences for the value of the strong coupling constant $\alpha_s$ determined in previous analyses. The leading log QCD model provides a good description of the yield of spherical, 4-jet like events of the data. The differences between the two types of models are also observed on the level of quarks and gluons, and thus are most likely due to the different QCD calculations.
In the year 1982, the JADE collaboration published the first experimental evidence for the observation of 4-jet like events in $e^+e^-$ annihilations around 33 GeV c.m. energies\cite{1}. In particular, data required the production of $(7.2 \pm 1.2)\%$ 4-jet events additional to the 1st order 2- and 3-jet QCD models. The corresponding theoretically expected 4-jet rate due to second order QCD calculations was 5\%, depending on the value of $\alpha_s$. Since that time, no specific analysis on the production of multi-jet events in $e^+e^-$ annihilation has been published. There were only indirect statements that the experimentally observed rate of 4-jet like events is higher than predicted by QCD models\cite{2}. However, experimental results on rates of observable multi-jet event production (jet multiplicities) at the PETRA energy range are of interest for tests of different QCD calculations. In this analysis the experimental jet multiplicities, defined by a cluster algorithm, are presented and are compared to the predictions of different QCD model calculations.

A detailed description of the JADE detector, the trigger conditions and of the selection of hadronic events are given in \cite{3}. In this analysis, the data which have been taken until the end of 1985 at 14 GeV, 22 GeV, from 32.0 to 36.7 GeV and from 40.0 to 46.7 GeV center of mass energies, are included. Both charged and neutral particles with momenta exceeding 100 $MeV/c$ and 150 $MeV/c$, respectively, are used in the analysis. In addition to the selection cuts mentioned in \cite{3}, further cuts have been applied in order to reject the background from 2$\gamma$ processes as well as events with hard initial state $\gamma$-radiation or significant particle losses around the beampipe.

We compare the data to the predictions of Monte Carlo models which are based on different types of QCD calculations for the production of quarks and gluons and on different phenomenological fragmentation schemes for the transition of quarks and gluons into final state hadrons. The Lund string fragmentation model \cite{4} is used in connection with the 2$^{nd}$ order perturbative QCD calculations of Gutbrod, Kramer and Schierholz (GKS) \cite{5}. In these calculations, however, certain 2$^{nd}$ order corrections to the 3-jet production had been neglected. New evaluations by Gottschalk and Shatz \cite{6} showed that these terms give sizable corrections for 3-jet events at large thrust values. Recently, the complete 3- and 4-jet matrix elements based on this new calculations became available \cite{7} and have been installed into the Lund fragmentation model. We found that $\alpha_s$ has to be decreased by $\approx 10\%$ compared to the model with the GKS matrix elements, but the overall description of the data in general does not change.

In this analysis, we basically present results of the Lund model with the original GKS matrix elements, which we will refer to as "2$^{nd}$o.MC", but will also discuss results of the version with the Gottschalk and Shatz matrix elements separately. The fragmentation parameters of the Lund model are those as deduced in previous analyses on energy correlations and the determination of the strong coupling constant $\alpha_s$ \cite{8}, on $K^-$ and $\rho^0$ production \cite{9}, and on a combination of measurements of the heavy quark fragmentation functions \cite{10}. In particular, $A_{QC\ell D}$ was set to 500 $MeV$, and $y_{min}$, the lower cutoff in the scaled invariant masses between any pair of partons in the event generator, defining the separation between the produced jet classes, is taken to $y_{min} = 0.015$ in order to provide a good description of the data\cite{8}. In the case of perturbative QCD, the relative production of 2-, 3- and 4-parton events is determined by the integrated matrix elements and thus depends only on the scale parameter $A_{QC\ell D}$ for a fixed definition cut $y_{min}$.

The second model we compare to our data is the QCD shower model of Webber and Marchesini \cite{11}, based on leading log QCD approximations including soft gluon interferences.
The parameters of this model have been mainly optimized by requiring a good description of the multiplicity and momentum distributions of the charged particles, in particular $\Lambda_{QCD}$ was set to 300 MeV. With these parameters, the model provides a reasonable good overall description of the data, besides that it produces too few 3-jet like events. It should, however, be noted that recent changes\textsuperscript{1} of the model [12] increase the number of visible 3-jet like events, so that also in this respect the data may be well described. For the changes of this model are still under study, in this analysis we will mainly present the results of the original version, which we will refer to as "l. log MC", and will only shortly discuss the effects of the recent changes.

All generated Monte Carlo events have been tracked by a computer simulation program through the JADE detector and underwent the same selection procedure as applied to the data.

The cluster algorithm we use in this analysis in order to define and reconstruct jets works in the following way: In each event, the two particles $i$ and $j$ with the smallest scaled invariant mass $y = M_{ij}^2 / E_{i\ell\ell}$ are combined to one "cluster" by adding the two 4-vectors, if $y$ is smaller than a fixed cutoff $y_{cut}$. This procedure is repeated until all possible combinations of particles or clusters satisfy the relation $y \geq y_{cut}$, and the resulting number of clusters is called the jetmultiplicity of this event. For calculating the invariant mass $M_{ij}$ we use the expression

$$M_{ij}^2 = 2 \cdot E_i \cdot E_j \cdot (1 - \cos \Theta_{ij}),$$

which is the definition of the invariant mass in case of massless clusters. Though the clusters, by adding 4-vectors of particles, in general have invariant masses of the order of a few GeV, we chose the described expression for $M_{ij}$ in order to obtain the closest agreement between the definition of massive clusters and massless partons at comparable $y$-values. This has been investigated in studies of Monte Carlo events, where the relation between partons and particles is known.

In table 1 are shown the resulting cluster-multiplicities for data and model calculations at $E_{cm} = 34$ GeV with $y_{cut} = 0.040$, which corresponds to minimum invariant jetmasses of 6.8 GeV and is a reasonable choice defining experimentally resolvable jets.

<table>
<thead>
<tr>
<th>&lt; $E_{cm}$ &gt; 34GeV</th>
<th>Data</th>
<th>2ndo. MC</th>
<th>l. log MC</th>
</tr>
</thead>
<tbody>
<tr>
<td>2-cluster</td>
<td>56.1 ± 0.4</td>
<td>53.2 ± 0.3</td>
<td>58.5 ± 0.4</td>
</tr>
<tr>
<td>3-cluster</td>
<td>40.2 ± 0.4</td>
<td>44.0 ± 0.3</td>
<td>37.8 ± 0.4</td>
</tr>
<tr>
<td>4-cluster</td>
<td>3.75 ± 0.16</td>
<td>2.85 ± 0.12</td>
<td>3.69 ± 0.15</td>
</tr>
<tr>
<td>$R_4/R_3$</td>
<td>0.093 ± 0.004</td>
<td>0.065 ± 0.003</td>
<td>0.098 ± 0.004</td>
</tr>
</tbody>
</table>

Table 1. Cluster multiplicities obtained with the described cluster-algorithm at $y_{cut} = 0.040$ (numbers in %).

\textsuperscript{1}The energy splitting on the first $q - \bar{q}$ vertex has been changed to the appropriate Altarelli-Parisi function and the event cascade is generated in a Lorentz system with a smaller opening angle between the first $q\bar{q}$ pair.
The 2nd order perturbative QCD model, though $\alpha_s$ is optimized to the data in different analyses, produces too many 3-jet like events and too few 4-cluster events. In particular, the ratio of 4- to 3-cluster events $R_4/R_3$ cannot describe the data. Using the 2nd order QCD matrix elements of refs. [6,7], we find a similar behaviour, and $R_4/R_3$ with a value of $0.062 \pm 0.004$ is too low to describe the data. The 3-cluster rate of the leading log QCD model comes out too low, but the rate of reconstructed 4-cluster events and the ratio of 4- to 3-cluster events describes data well within the statistical errors.

![Cluster-multiplicities of the data and model calculations at $E_{cm} = 34\,\text{GeV}$, determined with the cluster algorithm described in the text, as a function of $y_{cut}$.](image)

In order to test the structure of the different jet classes, fig. 1 shows the cluster-multiplicities of data and the models as a function of $y_{cut}$. It can be seen that the discrepancy in the 2- and 3-jet production between data and both models is almost the same in the $y_{cut}$ range between 0.015 and 0.060. The fraction of 4-cluster events predicted by the 2nd order QCD model falls below the data especially at high values of $y_{cut}$, whereas the leading log QCD model describes the data remarkably well. For the background of 2- and 3-parton events to the reconstructed 4-cluster events decreases with increasing $y_{cut}$ and is less than 25% at $y_{cut} = 0.06$, the deficiency of 4-cluster events at high $y_{cut}$-values seems to be a severe problem of the 2nd order QCD predictions for the production of 4-parton events.

Ordinary analyses of the strong coupling constant are influenced by this fact in the way that an increased production of 3-jet events, which means a higher value of $\alpha_s$, partly cancels the effects of missing 4-parton events in the analyzed observables. This is the reason why in
fig. 1 the 3-cluster rate of the $2^{nd}$ order QCD model, with $\alpha_s$ determined in an analysis of the asymmetry of energy-correlations [8], lies systematically above the data and matches the data only at high values of $y_{\text{cut}}$, where most of the original 4-parton events are reconstructed as 3-cluster events and are not resolved separately any longer. Consequently, the matching of the $2^{nd}$ order model 4-cluster fraction to the data at small $y_{\text{cut}}$-values is caused by the increased number of 3-parton events which, due to the dominant effects of fragmentation in this regime, have a high probability to be reconstructed as 4-cluster events.

Determinations of $\alpha_s$ by adjusting the rate of observable 3-cluster events in the region of $y_{\text{cut}} = 0.03$ to 0.05, where they are largely decoupled from the production of 4-jet events, result in values which are 10 % less than what has been obtained in ordinary analyses of the strong coupling constant. Consequently, the deficiency of observable 4-cluster events gets even more pronounced in this case. Similar results we obtained using the $2^{nd}$ order QCD matrix elements of refs [6,7]. The leading log QCD model, if modified as described above, may provide a good description of the observed 2-, 3- and 4-cluster multiplicities of the data within the whole range of $y_{\text{cut}}$ presented in fig. 1.

Fig. 2 Cluster-multiplicities of the data and model calculations as a function of $S = E_{cm}^2$, determined with the cluster algorithm described in the text. $y_{\text{cut}}$ was set to 0.240, 0.096, 0.040 and 0.024 at $E_{cm} = 14, 22, 34$ and 44 GeV, respectively.

In fig. 2 the observable cluster-multiplicities of data and model calculations are shown for different c.m. energies between 14 GeV and 44 GeV. Here, $y_{\text{cut}}$ at each energy bin was set to $y_{\text{cut}} = (6.8\text{GeV})^2/E_{cm}^2$ in order to keep the same mass-definition of "resolvable jets". At 14 GeV only 2-jet like events and at 22 GeV mainly 2- and also some 3-jet events are observed. The observed relative difference between the data and both model calculations is present in the whole energy range. The rate of observable 4-cluster events increases by more than a factor of two going from 34 GeV to 44 GeV center of mass energy.

An important question is whether the observed differences in the cluster-multiplicities between data and models, as well as between the two models itself, are due to uncertainties
in the phenomenological fragmentation parts of models, or whether they indicate differences of the QCD calculations used in the models. For the Lund model, a detailed study on the possibility of changing model parameters has been done, though there is not much freedom left due to the measurements and adjustments [8-10] mentioned above. No possibility has been found in order to increase the observable 4-jet and to decrease the 3-jet rate at the same time. A comparison between the 2nd order and the leading log QCD models on the level of quarks and gluons (partons) has also been done in order to evaluate whether the observed differences in jet-multiplicities are already present at this stage and thus are due to the different QCD calculations implemented in these models. Such a comparison cannot be done directly, for the leading log model produces mostly multi-parton configurations with up to 10 quarks and gluons, whereas the 2nd order QCD model generates only 2-, 3- and 4-parton events. Therefore a procedure has been adopted to recombine those partons of the cascade of a leading log model event, which have a lower invariant mass than a certain cutoff \( y_{\text{min}} \) (the same definition as described for the cluster algorithm has been applied). This procedure is equivalent to an earlier stop of the parton shower cascade at the given \( y_{\text{min}} \)-cutoff. The remaining recombined partons are then comparable to those of the 2nd order generator with the same value of \( y_{\text{min}} \).

\[ E_{\text{CM}} = 34 \text{ GeV} \]

**Fig. 3** Parton-multiplicities at \( E_{\text{cm}} = 34 \text{ GeV} \) of the 2nd order and the leading log QCD models as a function of \( y_{\text{min}} \).

In fig. 3, the parton-multiplicities of the 2nd order QCD model and the recombined parton-multiplicities of the leading log model are shown as a function of \( y_{\text{min}} \) at \( E_{\text{cm}} = 34 \text{GeV} \). The difference between both models is qualitatively the same as observed in the comparison of data and the fully fragmented and tracked model calculations. From these observations we conclude that the fact that both models cannot describe the experimentally observed rate of jet-events, is already due to the parton dynamics of the models and thus reflects the imperfections of the different QCD calculations.
Further studies using different cluster algorithms and definitions of jet-classes, as the sphericity tensor, always resulted in the same conclusions as presented above. A detailed analysis of acoplanarity distributions, which are especially sensitive to the production of spherical and 4-jet like events, also showed that the experimentally observed production of 4-jet like events is inconsistent with the expectations of 2nd order perturbative QCD calculations [13].

**SUMMARY**

The relative cluster multiplicities of the data as a function of the mean c.m. energy squared as well as of the minimum invariant cluster masses have been presented. The results are compared to model calculations based on 2nd order QCD perturbation theory and on leading log QCD calculations including soft gluon interference. The 2nd order QCD model produces too few spherical, 4-jet like events. At the same time, this model overestimates the production of 3-jet like events. This is due to the fact that in recent studies on the strong coupling constant $\alpha_s$, the missing rate of spherical and 4-jet like events has been partly compensated by an increased production of 3-jet events. The observed discrepancy in the ratio of 4-jet to 3-jet production cannot be explained in the framework of 2nd order perturbative QCD. The leading log QCD model provides a better description of the observed rates of 4-jet events, but underestimates the production of 3-jet events. The differences between the two types of models have also been observed comparing the event structures on the level of quarks and gluons. From these observations we conclude that the discrepancy between the models and the data are due to the different QCD calculations.

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Determination of $\alpha_s$

Using the Planar Triple Energy Correlation
and the Energy Correlation Asymmetry

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The strong coupling constant $\alpha_s$ is determined with two independent dynamical variables, which are shown to be free from infrared divergence, to be $0.12 \pm 0.07 \pm 0.01$ at 44 GeV using hadronic events from electron-positron annihilation.
Determination of $\alpha_s$ Using the Planar Triple Energy Correlation and the Energy Correlation Asymmetry

A. From 3-JETS to $\alpha_s$ determination.

Electron-positron annihilation into hadrons is a fruitful testing ground for Quantum Chromodynamics (QCD). In 1979 it was discovered\(^1\) at PETRA that about 10% of the events have most of their energy concentrated in three distinct and narrow cones. The angular separation between pairs of the cones is much larger than the width of each cone. This discovery of three jet events and the identification of this phenomenon with hard and isolated gluon bremsstrahlung from the quark or the antiquark provides the simplest and strongest argument for the validity of QCD.

In QCD, gluons are the carriers of interactions between quarks. The strength of this interaction is measured by the coupling constant, $\alpha_s$, commonly defined in the $\overline{\text{MS}}$ scheme\(^2\).

Since the discovery of gluon jet, much effort has gone into measuring $\alpha_s$. The chronological compilation of all the data excluding the total cross section measurement, from PEP and PETRA on the measurement of $\alpha_s$ to the second order of QCD is shown in Fig. 1. These conflicting results once raised doubts\(^4\) on the accuracy of $\alpha_s$ determination using the $e^+e^-$ annihilation data. Since 1983, we\(^5\) and others\(^6\) have shown in detail how the apparent discrepancy originated from the following sources:

1. the rather trivial approximations made by all experiments with the exception of MARKJ. This conclusion is confirmed by the recent results of other groups\(^7\);

2. the use of some variables such as thrust and $R_\perp$ both of which depend strongly on the details of fragmentation models and cut-off values, without citing the proper systematic errors due to this dependence;

3. the data on energy asymmetry has large corrections ($\approx 50\%$) due to radiative effects and detector resolution. Most data, with the exception of MARKJ, were fitted and presented in the form of the corrected data. There is a sizeable correction\(^8\) to the corrected data when the parameters used to correct the data are significantly different from the values obtained from the fit. I would like to warn people not to use the corrected data without including this correction to the corrected data.

After these corrections, the data of all experiments agree within errors. Since so far only one infrared cut independent variable, the energy asymmetry\(^10\), has been used in the determination of $\alpha_s$, the main question is whether there is another independent method which is also infrared independent and fragmentation insensitive, to measure $\alpha_s$ and thus check the systematic errors.

B. Experimental Variables.

It is necessary to choose methods and experimental variables which are insensitive to the model details. Although the variables commonly used are variations of the same theme, measure of the 3-jet rate, they differ in the degree of dependence on the fragmentation details, the infrared cut-off and the sensitivity to higher order corrections. Linear quantities such as the planar T.C., the asymmetry functions, and the oblateness are less sensitive to the details of the Monte Carlo (such as multiplicity) than are quadratic variables, for example, $<P^2_\perp>$. The latter depends strongly on the multiplicity
of the gluon jet which is poorly known. Furthermore, since finite order perturbative calculations after incorporating fragmentation, require non-physical "cut-offs" in the infrared and collinear particle configurations, it is vital that the results do not show significant dependence on these cuts. Otherwise the systematic error will be at least as large as the cut dependence itself. This has been a source of confusion in the discrepancy among earlier measurements of $\alpha_s$. Since different variables have different dependence on the higher order QCD, QED and detector corrections, the difference in the results obtained using different infrared cut independent variables shows the size of possible systematic errors in the calculation.

C. Triple Energy Correlation and Infrared Cut Independence.

For each hadronic event we measure the energy deposited in the individual calorimeter element. We then compute the triple products of the measured energy vectors versus the spatial angles $\xi$ between two of the three energy pairs. The planar triple energy correlation (P.T.C.)\(^{(9)}\) is defined by,

$$ I = \frac{\sum \sum \sum E_i E_j E_k}{\sum \sum \sum E_{\text{vis}}^2} $$

where the sum is over all hadronic events; $E_i$ is the energy measured in the solid angle element $i$; $E_{\text{vis}}$ is the measured total event energy; $\xi$ is the angle between $E_j$ and $E_k$ and $\delta \xi$ is the range of integration of $\xi$, etc.; the three energy vectors $E_i$, $E_j$, and $E_k$ are approximately in a plane, i.e. when $\xi < \pi - \xi_1 - \xi_2 - \xi_3 \leq \theta_p = 0.1 \text{ rad.}$

The P.T.C. has the following properties:

1. a detailed event by event analysis is not required, just as in the case in the energy correlation asymmetry.

2. it is a mathematically independent quantity (its correlation with energy asymmetry is less than that between asymmetry and thrust), hence provides a new method to determine $\alpha_s$ reliably;

3. it is a natural measure of three jets as all three jets are treated equally. This is quite different from the energy asymmetry, when one jet is always integrated over; and

4. the correction due to radiative and detector effects is small.

The integrated P.T.C., $I(\beta > 40^\circ)$, with the integral is carried out with $\xi_1$, and $\xi_2 < 180^\circ - \beta$ and $\xi_1 + \xi_2 > 180 + \beta$ has only a small contribution from the two jet events, and measures the contribution of three jet events. Our results are insensitive to the exact values of the cuts, in the range for $50^\circ > \beta > 30$ and $0.12 > \theta_p > 0.08 \text{ rad.}$

We have calculated several QCD measures, both before and after the imposition of the fragmentation model and the detector simulation. We show three such measures in Fig. 2, where the solid curves are the cross sections evaluated from the fragmented hadrons after detector simulation. The horizontal scale is the relevant cutoff used in the QCD calculation. One commonly used measure, thrust, shows considerable dependence upon the cutoff value used at both the parton level and after fragmentation and detector simulation. The other two measures, the P.T.C., $I(\beta = 40^\circ)$ and the asymmetry of the energy energy correlations\(^{(10)}\), shows a remarkable stability over the entire range, and thus are suitable for the determination of QCD parameters.

D. Calculation Procedure and Data Analysis.
The procedure used by MARKJ is:

1. We begin with a complete calculation of second order QCD(11), via the generation of 2, 3 and 4 parton final states following Kuntz and Ali(12). The 3 parton state includes terms linearly dependent on $\alpha_s$ and terms quadratically dependent on $\alpha_s$. This allows the simultaneous calculation of QED correction, currently carried out to the third order of $\alpha$, as well as the fitting for $\alpha_s$;

2. Soft parton and collinear parton emission have been included in the fragmentation process. In order to avoid infrared divergences and the double counting due to fragmentation we require that either the Sterman-Weinberg parameters, the energy ($E=\text{energy of parton}/\sqrt{S}$) and angles between the partons ($\delta$) or the invariant masses of each pair of partons ($M$) must be above some minimum values. Those partons which fail the cuts are combined either with the closest parton which passes the cuts, or with a parton passing the cuts and the pair having the smallest invariant mass. The combination can be carried out either by adding up their momentum or their energies. Our results do not depend on the choice of the method used. All other quantum numbers are preserved in the combination and the subsequent fragmentation processes. The soft parton cannot be discarded since in the data, one measures hadrons not partons, thus cannot throw the corresponding quantity away. The 2, 3 and 4 parton states are kept separately until step 4;

3. A fragmentation model is used to transfer partons into hadrons which are simulated through the detector.

4. The relative weights of 2, 3 (linear), 3 (quadratic) and 4 parton states depend on $\alpha_s$. The comparison of the resultant distribution with data using a fit determines $\alpha_s$.

It is important to carry out the complete integral in one step since the integral of the products is not equal to the product of the integrals when the variables over which one integrates are correlated.


Since the partons to be fragmented are light in comparison to the masses of the hadron jets resulting from the fragmentation, in order to conserve momentum and energy in the procedure, energy must be transferred from one jet to another, a flow which is typically in the direction of the least energetic parton. We consider two different models for fragmentation, which are primarily distinguishable in how this is done. In models such as that of Ali, et al.(14), a Lorentz boost is performed to restore momentum balance, effectively transferring energy within the full event. In the string model(15), strings ending on partons are the objects fragmented, which transfers energy between the partons at the two ends. We use the difference in results from these two models to estimate the uncertainty due to modelling.

F. Results and Discussion.

The essence of the MARKJ results are the use of two variables to analyze the data over a wide range of energies. Fig. 3 shows the measured P.T.C. as a function of $x_2$ for several different values of $x_1$ for the high energy (39.7 to 46.78 GeV) and lower energy (35 GeV) data. The histograms are Monte Carlo simulations using a complete second order QCD calculation ($A=100$ MeV) for both the Lund and Ali fragmentation models. Both models fit our data equally well. The effects of QCD radiative corrections, finite detector acceptance, and resolution are imposed upon the models to facilitate direct comparison with the data. The P.T.C., integrated over $x_1$ and $x_2$ as in Eq. 4 for $\beta=40^\circ$, is shown in Fig.4 as a function of $\sqrt{S}$. The predictions of the Ali and Lund fragmentation models with, $A=100$ MeV, are shown to be good descriptions of the data.
Several fits were performed, varying the region in $\chi_1$ and $\chi_2$ and values of $\beta$. In all cases a $\chi^2$ approximately equal to the number of degrees of freedom is obtained with a small variation in the central value. We estimate a systematic error of 0.006. In Fig (5), the resulting values of $\alpha_s$ from the utilization of both P.T.C. and the asymmetry function are displayed, using data samples centered at 22, 35, 41, 44 and 46 GeV. The low energy data from $\tau$ decays(16) and from deep inelastic scattering (DIS)(17) are also shown in Fig. 5. Separate determinations, corresponding to the two fragmentation models are displayed. We note that:

1. The values of $\alpha_s$ obtained using T.C. agree well with these using the asymmetry function;

2. Both sets of data are in agreement with a decrease of $\alpha_s$, as expected from QCD, and with values of $\Lambda = 55$ MeV for Ali and 160 MeV for Lund.

The combined result from the P.T.C. and asymmetry function yields

$$\Lambda = 105 \pm 25 \pm 6 \, \text{MeV}$$

or $\alpha_s = 0.12 \pm 0.007 \pm 0.01$ at 44 GeV

and $\alpha_s = 0.124 \pm 0.004 \pm 0.012$ at 35 GeV.

The error is dominated by systematic errors from the fragmentation models. These results were obtained using a complete second order QCD calculation (11) together with a rigorous Monte Carlo method which integrates over the radiation effects, jet dressing, fragmentation, and detector resolution in a single step. These results should be compared with other results(18), which used approximate calculational methods(19), only after the latter are properly renormalized(5). The data are consistent with a logarithmic $\log s$ dependence of the coupling constant $\alpha_s$, as expected from QCD and the low energy data from the $\tau$ decays and from the deep inelastic scattering data.

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   These studies conclude that the approximate formulae commonly used (19) will yield $\alpha_s$ values $= 25\%$ too high for the $E_\text{cut}$ and 5 to 10 $\%$ too high for $M_\text{cut}$, when the soft partons are discarded. We emphasize that the calculation used here and in our previous work do not involve this approximation.


FIGURE CAPTIONS

Fig.1 The measured values of second order $\alpha_s$ from PEP and PETRA versus time for a. the LUND and b. the ALI models from references 5,7 and 18.

Fig.2 Infrared cut dependence of several commonly used experimental measurables: fraction of low thrust events as a function of the invariant mass $M$, integrated energy asymmetry and planar triple energy correlation as functions of the parton energy fraction $E$. The curves, corresponding to the right-handed scales, are Monte Carlo results of the partons while the points, corresponding to left-handed scales, are the results after fragmentations and detector simulation.

Fig.3 P.T.C. data as a function of $X_2$ for several values of $X_1$ are shown as solid points, at 35 and 44 GeV. The histograms are full Monte Carlo simulations including radiative effects, detector simulation and fragmentation based on the Lund (solid) and the Ali (dashed) hadronization models with $\Lambda = 100$ MeV.

Fig.4 The integrated planar T.C., $\langle I(\phi=40^\circ)\rangle$, is shown as solid points as a
function of cm energy. The calculated $q\bar{q}$ contribution is shown as the dashed curve. The QCD curves are for $\Lambda = 100$ MeV for both the Ali and Lund models.

Fig.5 The measured values of the strong coupling constant $\alpha_s$ as a function of cm energy. The data above 20 GeV are the MARKJ data obtained using both planar T.C. and energy asymmetry, while the data below 10 GeV are from the $\Upsilon$ decays. The deep inelastic scattering data (DIS) is also shown. The curves are the prediction of QCD with $\Lambda = 160$ and 55 MeV with $N_F = 5$ above 20 GeV and = 4 below 9 GeV. The curves are smoothly connected between 9 and 20 GeV.
Figure 1
Figure 2

Figure 3
Figure 4

Figure 5
DETERMINATION OF $\alpha_s$ FROM $\Upsilon(1S)$ DECAYS

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Abstract

We have measured the spectrum of direct photons from the decay of the $\Upsilon(1S)$. We compare this spectrum with several theoretical models and evaluate the strong coupling constant ($\alpha_s$) in each model.
QCD predicts that the dominant decay mode of the $\Upsilon(1S)$ is through three gluons which then hadronize (Fig. 1a). The theory also describes the decay of the $\Upsilon(1S)$ when one of the three gluons is replaced by a photon (Fig. 1b). The similarity of these processes can be exploited to normalize the strong coupling constant ($a_s$) to the known electromagnetic coupling constant ($\alpha_{em}$). The relative rate of these decay processes is given by

$$B_\gamma = \frac{\Gamma(\Upsilon^+\gamma gg)}{\Gamma(\Upsilon^+ggg)} = \frac{36 \alpha_{em} a_s}{5} \cdot \frac{1}{2} \left[ 1 + (2.2 \pm 0.6) a_s / \pi \right]$$

(1)

where $q$ is the charge of the $b$ quark ($-1/3\ e$). We can relate $a_s$ to the fundamental energy scale of QCD ($\Lambda_{QCD}$) by

$$a_s = \frac{4\pi}{\beta_0 \ln(x)} - \frac{1}{\beta_0^2 \ln^2(x)}$$

(2)

where $\beta_0 = 11 - 2n_f / 3$, $\beta_1 = 102 - 38n_f / 3$, $x = Q^2 / \Lambda_{QCD}^2$, $n_f$ is the number of active quark flavors (4 for the $\Upsilon(1S)$), and $Q = 0.157\ \Lambda_{QCD}$.

The CLEO detector has been described in detail elsewhere. Here we will briefly mention those detector elements used in finding photons. Our main photon detector is a twelve radiation length lead/PWC sandwiched shower counter. This device covers about 47% of $4\pi$ solid angle and has an energy resolution $\sigma_E / E = 0.21 / \sqrt{E}$ (E in GeV) and an angular resolution of $10^\prime\ mr$. Our central tracking chamber is used to veto those showers which are associated with charged particles. This device consists of a 17 layer cylindrical drift chamber inside a 10 kG solenoidal magnetic field and has a momentum resolution $\sigma_p / p = 0.007p$ (p in GeV/c). Photon candidates are vetoed if they are matched to a charged particle track projected from the central drift chamber. To reduce the background from $\tau$'s produced in the decays of $\tau$'s and $\eta$'s, we reject any candidate which lies within 320 $\text{mr}$ of another candidate.

The data used in this analysis consists of 12 pb$^{-1}$ on the $\Upsilon(1S)$ resonance and 4.1 pb$^{-1}$ taken on the continuum just below the $\Upsilon(1S)$ resonance. This sample contains 223000 observed hadronic decays of the $\Upsilon(1S)$. Our standard hadronic event selection criteria were applied. From Monte
Carlo simulation we find that our event selection criteria accept about 90% of the hadronic decays of the $\Upsilon(1S)$ (80% for continuum decays).

The signal of direct photons suffers from a number of backgrounds. The first is due to non-resonance hadron production. We use our sample of below-resonance continuum data, appropriately scaled for relative luminosities and the energy dependence of the continuum cross section to evaluate this background. The second is due to the decay of the $\Upsilon(1S)$ through a virtual photon into $q \bar{q}$ (Fig. 1c). We estimate that this process accounts for about 11% of our $\Upsilon(1S)$ decays\textsuperscript{4}. This process is similar to that for non-resonant hadron production, except that it has no contribution from initial state radiation of photons. We have used an EGS Monte Carlo event simulation\textsuperscript{5} to correct our continuum data for initial state radiation and then used this corrected continuum data to subtract off the $\Upsilon(1S)+q \bar{q}$ contribution to the resonance data.

The final background comes from the decays of $\pi^*$ and $\eta$ mesons into photons, where the photon pairs which are produced appear to merge into a single shower in our detector or where one of the two photons is missed by our detector. This is the major source of background at low momenta and limits our measurement to values of $x_\gamma > 0.55$ ($x_\gamma = E_\gamma / E_{beam}$). Most of this background is attributable to photons from $\pi^*$'s; photons from $\eta$'s amount to only about 10% of the background in the region $x_\gamma > 0.55$. Given the momentum spectrum for $\pi^*$'s, we can determine the spectrum for the photons they produce, using a Monte Carlo simulation of our detector to account for the probability that they merge into a single shower or that they are lost. Using our measured $\pi^*$ momentum spectrum\textsuperscript{6} we have found that the merging probability is negligible for $x_\pi < 0.2$ ($x_\pi = p_\pi / E_{beam}$) and rises linearly to 70% at $x_\pi = 1.0$. We have checked this by assuming the $\pi^*$ spectrum to be 1/2 our measured spectrum\textsuperscript{4} for $\pi^*$ and have obtained consistent results. Using these $\pi^*$ spectra and a momentum dependent $\pi^*$ angular distribution from a Monte Carlo simulation of $\Upsilon(1S)+ggg$ events, we calculated the expected background level and correct our resonance data. We correct for photons from $\eta$'s by assuming $\eta$ production to be 0.32 times that for $\pi^*$'s\textsuperscript{7}.

Our observed signal and these backgrounds are shown in Fig. 2. Also shown is the final signal after correcting the observed signal for the various backgrounds. At low values of $x_\gamma$, the background level due to photons from $\pi^*$ and $\eta$ decays is comparable to the level of the corrected spectrum. At high values of $x_\gamma$, the background levels are negligible. We observe photons apparently having $x_\gamma > 1.0$ due to the energy resolution of the shower counter.

To determine a value for $a_s$, we must correct our data for the contribution below $x_\gamma = 0.55$. We have fit three theoretical spectra to our data in order to estimate this. The first is derived from a lowest order perturbative calculation\textsuperscript{8} which gives a result like the photon spectrum expected from the
The decay of orthopositronium into three photons. The second is from a calculation by Photiadis, which corrects the lowest order spectrum near \(x_\gamma = 1\) by summing the leading log contributions to all orders in perturbation theory. The third is obtained from a cluster model Monte Carlo by Field, which attempts to correct the lowest order spectrum for effects due to the hadronization process.

We modify these theoretical spectra to reflect the resolution and efficiency of the CLEO detector and fit them to our background-subtracted data by varying only the overall normalization. These spectra are shown with the data in Fig. 3. For the lowest order QCD spectrum, the Photiadis spectrum, and the Field spectrum, we obtain \(\chi^2\)'s of 14.2, 10.8, and 8.1, respectively, for 11 degrees of freedom. The resolution of our detector does not allow us to make a definitive choice among these models, though we tend to favor the Field model prediction. The detection efficiency used in smearing the theoretical spectra was determined from a Monte Carlo simulation of the events and the detector. For most of the available range in \(x_\gamma\), this efficiency is constant at 33%. Above \(x_\gamma = 0.9\), the efficiency falls sharply; however, this has little effect on our results due to the small fraction of the photon spectrum which it affects.

Table I shows the values of \(B_\gamma\), \(a_\gamma\), and \(A_{\pi S}\) obtained from these spectra. The first error is statistical. The second is systematic and includes the uncertainties in our determination of the \(\pi^*\) spectrum (4.1%), the amount of
Table I. Values of $B_\gamma$, $a_s$, and $A_{\eta\gamma}$ for three assumptions about the shape of the photon energy spectrum below $x_\gamma^T=0.55$.

<table>
<thead>
<tr>
<th>Model:</th>
<th>QCD</th>
<th>Photiadis</th>
<th>Field</th>
</tr>
</thead>
<tbody>
<tr>
<td>$B_\gamma$ (%)</td>
<td>1.88$\pm$0.14$\pm$0.17</td>
<td>2.03$\pm$0.15$\pm$0.16</td>
<td>2.54$\pm$0.18$\pm$0.14</td>
</tr>
<tr>
<td>$a_s$</td>
<td>0.40$\pm$0.04$\pm$0.06</td>
<td>0.36$\pm$0.04$\pm$0.04</td>
<td>0.27$\pm$0.03$\pm$0.03</td>
</tr>
<tr>
<td>$A_{\eta\gamma}$ (GeV)</td>
<td>0.37$\pm$0.05$\pm$0.07</td>
<td>0.32$\pm$0.05$\pm$0.06</td>
<td>0.19$\pm$0.04$\pm$0.04</td>
</tr>
</tbody>
</table>

$\Gamma(1S)^{-}\bar{q}q$ present (1.7%), and the ratio of $\eta$ to $\pi^+$ production (1.3%). We have also included the uncertainty in our determination of the efficiency of our hadronic event selection criteria. This uncertainty is dependent on the shape of the momentum spectrum. We estimate this uncertainty to be 7.7%, 6.2%, and 3.3% for the spectra from lowest order QCD, Photiadis, and Field, respectively. The systematic errors for $a_s$ and $A_{\eta\gamma}$ reflect the theoretical uncertainty expressed in Eq. 1 (5.5%).

Other experiments report a wide range of values for $A_{\eta\gamma}$ with typical values in the region of 0 to several hundred MeV. These other methods, however, are often subject to large uncertainties due to non-perturbative effects. In our case the major uncertainty is due to the correction which we must make as a result of the limited range of our measured spectrum. In principle this difficulty can be resolved experimentally.

In conclusion we have measured $B_\gamma$, $a_s$, and $A_{\eta\gamma}$ using three theoretical models to extrapolate below the region over which we have useful data ($0.55<x_\gamma^T<1.0$). Our data tend to favor a softer spectrum than that predicted by lowest order QCD though it is not possible to make a definitive selection among these models due to the level of the uncertainties in our data.

References

6] S. Behrends et al., op. cit. The $\tau^+$ measurement we report here is based on an improved data sample relative to that given in this reference.
7] We have measured $\eta_{/\tau^+}$ on the continuum for $x>0.6$ to be 0.32$\pm$0.10$\pm$0.05.
QCD-EFFECTS IN JET PROFILES IN $\mu$-p SCATTERING

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ABSTRACT

QCD-effects in jets produced in $\mu$-p deep inelastic scattering (DIS) have been investigated. Data is found to follow the QCD-predictions including contributions from soft gluon emission. Evidence is obtained that data favour a string-alike fragmentation over an independent jet model.
1. INTRODUCTION

Besides e⁺e⁻ collisions, deep inelastic μ-p scattering (DIS) is an appropriate process to study QCD-modifications of the simple Quark-Parton-Model predictions. The interaction of the virtual photon with the constituents of the target nucleon results in a forward quark jet and a backward target jet. The direction of the virtual photon, the momentum transfer $Q^2$ and the energy transfer $\nu = k_\mu - k_\mu'$, can be well determined from the kinematics of the muon. The total hadronic energy $W$ is related to $Q^2$ and $x_{Bj}$ by

$$W^2 = M^2 + Q^2 \left( \frac{1}{x} - 1 \right)$$

and ranges from about 4 GeV to 24 GeV in this experiment.

QCD effects can be studied using single particle distributions (e.g. $z$, $x_F$, $p_t$, particle and energy-flows) as well as event shape variables (like Thrust, Oblateness, Spurrity). For example, the $Q^2$-dependence of the differential multiplicity (1) of charged hadrons and the increase of $<p_t>$ relative to the virtual photon direction (2) with $W$ are effects predicted by QCD.

The results presented are based on an analysis of 30000 μ-p events from a hydrogen target using the 280 GeV μ-beam and the EMC double magnet spectrometer. The apparatus (3) consists of a forward spectrometer to measure the scattered muon and fast hadrons and a vertex spectrometer part including a streamer chamber. Tracks are measured in a momentum range between .2 and 280 GeV/c, this corresponds to a nearly 4π acceptance in the CMS.

In μ-p DIS, the following diagrams contribute to additional $p_t$ for forward going hadrons.
These QCD-effects have to be isolated from other sources of $p_t$, such as transverse motion of the quarks inside the nucleon (primordial $k_t$), the fragmentation itself and the decay of unstable particles.

Contributions from fragmentation and from $k_t$ do not explicitly depend on $W$ and are expected to be roughly forward/backward symmetric, whereas hard QCD-effects as well as effects from the emission of soft gluons (as in the Lund 4.3-model) increase with $W$ and occur only in the forward jet. (4)

2. JET PROFILES IN $\mu$-p SCATTERING

A possibility to separate QCD-effects in jets from purely kinematic $W$-effects is the analysis of jet-profiles which are defined as:

$$\frac{d\varepsilon}{d\lambda} = \frac{1}{N_{ev}} \sum \frac{\Delta\varepsilon^i}{\Delta\lambda} \quad \varepsilon^i = \frac{E_{i \text{charged}}}{E_{\text{jet}}}$$

$$E_{\text{jet}} = \sum E_{i \text{charged}}$$

(summation: $x_F < 0$ backward jet
$x_F > 0$ forward jet)

$$\lambda = \frac{x_F}{p_t} = \frac{\cot\theta}{W/2}$$

Scaling of jet profiles (e.g. being independent of jet energy) was found to be valid at low energy $pp$-reactions (5) and in neutrino-p scattering (6).

Using charged particles only, EMC data has been analysed in 4 $W$-intervals and is compared to various model predictions.

$W_1$: 4 - 8 GeV  $W_2$: 8 - 12 GeV  $W_3$: 12 - 16 GeV  $W_4$: 16 - 20 GeV
As is shown in Fig. 3, the independent jet model without QCD shows very good scaling, except for the first bin. The ratios of jet-profiles at different energies, which are defined as:

\[ W_{ik}(\lambda) = \frac{d\varepsilon/d\lambda (W_i)}{d\varepsilon/d\lambda (W_k)} \]

are very close to unity both for the backward and the forward jet (Fig. 5a). However, data jet profiles do not show this scaling behaviour in the forward jet (Fig. 4).

Ratios of energy-profiles as expected from the Lund model are compared to data in Fig. 5b. While the backward jet ratios essentially remain flat (also with QCD included), the forward jet shows sizeable deviations increasing with \( W \). They are due to hard QCD as well as to multiple emission of soft gluons, both effects being almost equally large.

In contrast to the independent jet model, there are scaling deviations present already in the Lund-model without QCD (Fig. 5b, dotted line). In fact this behaviour is expected as a "string-effect" due to a "crosstalk" between the backward and the forward jet, it depends (for the forward jet) on the transverse mass of the backward going particles (baryons from the target remnant) and contributes to the scaling deviations in the string model (7). By definition, there is no "crosstalk" in the independent jet fragmentation. Good agreement of data and Monte-Carlo can only be obtained using the Lund-fragmentation with full QCD.
3. THE RADIATION PROFILE FUNCTION

A radiation profile function parametrizing the jet profile has been suggested (5) and was tested at low energies and in Monte Carlo calculations for $e^+e^-$ annihilation (8).

$$\rho(\lambda) = \frac{M}{1 + M^2 \lambda^2}^{3/2}$$

The parameter M corresponds to the average transverse mass in the central rapidity plateau. To test the validity of $\rho(\lambda)$ in $\mu$-p scattering a LSQ fit of M to the data (see Fig. 4) and various Monte Carlo models is summarized in Fig. 6. The resulting M reflects the "broadness" of the jets with fragmentation effects (limited $p_t$) being eliminated. In general, $\rho$ is found to describe the jet profile for the forward jet satisfactorily, the backward jet, however shows sizeable deviations. For the Lund-model with QCD and the "string-effect" included, the increase of M with W roughly agrees with the data and is much stronger than in the independent jet model.

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Fig. 4: Forward jet profiles for different W-intervals (data and fit)

Fig. 5a
Ratios of jet profiles for IJ model without QCD (5a) and the Lund-model with QCD compared to data (5b)

Fig. 6: Fitted M-values for data and Lund-model
MEASUREMENT OF ASSOCIATED PRODUCTION
OF W AND Z WITH JETS IN UA2

The UA2 Collaboration

Bern  –  CERN  –  Copenhagen (NBI)  –  Heidelberg  –  Orsay(LAL)  –
Pavia  –  Perugia  –  Pisa  –  Saclay (CEN) Collaboration

presented by Valerio Vercesi

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ABSTRACT

Following the 1985 Collider run at \( \sqrt{s} = 630 \) GeV, the UA2 Collaboration has collected data for W and Z production and subsequent leptonic decay, corresponding to an integrated luminosity of 863 nb\(^{-1}\). A preliminary analysis of the production of W and Z bosons, which are accompanied by jets of high transverse momentum, is presented for the full data sample. The data are in good agreement with previously published results and with the expectations of QCD models.
1. Introduction

In previous publications\textsuperscript{1)\textemdash}2) the UA2 Collaboration reported experimental results on the processes
\[
\begin{align*}
\bar{p} + p & \rightarrow W^\pm \rightarrow e^\pm + \nu(\bar{\nu}) \\
\bar{p} + p & \rightarrow Z^0 \rightarrow e^+ + e^-
\end{align*}
\]
where \( W \) and \( Z \) are the Intermediate Vector Bosons (IVB). In the QCD framework of \( W \) and \( Z \) production a significant fraction of IVB's are produced at high \( p_T \) and are accompanied by hadronic jets. In this presentation we report on the associated production of \( W \) and \( Z \) with jets of large transverse momentum, and we include in this analysis a preliminary data sample collected in the period September\textendash December 1985. Following this last successful operational period of the CERN \( \bar{p}p \) Collider, the total luminosity collected by the UA2 experiment is 863 nb\textsuperscript{\textminus 1}.

The UA2 detector has been described elsewhere\textsuperscript{3)} and therefore its features will not be repeated here. A detailed explanation of the selection criteria used to identify electrons in UA2 is given in ref. 1 and 2.

2. Electrons with associated jets

The main consequences of QCD contributions to \( W \) production are to increase the production cross section by \( \sim 30\% \) and to give the \( W \) a sizeable average transverse momentum \( < p_T^W > \). Moreover, for relatively large values of \( p_T^W \), the \( W \) bosons are expected to recoil against hadronic jets which can be observed experimentally.

To study the production of hadronic jets in association with high \( p_T \) IVB's in an unbiased way, we select a sample of events with large \( p_T^{e^+} \) and large \( p_T^{e^-} \) in the case of the \( W \), without any restriction on the associated event topology. For the \( Z \) boson we evaluate the invariant mass of the \( e^+ e^- \) system, and a selection about the \( Z \) mass provides an essentially background free sample.

In the case of the \( W \), we select events on the basis of the missing transverse momentum, defined as
\[
\vec{p}_T^{\text{miss}} = \vec{p}_T^{\nu} = -\vec{p}_T^{e^+} - \Sigma \vec{p}_T^{j} - \lambda \vec{p}_T^{SP}
\]
The sum runs over all the hadronic jets $j$ with $p_T^j > 5$ GeV and $\sum p_T^j$ represents the total transverse momentum carried by particles not belonging to jets. The correction factor $\lambda$ accounts for the non-linearity of the calorimeter and the limited acceptance of the UA2 detector. We obtain a value of $\lambda = 1.5 \pm 0.6$ by minimising $< p_T^{\text{miss}} >$ for the sample of $Z^0$ events observed in the experiment. The distribution of $p_T^\nu$ for the 1985 data sample only is shown in fig. 1 for $p_T^e > 17$ GeV, together with an estimate of the background (see ref. 2 for details of the background calculation). The Jacobian peak at $p_T^\nu = 40$ GeV/c, as expected from $W \to e\nu$ decay, is clearly evident in fig. 1. For $p_T^\nu > 25$ GeV/c we obtain 110 events in the 1985 data sample, with $11.6 \pm 2.7$ events due to background contributions.

We now consider the total data sample satisfying $p_T^e > 15$ GeV/c (17 GeV/c for the 1985 data) and $p_T^\nu > 25$ GeV/c and we compare this sample with QCD expectations. After applying a cut on the transverse mass of the $e\nu$ pair, $M_T > 20$ GeV/c, the sample contains 232 events (the background estimate is $22.2 \pm 3.3$ events). Fig. 2 shows the distribution of the $W$ transverse momentum $p_T^W$, defined as $p_T^W = p_T^e + p_T^\nu$, without background subtraction. The shaded area identifies events in which the $W$ is produced in association with at least one jet of $E_T > 5$ GeV and, as expected, the majority of high-$p_T W$'s have associated jet activity. A QCD prediction of Altarelli et al. is in good agreement with the data. For $p_T^W > 30$ GeV/c we measure 7 events (including 1.3 background events), to be compared with 5 events predicted by QCD calculations.

Fig. 3 shows the transverse momentum distribution for the $Z$ events in the total corresponding sample of 30 events, together with a QCD prediction. The average value is $< p_T^Z > = 7.5 \pm 1.6$ GeV/c, comparable with the value of $< p_T^W > = 8.8 \pm 1.7$ GeV/c.

3. The transverse energy of the hadronic system

When more than one jet is present in the final state the study of the scalar $E_T$ of the hadronic system brings additional information on the $W$ production mechanism. Following the analysis of ref. 2 we compare the probability $P_W(J)$ for a $W$ to be produced in association with a system of jets $J$ with the corresponding $P_{j_1j_2}(J)$ of QCD events containing 2 jets $j_1$ and $j_2$ and an additional system $J$. Because of the larger number of relevant subprocesses we would naively expect larger values for $P_{j_1j_2}(J)$. We
retain only those \( (j_1,j_2) \) pairs that satisfy the same kinematic cuts as \( (e,\nu) \) pairs from W decay and with an invariant mass compatible with that of the W. In fig. 5 we compare \( F_{j_1,j_2}(E_T^0) \), defined as the probability that a pair \( j_1,j_2 \) is associated with a system \( J \) having \( E_T^J > E_T^0 \), with the corresponding function \( F_W(E_T^0) \) of W→eν events.

We also compare this data sample with the predictions of a Monte Carlo program\(^5\) which generates W bosons with jets according to perturbative QCD up to order \( \alpha_s \), and which includes a full simulation of the UA2 detector. The two curves superimposed on fig. 5 are for \( \alpha_s \) values of 0.14 and 0.20. Both the measurements of \( F_{j_1,j_2} \) and the QCD expectations of \( F_W \) are in good agreement with the full sample of data collected between 1983 and 1985. For \( E_T^0 = 30 \) GeV the QCD model predicts \( 5.0 \pm 0.6 \) events, to be compared with 4 observed events.

We have also studied differences between the underlying event structure obtained excluding the electron from the W decay, and the event structure of non-diffractive pp interactions ("minimum bias" events). To remove expected contributions from QCD hard processes we consider only those events having no associated jet. In fig. 4 we compare the quantity \( E_T \), defined as the total transverse energy deposited in the calorimeter of particles of the underlying event, in the pseudo-rapidity range \( |\eta| < 1 \). We obtain

\[
\langle E_T \rangle = 8.5 \pm 0.4 \quad \text{(W→eν events with no jet)}
\]

\[
\langle E_T \rangle = 6.5 \pm 0.3 \quad \text{(minimum bias events with no jet)}
\]

We note, however, that these measurements are sensitive to the \( E_T \) threshold used to define the jets and that 2.3% of minimum bias events have at least one cluster of \( E_T > 5 \) GeV in the pseudo-rapidity range \( |\eta| < 1 \).

REFERENCES


W AND Z PRODUCTION PROPERTIES

UA1 Collaboration

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ABSTRACT

The production properties of the Intermediate Vector Bosons produced at the CERN SPS collider are presented. The longitudinal and transverse momentum distributions and the jet activity associated to the W and Z events are well described by the QCD improved Drell-Yan mechanism of quark-antiquark annihilation.
I - Event samples and cross sections

Between 1982 and 1985 the UA1 experiment \[1\] has accumulated data corresponding to an integrated luminosity of 729 nb\(^{-1}\), where 136 nb\(^{-1}\) were taken at the energy of \(\sqrt{s} = 546\) GeV \[2,3,4,5\] and 593 nb\(^{-1}\) at \(\sqrt{s} = 630\) GeV. The number of Intermediate Vector Bosons detected until now is 82 \(W \to \mu \nu\), 20 \(Z \to \mu^+\mu^-\) events and \(\sim 300\) \(W\)'s and \(\sim 40\) \(Z\)'s in the electron channel. These numbers are not yet final as the analysis of data obtained in 1985 is still in progress. The corresponding \(W\) and \(Z\) production cross sections for the two energies are shown in table 1 together with the theoretical predictions of Altarelli et al. \[8\]. The results for the muon and electron channels are in good agreement, and both are compatible with the theoretical expectations. In particular, the data show a 15% \(\pm\) 18% increase of the \(W\) production cross section from \(\sqrt{s} = 546\) GeV to \(\sqrt{s} = 630\) GeV, which, within the error, agrees with the theoretical prediction of 24%.

<table>
<thead>
<tr>
<th>measured cross-sections</th>
<th>(\sqrt{s} = 546) GeV</th>
<th>(\sqrt{s} = 630) GeV</th>
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</thead>
<tbody>
<tr>
<td>(\sigma_B(W^\pm \to \mu^\pm \nu))</td>
<td>0.56 (\pm 0.18 \pm 0.12) nb</td>
<td>0.66 (\pm 0.12 \pm 0.14) nb</td>
</tr>
<tr>
<td>(\sigma_B(W^\pm \to e^\pm \nu))</td>
<td>0.55 (\pm 0.08 \pm 0.09) nb</td>
<td>0.63 (\pm 0.05 \pm 0.09) nb</td>
</tr>
<tr>
<td>(\sigma_B(Z^0 \to \mu^+\mu^-))</td>
<td>100 (\pm 50 \pm 15) pb</td>
<td>70 (\pm 29 \pm 13) pb</td>
</tr>
<tr>
<td>(\sigma_B(Z^0 \to e^+e^-))</td>
<td>42 (\pm 25 \pm 18) pb</td>
<td>74 (\pm 23 \pm 11) pb</td>
</tr>
</tbody>
</table>

Table 1

<table>
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<tr>
<th>theoretical predictions</th>
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<tr>
<td>(\sigma_B(W))</td>
</tr>
<tr>
<td>(\sigma_B(Z))</td>
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</tbody>
</table>

From now on, we discuss the \(W\) and \(Z\) data in the electron channel only. The muon channel has been presented in this conference by Steve Wimpenny \[14\]. We emphasise that the final selection and analysis of the data from the 1985 run (representing \(\sim 45\%\) of the total) are not yet definitive. The selection of \(W \to e\nu\) is by now rather standard: we require events with a large (and validated) missing transverse energy (\(E_T^{\text{miss}} > 15\) GeV) for the neutrino and an isolated electromagnetic cluster (\(E_T > 15\) GeV) with an associated track for the electron. We have at present a sample of 262 \(W \to e\nu\) events. In the selection of \(Z \to e^+e^-\) events, for which we require two electromagnetic clusters, we are left with 32 \(Z \to e^+e^-\) candidates. The first cluster has to be isolated with an associated track and cluster \(E_T > 15\) GeV, while for the second cluster we require \(E_T > 8\) GeV.
II - Mass distributions

For this data sample we show the transverse mass distribution for the $W \to e\nu$ events in fig. 1 and the mass distribution for the $Z \to e^+e^-$ sample in fig. 2. In fig. 1 only $W$'s (202 events) where neither the electron nor the neutrino is within $\pm 15^\circ$ from the vertical plane or the horizontal plane along the beams are taken into account, to avoid a region where the electron energy is not well measured. The hatched region is the estimated background contribution (10%) [7], which comes from: i) jet-jet fluctuations and ii) $W \to \tau\nu$ events, where the $\tau$ decays leptonically or hadronically. Superimposed to the data is the transverse mass distribution expected from the published mass fit of 1984 ($M_W = 83.5$ GeV/$c^2$) [7], which describes very well the present data. Notice that no event is observed beyond the $W$ mass region thus setting a lower limit (90% CL) of 230 GeV/$c^2$ for a more massive $W$, assuming same coupling and $e\nu$ branching ratio.

For the study of the $Z \to e^+e^-$ mass distribution, events are now required to have also a second electromagnetic cluster with a matching track with as stringent criteria as for the first cluster, and no missing energy ($E_T^{\text{miss}} < 15$ GeV). The remaining 30
events are shown in fig.2, and are compared with the expected distribution using the published Z mass fit of 1984 \((m_Z = 93.0 \text{ GeV/c}^2)\) [7]. No \(e^+e^-\) events have been detected above the \(Z^0\) region which sets the lower limit (90% CL) for a second \(Z^0\) at 180 GeV/c\(^2\) under the assumption of the same coupling as \(e^+e^-\) branching ratio.

III - \(W,Z\) longitudinal momentum

The Feynman \(x\) of the \(W\) is obtained from the longitudinal momenta of the neutrino and electron as

\[
x_w^- = \frac{p_L^\nu + p_L^e}{\sqrt{s}/2} = x_q^\nu - x_q^e
\]

which is, per definition, the difference of the fractional momenta of the annihilating quarks. Unfortunately, the longitudinal momentum of the neutrino cannot be directly measured by the experiment. But by imposing the \(W\) mass on the electron-neutrino system

\[
x_w^2 = (E_e + E_\nu)^2 - (p_{L\nu} + p_{Le})^2
\]

we can calculate the longitudinal momentum of the neutrino. The two solutions for the longitudinal neutrino momentum, one corresponding to the neutrino being emitted forwards in the \(W\) rest frame and the other backwards, leave us with a two-fold ambiguity for \(x_w^-\) in about one-half of the events (in the other cases one of the two solutions is unphysical \(|x_w^-| > 1\) or both give the same \(x_w^-\)). When the constraint of energy conservation in the overall interaction is considered, the smallest of the two \(x_w^-\) solutions is almost always preferred, so we choose this one for all the events.

Using energy/momentum conservation \(x_q^\nu x_q^e = x_w^2/s\) we can extract the values of \(x_q^\nu\) and \(x_q^e\), and, for the \(W\)'s with a well determined charge for the decay \(e^+\), (better than 2 standard deviations from infinite momentum), an extraction of the longitudinal momentum fractions of the \(u\) and \(d\) quarks separately is possible. Since the \(W^+\) is supposed to result from an \(u\bar{d}\) annihilation, we have \(x_u^\nu = x_q^\nu\) and \(x_d^\nu = x_q^e\), and for the \(W^-\) vice versa (\(u\bar{d}\)).
These longitudinal momentum distributions are expected to reflect the structure function of the annihilating partons. In fig. 3a, the $|x_W|$ distribution is shown. It is well described by the structure functions of Eichten et al. [9] with $\Lambda = 0.2$ GeV. One can also see in the data some indication of the expected softening of $|x_W|$ with increasing $\sqrt{s}$. The longitudinal momentum of the $Z$ is shown in fig. 3b, compared with the same prediction. Notice that in the case of $x_Z$ there is no kinematical ambiguity as for $x_W$.

![Fig. 3a](image)

**Fig. 3a**
Feynman $x$-distribution for $W$'s produced at $\sqrt{s} = 546$ GeV and $\sqrt{s} = 630$ GeV. The curves show the predictions using the structure functions of Eichten et al. [9] modified in the appropriate way to take into account selection biases, experimental resolution and the analysis bias of choosing always the lowest $x_W$ solution.

![Fig. 3b](image)

**Fig. 3b**
Feynman $x$-distribution for $Z$'s, compared to the structure functions of Eichten et al.

In fig. 4 are shown the $x$ distributions of the quarks and antiquarks sampled by $W$ production. The data are again in a satisfactory agreement with the expectation resulting from the structure functions of Eichten et al. [9]. The separate $x_u$ and $x_d$ distribution in fig. 5 and fig. 6 are also well described by the theoretical expectations [9].
Fig. 4
Feynman x-distribution for the proton and antiproton partons making the W. The curve shows the prediction of Eichten et al.

Fig. 5
Feynman x-distribution for u-quarks sampled in W production. The curve shows the prediction of Eichten et al.

Fig. 6
Feynman x-distribution for d-quarks sampled in W production. The curve shows the prediction of Eichten et al.
IV - $W$, $Z$ transverse momentum and associated jet activity

Before discussing the transverse momentum distribution we must mention briefly some points on the detection of jets in UA1. At lowest order the $W$'s and $Z$'s are produced by the Drell-Yan mechanism of quark-antiquark annihilation. In the QCD improved picture of the production mechanism coloured quarks and antiquarks can radiate gluons thus generating a recoiling transverse momentum for the $W$. So, higher-order corrections must be taken into account for the correct description of the observed $P_T^{W,Z}$ distributions, especially at large values of $P_T$.

Now, the radiated gluons, responsible for the $P_T$ of the $W$ or $Z$, may have a sufficiently high transverse energy to produce observable jets in our apparatus. These jets are reconstructed by the standard UA1 jet algorithm [10], and we consider here jets with $E_T^{\text{jet}} > 5$ GeV in the pseudorapidity range $|\eta^{\text{jet}}| < 2.5$. This jet algorithm associates calorimetric cells in a cone of radius $R$ in an azimuthal and pseudorapidity space

$$\Delta R = 1, \quad \Delta R = \sqrt{\Delta \eta^2 + \Delta \phi^2}$$

starting from an initiator cell with $E_T > 2.5$ GeV.

In fig. 7a the observed transverse momentum distribution of the $W$ is shown. The distribution has a peak at $\sim 4$ GeV/c, reflecting primarily the neutrino transverse energy measurement resolution, and a long tail well described by the QCD calculations of Altarelli et al. [8]. (The superimposed curve uses the structure functions of GHR[11] with $\Lambda = 0.4$ GeV). The average value of the experimental distribution is $< P_T^W > = (7.7 \pm 1.0)$ GeV/c (statistical error only). Events with at least one detected jet (shaded events in fig. 7a)), have a higher mean value $< P_T^W > = (12.5 \pm 2.0)$ GeV/c (statistical error only) in qualitative agreement with the QCD picture, in which one should observe jets associated with higher transverse momentum $W$'s. Notice, that all the events with $P_T^W > 20$ GeV/c have at least one reconstructed jet. During the run of 1985 we have also detected one $W$ event with an extremely high transverse momentum of 89 GeV/c. In the $W \rightarrow \mu \nu$ sample a similar event with $P_T^W = 66$ GeV/c has been observed. A graphical representation of these events is shown in fig. 8a and b. These $W$ events of highest $P_T$ have both two jets recoiling against the $W$. 
Fig. 7a
The $W$ transverse momentum distribution. The curve shows the QCD prediction of Altarelli et al. [8]. The shaded subhistogram shows the contribution from events with at least one reconstructed jet with $E_T > 5$ GeV.

Fig. 7b
The $Z$ transverse momentum distribution. The shaded area represents events with at least one jet with $E_T > 5$ GeV.

Fig. 8
A graphic display of the UA1 event containing the highest transverse momentum $W$ decaying in the $e\nu$ (fig. 8a) and $\mu\nu$ (fig. 8b) channels, observed in the 1985 data taking period.
The transverse momentum distribution of the Z events (for which the resolution in $p_T$ is better) is shown in fig. 7b, it has a similar shape, with a mean value $< p_T^Z > = (8.5 \pm 1.8) \text{ GeV/c}$ (statistical error only). Again, events accompanied by at least one jet have in average a higher $p_T^Z$ (shaded events in fig. 7b).

The jet multiplicity distributions for W and Z events are shown in Figs. 9a and b. The number of events decreases with jet multiplicity by a factor $\approx 0.3$ per additional jet, for jets of $E_T \geq 5 \text{ GeV}$. The observed jet activities in W and Z events are very similar and agree well with simple QCD expectations [12]. The inclusive transverse energy distributions of the 121 jets associated to the W production and the 20 jets for the Z's in figs. 10a and b, which primarily reflect the transverse momentum of the events, are also well described by the QCD calculation of Stirling et al. [12].

![Jet multiplicity distribution for hadronic jets produced in association with a) W's and b) Z's. The QCD curve combines the ISAJET [13] prediction for the rate of (W+1 jet) events with the assumption that the probability of producing two jets along with the W is equal to the square of the probability to produce one jet [12].](image-url)
Jet transverse momentum distribution for:

- a) jets produced in association with W's. The curve shows the QCD prediction [12].
- b) jets produced in association with Z's.

Fig. 10

The angular distribution for jets reconstructed in W events. The distribution of \( \cos \theta' \) is shown, \( \theta' \) being the angle between the jet and the average beam direction in the rest frame of the W and the jet(s). The curve shows the QCD prediction [12]. The dashed line is the angular distribution for QCD scattering.

Fig. 11
As these jets are supposed to arise mainly from initial state gluon bremsstrahlung, their angular distribution is expected to be peaked in the beam direction. QCD predicts [12] a distribution of the form:

\[ \frac{d\sigma}{d\cos\theta_{\text{jet}}} \sim (1 - |\cos\theta_{\text{jet}}|)^{-1}, \]

where \( \theta_{\text{jet}} \) is the angle between the jet and the beam direction in the (W, jet) rest frame. Again, the experimental distribution in fig. 11 is in good agreement with this prediction [12], in contrast to the angular distribution observed in jet-jet systems at comparable C.M. energy (dashed line).

\section*{V - Conclusions}

i) The total W and Z production cross sections observed in the muon and electron decay modes are compatible and in fair agreement with theory.

ii) The observed longitudinal motion of the W is consistent with the parton model expectation for annihilating quarks and antiquarks (u\overline{d}, \overline{u}d).

iii) The measured transverse momentum distribution of the W/Z is in good agreement with QCD expectations even out to very large values of \( P_T \); events with \( P_T = M_W \) are observed.

vi) The jet activity associated to the W and Z events is well explained by initial state gluon bremsstrahlung from the incoming partons. There is no difference between the jet activities of W and Z events.

In conclusion, at the present level of sensitivity, W and Z production is well described by the QCD improved Drell-Yan mechanism of quark-antiquark annihilation.
References


ABSTRACT: Recent results in the study of high $P_T$ direct photons from the experiments NA24, R808, UA2 and WA70 are reviewed. Studies of correlation of particles associated with high $P_T$ direct photons and beam dependance of direct photons production are presented. Ratios of $\pi$ to $\pi^*$ and $\phi$ to jet productions as well as cross-sections are also presented.
I. INTRODUCTION

Since the discovery of prompt photon (often called direct photon) production at the ISR there has been a growing interest in studying this process.

The study of scattering processes with large transverse momentum ($P_T$) is a convenient way to extract information on hadron constituents (quarks and gluons) and their interactions\(^1\). The advantages of selecting high $P_T$ direct photon final states in the hard scattering of hadron constituents can be briefly recalled:

- the number of parton processes involved is small (2 at first order)
- the coupling of a photon to a quark is well understood
- the photon directly participates in the collision and can be detected whereas quarks and gluons fragment into hadrons
- no fragmentation effects are present making the results directly comparable to the theory with no need of a fragmentation model.
- the kinematics (angle and energy) of the outgoing parton, photon in our case, is measured with a good precision.

At first order in $\alpha_s$, direct photons are produced through two processes: the inverse QCD Compton process: $qg \rightarrow \gamma g$ and the annihilation: $qq \rightarrow \gamma g$. One feature of direct photon production is that a gluon is present in both processes, either in the initial state (Compton) or in the final state (Annihilation).

If the contributions from these two first order graphs can be separated, prompt photon production may allow a study of the gluon distribution within hadrons and jet fragmentation.

Prompt high $P_T$ photons may also be produced by many bremsstrahlung diagrams like the following:

These contributions as well as the next-to-leading order in $\alpha_s$ contributions have been calculated\(^2,3\). All together these contributions increase the prompt photon cross-section by a factor of the order of 1.5 to 2, depending on the kinematic configuration.
Experimentally the detection of direct photons is not easy for several reasons. The yield of single photon relative to jets is smaller by several orders of magnitude ($\approx 10^{-4}$) because the photon production is reduced by a factor $\alpha/\alpha_s$ compared to jets and many more graphs and colour factors contribute to the jet yield. There is a large background resulting from $\pi^\circ$ and $\eta$ decays into photons, which implies a difficult measurement of the direct photon signal. This background follows from two effects. The first effect is that one of the photons from a $\pi^\circ, \eta \rightarrow \pi^\circ \pi^\circ$ decay may not be detected, either because it is outside the detector acceptance or because it has a too low an energy; the other photon then appears as a direct photon. This is the main source of background at low momentum. At high momentum, the two photons from a $\pi^\circ, \eta$ decay (or at very high $P_T$ several neutral particles belonging to the same jet), have too small a spatial separation to be resolved in the detector and appear as a single photon. Beside this background in nearly all direct photon experiments, beam halo or $\mu$ bremsstrahlung (in fixed target experiments) fakes direct photons.

However because prompt $\pi^\circ$ contrarily to $\pi^\circ$ and eta do not issue from a parton fragmentation, we expect, that at fixed $\sqrt{s}$ the ratio $\pi^\circ/\pi^\circ$ should increase with $P_T$. Then working at high $P_T$ makes the measurement easier (although in this region the statistics are poor). Several other effects contribute to the increase of the $\pi^\circ/\pi^\circ$ ratio with $P_T$: the fragmentation function becomes softer as the parton process becomes harder, the gluon distribution within hadrons becomes smoother, reducing the contribution of graphs like $gg\rightarrow gg$ to the jet ($\pi^\circ$) yield and the strong coupling constant $\alpha_s$ decreases with the hardness of the collision.

Two different techniques have been employed to extract the direct photon signal and suppress the huge background resulting from non-prompt photon sources (e.g. $\pi^\circ, \eta$ decays, non-resolved neutral multi-particle states) and hadrons misidentified as $\pi^\circ$ (e.g. $K^*_{1L}$, neutrons ...).

In the first one, called the "Direct Method" a prompt photon is defined as a photon not coming from an identified decay (e.g. $\pi^\circ, \eta$ ... decays, non-resolved neutral multi-particle states). A large fraction of $\pi^\circ$ and $\eta$ are recognised by reconstructing the two photon mass. From a knowledge of the detector acceptance and the observed rate of $\pi^\circ$ and $\eta$, most of the background can be evaluated. The direct photon signal is then calculated from the excess of observed photons over the calculated background. Detectors should be designed to have good efficiency in resolving $\pi^\circ$'s or $\eta$'s and good acceptance in order to minimize the number of lost decay photons. This method is limited at high $P_T$, where the two photons separation is comparable to the detector granularity, most of the decays cannot be resolved.

In the second technique called the "Conversion Method", $\pi^\circ$'s and $\eta$'s are not reconstructed and a direct photon candidate is defined as a neutral energy deposition in the calorimeters. The fraction
of direct photons is determined statistically by measuring the conversion rate in a preshower detector. A signal, in this detector more likely occurs for two coalescing photons than for a single one. This technique requires a precise knowledge of the conversion probability in the preshower detector and of the energy loss in the converter. The later requirement is needed when comparing the steeply falling spectra of converted and non-converted photons.

In the following, I shall review recent prompt photon production results exclusively from four experiments. NA24 and WA70 are two SPS fixed target experiments using 300 GeV/c and 280 GeV/c beams ($\pi^-, \pi^+, P$) respectively interacting on a hydrogen target. R808 is an experiment studying $p\bar{p}$ and $\bar{p}p$ collisions at a center of mass energy $\sqrt{s} = 53$ GeV at the ISR using the A.F.S. detector. UA2 operates at the SPPS collider studying $p\bar{p}$ collisions at $\sqrt{s} = 630$ GeV.

These four experiments measure prompt $\gamma$ production in the complementary kinematic regions given in Table 1.

Before presenting the results, I shall briefly describe the experimental features relevant for the analysis presented here.

**II. EXPERIMENTAL SET-UP**

Three among the four experiments mentioned above have observed direct photon production by employing the direct method (NA24, R808, WA70), while UA2 uses the conversion method.

The three experiments using the direct method require a fine granularity photon detector in order to separate the two $\gamma$'s from $\pi^+$ or $\eta$ decays (especially for fixed target experiments).

The NA24 photon detector is 9.6 radiation lengths ($X_0$) thick and consists of alternate layers of lead sheets and proportional tubes ($0^\circ, 90^\circ$ and $45^\circ$ inclination). The photon detector is followed by the NA5 calorimeter which consists of an electromagnetic ($16X_0$) and a hadronic ($6X_0$) part. An iron wall and a veto counter array located upstream of the detector safeguard the experiment against upstream interactions and muon background.

The WA70 photon detector consists of an electromagnetic calorimeter made of lead sheets interleaved with teflon tubes filled with liquid-scintillator. This detector is $24X_0$ deep and segmented longitudinally in 3 parts ($8X_0$ each). Tubes are arranged orthogonally and an electronic time of flight system is used to resolve spatial ambiguities and reject backgrounds.

The R808 photon detector is provided by two opposite walls of NAI blocks, 5.3 $X_0$ deep. This detector is located inside an uranium calorimeter consisting of a 6 $X_0$ electromagnetic part and a 3.6 absorption length hadronic part.

Table 1 shows, the energy resolution and the $\pi^+$ mass resolution achieved by these experiments.
In the UA2\textsuperscript{12}) experiment the conversion method is used to measure the direct photon signal. Two measurements have been performed: in the central region at a mean pseudorapidity $\eta=0$ and in the Forward/Backward regions : $1.1 \leq |\eta| \leq 1.7$. These measurements rely on preshower detectors which consist of a $1.5 X_0$ converter followed by a multi-wire proportional chamber in the central part and a $1.4 X_0$ converter followed by a multi-tube proportional chamber in the Forward/Backward regions. The chamber thresholds are set to 2 minimum ionizing particles deposit (mips) and 6 mips for the central and the Forward/Backward detectors respectively. In both cases, the photon energy is measured using calorimeters located behind the preshower detectors.

### III. DATA SELECTION AND ANALYSIS

For NA24, R808 and WA70 experiments details of the analysis can be found in previous publications\textsuperscript{13)}, and they are only briefly recalled here. These experiments using the direct method have similar criteria to select prompt photons. In each event, electromagnetic showers, are paired; if the invariant mass of the pair is found to be consistent within errors with the $\pi^0(\eta)$ mass, the pair is assumed to originate from a $\pi^0(\eta)$ decay. An electromagnetic shower which could not be paired is considered to be a direct photon candidate. The background to direct photon candidates is estimated from Monte-Carlo generated $\pi^0$ and $\eta$ events. Depending on $P_T$ it is due to: i) photons from $\pi^0$ and $\eta$ decays for which one of the photons escapes detection (limited detector acceptance) or is not reconstructed due to its low energy and ii) coalesced showers from unresolved $\pi^0$ decays. The observed $\pi^0$ and $\eta$ samples must finally be corrected by their relative detection efficiency to give the final observed $\pi^0/\eta$ ratio.

#### NA24 EXPERIMENT

Events are selected off-line by requiring that the direction of the triggering shower (as determined by the shower development measurement in the photon detector) points to the target, the total energy measured be consistent with the beam energy and that the calorimeter timing (using Flash ADCs) agrees with the time at which the interaction took place. These cuts remove most of the pile up and muon background.

#### WA70 EXPERIMENT

Cuts are used to remove pile up and muon background. The shower direction as given by the calorimeters must point to the target and the time-of-flight information must agree with the time of the incident particle.

The ratio $\pi^0/\eta$ and estimated background are shown in Fig. 1.a for WA70 and Fig. 1.b for NA24.
R808 EXPERIMENT

Background originating from cosmic rays and beam-gas interactions is rejected by requiring a vertex from charged particles in the crossing region and correct timing in hodoscopes located near the beam. The background for the detection of direct photons has been evaluated from Monte-Carlo calculation. The different contributions are shown in Figure 2 as a function of $P_T$.

**Fig. 1.a:** Ratio $\bar{\nu}/\nu$ for incident $\pi^-$. Background contribution due to halo (…) and total (---) background are shown.

**Fig. 1.b:** The ratio of $\bar{\nu}$ candidate/$\nu$ and the Monte-Carlo estimated background (dashed area).

**Fig. 2.a:** Contributions to $\bar{\nu}/\nu$ background.

**Fig. 2.b:** Observed $\bar{\nu}/\nu$ ratio. The dashed line is the sum of all background contributions.

UA2 EXPERIMENT

In the $P_T$ range considered (10-50 GeV/c), $\pi^0$ decays cannot be resolved and a large background from jets containing unresolved neutral multi-particles is present. However such backgrounds are generally accompanied by other jet fragments whereas direct photons are expected to be well-isolated.

**Fig. 3:** $\bar{\nu}$-$\nu$ energy flow plot of a direct photon candidate in UA2. The energy of the direct photon candidate is shown in black, the away side jet is clearly seen (white towers).
Therefore a direct photon candidate is defined as an electromagnetic energy deposition in the calorimeters without any associated charged track and satisfying isolation criteria. Energies are corrected for losses in the preshower detector.

Figure 3 shows the "lego" plot of a typical direct photon candidate in UA2. Beam halo background is removed by requiring: i) that small angle hodoscopes and calorimeter signals occur at the correct time, ii) that the missing transverse momentum of each event be less than 80% of the photon candidate transverse momentum.

The isolation criteria are:

In the central region:
- no charged track and at most one preshower signal may be found in a cone of \( \sqrt{\Delta x^2 + \Delta y^2} < 0.35 \) about the cluster direction
- the pattern of calorimeter photomultiplier signals must be consistent with the pattern expected for a single photon

In the Forward/Backward regions:
- no charged track may point to the cluster
- the total energy of all charged and neutral particles impinging on the cells adjacent to the cluster must not exceed 0.3 GeV.

The isolation criteria are intended to reject a large fraction of background and select direct photons. Evidence for this effect can be seen in Fig. 4, where the conversion rate observed in the preshower detectors, for isolated events (dots) and non-isolated events (triangles) is plotted, as a function of energy. \( \epsilon_{2\gamma} \) and \( \epsilon_\gamma \) are the conversion probabilities for two unresolved photons and a single photon, respectively. These probabilities are calculated from a Monte-Carlo simulation of the preshower detectors using the EGS\(^{14}\) program.

As a result of the different thresholds used to define a preshower signal in the central and Forward/Backward regions, \( \epsilon_\gamma \) behaves differently in the two regions. The simulations correctly describe the response of the preshower detectors to electrons from test beams. In the Forward/Backward regions, they also agree well with the extrapolation of data containing low-energy reconstructed \( \pi^0 \)s\(^{15}\). Since the observed conversion rate \( \alpha \) is between that expected for single and di-photons in both regions, it is clear that both samples have a substantial content of single photon events. The non isolated sample consists mostly of unresolved photon showers resulting from decays. Due to the effect of neutral multi-particle states, the observed conversion rate is slightly higher than \( \epsilon_{2\gamma} \). The calculation of \( \epsilon_{2\gamma} \) assumes that the ratio \( \eta / \pi^0 \) is 0.6\(^{16}\).

The direct photon sample contains a residual contamination of unresolved \( \pi^0 \) and \( \eta \) decays or multi neutral particles. Assuming that the contamination is due to single \( \pi^0 \) and \( \eta \) only, the
fractional contamination of the sample, \( b(P_T) \) is related to the converted fraction \( \alpha \) in the sample by:

\[
b(P_T) = \frac{(\alpha - \varepsilon_\gamma)}{(\varepsilon_{2\gamma} - \varepsilon_\gamma)}
\]

The conversion probability in the UA2 preshower detectors as a function of energy. The dots represent the conversion probability for the selected photon candidates and the triangles represent the conversion probability for candidates that fail the isolation criteria.

A pure sample of direct photons (\( b=0 \)) would give an observed conversion rate \( \alpha = \varepsilon_\gamma \), whereas a pure sample of \( \pi^0 \)'s and \( \eta \)'s (\( b=1 \)) would give \( \alpha = \varepsilon_{2\gamma} \). The contribution of unresolved multi-\( \pi^0/\eta \) states has the effect of increasing \( \varepsilon_{2\gamma} \). This contribution is evaluated using the ISAJET\(^{17}\) program. The resulting background fraction \( b(P_T) \) averages to \( 0.28\pm0.09 \) (stat)\( \pm0.04 \) (syst) in the central region and to \( 0.70\pm0.06 \) (stat)\( \pm0.04 \) (syst) in the Forward/Backward regions. The systematic errors come mainly from the evaluation of the multi-\( \pi^0/\eta \) states contribution.

Since hadronic collisions produce a number of spectator particles, the isolation criteria induce a loss of real events. The efficiency of the isolation criteria is evaluated from samples of minimum bias events and \( W \rightarrow e\nu \) events. The hadron contamination to the direct photon sample has been estimated statistically using the energy deposition in the hadronic calorimeters, to be less than 4% of the \( \gamma,\pi^0 \) contamination.

**IV. RESULTS**

**IV.I \( \gamma/\pi^0 \) RATIO**

Although it is not predicted by perturbative QCD, the ratio of direct \( \gamma \) production to \( \pi^0 \) production is one of the favourite quantities quoted by experimenters. This ratio explicitly exhibits
the direct $\gamma$ signal compared to the estimated magnitude of the background. As mentioned in the introduction, the $\gamma/\pi^0$ ratio should increase with $P_T$. The reported four experiments, for different reactions and over a large range of center of mass energy, observe such an increase.

Results from NA24, after background substraction, are shown as a function of $P_T$ in Figure 5.a for $\pi\pi$ interactions. Error bars correspond to the statistical error only. Corresponding results from WA70 are shown in Figure 4.b. One can already notice that, at a given $P_T$, this ratio is slightly higher for $\pi\pi$ interactions than for $\pi^+\pi^-$ interactions.

R808 has measured this ratio both for $p\bar{p}$ and $p\bar{p}$ collisions, it is compare to QCD predictions on Fig. 6. No significant difference is observed between $p\bar{p}$ and $p\bar{p}$ collisions.

The measurement of the $\pi^0$ cross-section made by UA2 in the Forward/Backward regions at $\sqrt{s}=540$ GeV\(^1\) has been repeated at $\sqrt{s}=630$ GeV. The $\gamma/\pi^0$ ratio measured at a mean pseudorapidity of $<\eta>=1.4$ is displayed, in Fig. 7, as a function of $x_T$.

This ratio is compared to the ratio measured in $p\bar{p}$ collisions by R808. The increase, at the same $x_T$, from ISR to collider energy is well predicted by a first order QCD calculation using the
parametrization of structure functions of Ref. 19 and the fragmentation function of Ref. 20. The calculation of the cross-sections ratio at $\sqrt{s} = 53$ GeV is shown as the dotted curve.

The same calculation at $\sqrt{s} = 630$ GeV is shown as the dashed curve, the effect of scaling violation in the fragmentation function is investigated by repeating the calculation with no such scaling violations (shown as the dotted-dashed curve).

![Figure 7: The ratio $\pi^-/\pi^+$ as a function of $x_T$. The solid points are the result of UA2. The result of R808 is shown as open circles. A set of QCD calculations is also shown (see text).](image)

IV.2. BEAM RATIOS

Using beams of particles and anti-particles, one expects to separate the two main contributions to direct photon production. Because of the large content of valence anti-quarks of $\pi^-$, in $\pi^-P$ collisions the annihilation process should be greatly enhanced at large $x_T$, compared to $\pi^+P$ collisions and should contribute substantially to the direct photon cross-section. However in the $x_T$ range covered by the experiments the Compton term (identical for $\pi^-P$ and $\pi^+P$ reactions) contributes predominantly to the cross-section. Differences are expected to be small and difficult to exhibit. NA24 has observed a larger direct photon cross section in $\pi^-P$ collisions than in $\pi^+P$ collisions (Fig.8.a). WA70 has observed a slightly higher $\pi^-/\pi^+$ ratio in $\pi^-P$ interactions compared to $\pi^+P$ (Fig.8.b). These ratios obtained by the two experiments are compatible with the theoretical expectation but deviation from unity has a small statistical significance.

![Figure 8.a: $\pi^-/\pi^+$ beam ratio for NA24. The upper curve is taken from ref. 21 while the dashed band is taken from ref. 22.](image)

![Figure 8.b: $\pi^-/\pi^+$ beam ratio for WA70. The upper curve is taken from ref. 21 while the dashed curve is taken from ref. 22.](image)
At the ISR, no significant difference is observed between PP and P\bar{P} collisions (Fig. 6). It should be mentioned that at the highest \(P_T\) measured the annihilation term should contribute at most for 25% to the direct \(\gamma\) cross-section (R808 has not observed any difference in the production of \(\pi^\pm\) between PP and P\bar{P} collisions).

In conclusion to this section, better statistics are needed for a clear observation of the annihilation contribution.

**IV.3. CORRELATIONS**

The measurement of the distribution of particles densities, as a function of the azimuthal angle \(\phi\) relative to the photon direction, should provide informations on the direct photon production mechanism. Except for bremsstrahlung contributions, prompt photons should not be accompanied by particles along their direction. The Compton process produces a direct photon in association with a quark jet in the opposite direction, whereas in the annihilation process a direct photon is associated with a gluon jet.

Both NA24 and WA70 have measured the particle density as a function of azimuth. Distribution are shown in Fig. 9 for direct \(\gamma\)'s and \(\pi^\pm\)'s. The difference on the trigger side (\(\Delta \phi \approx 0^\circ\)) suggests that bremsstrahlung contribution to direct photon production is small, and \(\gamma\)'s are well isolated whereas \(\pi^\pm\)'s are part of jet fragments. On the away side (\(\Delta \phi \approx 180^\circ\)) no significant difference is observe between \(\gamma\)'s and \(\pi^\pm\)'s. Similar results have been obtained by the R807 experiment at the ISR\(^{23}\).

![Figure 9: Particle densities associated with a high \(P_T\) direct photon as measured by NA24 and WA70.](image)
IV.4. CROSS-SECTIONS

As mentioned in the introduction, direct photon cross-sections have been calculated to next-to-leading order\(^2\,^3\). They are about a factor of 2 higher than leading order cross-sections. Prompt \(\gamma\) physics has become an excellent quantitative testing ground of QCD, though experimental measurements are difficult and the yield is small. Cross section measured by the NA3 and the R806 experiments are in agreement with next-to-leading order calculations\(^2\,^4\). Recently NA24 and UA2 have measured the direct-photon production cross-section.

NA24 results are shown on Figure 10, for \(\pi^-P\), \(\pi^+P\) and \(PP\) interactions, together with a next-to-leading order calculation of Aurenche et al.; a remarkable agreement between data and theory is obtained.

UA2 has measured the direct photon production cross-section in two different pseudorapidity intervals (\(|\eta| \leq 0.85\), \(1.1 \leq |\eta| \leq 1.7\)). Figure 11 compares the results and a next-to-leading order prediction of Aurenche et al.\(^2\,^5\) (upper curve). The lower curve corresponds to the same prediction but excluding bremsstrahlung contributions with an angle less than 45°, in order to simulate the effect of the isolation criteria which reject a photon too close to a jet or within a jet. Again a remarkable agreement is obtained. One can notice the high values of \(P_T\) obtained (up to 50 GeV/c); it corresponds however to low \(x_T\) (0.15 at maximum) a region where the Compton contribution dominates over the annihilation (up to \(P_T \approx 60\) GeV/c). On the same figure, the measured \(\pi^0\) cross-section in the Forward/Backward regions is also shown.

![Fig. 10: Direct photon cross-sections for three incident beams as measured by NA24. The solid lines represent a QCD calculation from ref. 21. A 2% systematic uncertainty on the energy scale has not been included in the error bars. The systematic errors (brackets) which are due to the uncertainty on the Monte-Carlo background estimation are added linearly to the statistical errors.](image-url)
The invariant inclusive cross-section for direct photon production (black dots) as at $\eta = 0$ and at $1.1 \leq |\eta| \leq 1.7$ as measured in UA2. Upper curves show prediction from ref. 25, lower curves result from excluding bremsstrahlung photon with quark-photon angle less than 45°. A 20% systematic uncertainty as not been include in the error bars. Also shown the invariant $\pi^+$ cross-section for production (open circles) at $1.1 \leq |\eta| \leq 1.7$.

Since many measurements are now available at different center of mass energies and different rapidities, one would like to compare them to calculations using different parametrizations of the gluon structure function in order to choose among them.

**IV.5. $\gamma$/JET RATIO**

UA2 has measured the ratio of the direct photon cross section to the jet cross section, in the $P_T$ range between 30 and 50 GeV/c:

$$\frac{\sigma_{\gamma}}{\sigma_{\text{jet}}} = 2.9 \pm 0.9 \text{ (stat.)} \pm 1.2 \text{ (syst.)} \times 10^{-4}$$

where the systematic errors are due to the uncertainty of the jet energy scale. In this $P_T$ range a lowest order QCD calculation using different parametrizations of the structure function predicts a ratio in the range between 3 and $4.5 \times 10^{-4}$. It is interesting to note that the $\gamma$/jet ratio found at the ISR, for different values of $P_T(x_T)$, has a similar value, approximately $10^{-4}$.

**V. CONCLUSIONS**

A clear direct photon signal has been observed by all experiments (NA24, R808, UA2 and WA70) over a large range of center of mass energies (from $\sqrt{s} = 23$ GeV to 630 GeV), for different rapidities and in many different reactions ($\pi^+P$, $\pi^+P$, $PP$ and $P\bar{P}$).

Results are in agreement with theory concerning: cross-sections, beam ratios and $\gamma$/jet ratio.
More statistics are needed to give information on the gluon structure and fragmentation functions, the bremsstrahlung contribution and to separate the annihilation from the Compton contribution.

ACKNOWLEDGMENTS
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O. Botner et al., Nucl. Inst. Meth. 196 (1982) 314;

5) UA2 Collaboration, M. Banner et al., CERN/SPSC/78-54

6) UA2 Collaboration, M. Banner et al., CERN/SPSC/78-54

7) WA70 Experiment:
M. Bonesini et al., Prompt photons production by π− and π+ on protons at 280 GeV/c, contributed papers P13-4; HEP85 Bari and 308, ISLEPH85 Kyoto.


15) For a pure sample of single photons $\epsilon_\gamma$ depends on the converter thickness and on the preshower counter threshold (T). In the central preshower detector T is set at 2 mips with a converter thickness of 1.5 X_0, $\epsilon_\gamma$ is then roughly independent of energy in the range considered. In the Forward/Backward regions, T is set at 6 mips, with a converter thickness of 1.4 X_0, resulting in a dependence of $\epsilon_\gamma$ upon energy which has to be taken into account.

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21) Ref. 2 and private communication.
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**TABLE 1**

<table>
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<th>EXPERIMENT</th>
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<th>WA70</th>
<th>R808</th>
<th>UA2</th>
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</thead>
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<td>300 GeV/c beams π^-, π^+, P on H₂ target</td>
<td>280 GeV/c beams π^-, π^+, P on H₂ target</td>
<td>√S = 53 GeV PP, PP</td>
<td>√S = 630 GeV PP</td>
</tr>
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<td></td>
<td></td>
<td></td>
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<tr>
<td>PT RANGE</td>
<td>4 - 7 GeV/c</td>
<td>4 - 7 GeV/c</td>
<td>3 - 6 GeV/c</td>
<td>12 - 50 GeV/c</td>
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<tr>
<td>σ/√E</td>
<td>0.28</td>
<td>0.15 + 0.04 √E</td>
<td>0.09 - 0.06</td>
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<td>METHOD</td>
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<td>CONVERSION</td>
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<td>σ(Mπ⁺)</td>
<td>15 MeV</td>
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<td>15 MeV</td>
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We present preliminary results on the invariant cross section for inclusive $\pi^*$ production, the $\eta/\pi^*$ production ratio, and the ratio of direct photons to $\pi^*$ production in $\bar{p}p$ collisions at $\sqrt{s} = 24.3$ GeV in the range $2.5 < p_T < 4.5$ GeV/c. These results are from data collected during a start-up run of the experiment in the fall of 1984 with an integrated luminosity of $4.0 \times 10^{34}$ cm$^{-2}$. 
The aim of experiment UA6 [1] is to study, in \( \bar{p}p \) and pp collisions at \( \sqrt{s} = 24.3 \text{ GeV} \), inclusive electron, \( \pi^+ \), \( \eta \) and \( \gamma \) production in the central region at large values of \( x_T = 2p_T/\sqrt{s} \), and low-t elastic scattering and diffractive dissociation. We report here preliminary results on detected electromagnetic showers from pp collisions.

The experiment, shown in Fig. 1, is situated in the main ring of the SPS. It consists of an internal hydrogen cluster-jet target followed by a double arm spectrometer. Both arms are identical, each arm with a conventional magnetic spectrometer with multiwire proportional chambers before and after a 2.4 Tm dipole magnet, a lead/proportional tube gas sampling calorimeter, an ionization (dE/dx) chamber, and a transition radiation detector. The acceptance of each arm is 20 to 100 mrad in the polar angle \( \theta \) and about 75° in azimuth for an integrated geometrical acceptance of 1.8 sr. At full jet intensity with \( 5 \cdot 10^{10} \) antiprotons in the main ring the luminosity is \( 5.2 \cdot 10^{29} \text{cm}^{-2}\text{s}^{-1} \). The luminosity is monitored by measuring the elastic cross section using a set of four silicon detectors set near \( \theta = 90° \) in the laboratory frame.

The electromagnetic calorimeter [2] consists of 30 lead plates, each 4.0 mm thick, interleaved with alternating planes of horizontal and vertical gas filled proportional tubes for a total of 24 radiation lengths and 0.75 nuclear interaction lengths. The transverse dimension of each tube is 1.0 cm. The calorimeter is longitudinally divided into three identical modules each containing 5 horizontal and 5 vertical planes of tubes. In each module the analog signals are summed in depth providing a longitudinal sampling of the shower every 8 radiation lengths in both horizontal and vertical views. A hodoscope of seven horizontal scintillation counters placed between the first two modules was used to trigger on electromagnetic energy deposited in the calorimeter.

Test beam results show the energy resolution of the calorimeter to be \( \sigma(E)/E = 33\%/\sqrt{E} \) with the response linear to within 3% from 10 to 100 GeV/c. Hadron rejection is approximately \( 10^4 \). The shower position resolution is \( \sigma = 2\text{mm} \). The calorimeter was calibrated in a test beam, the calibration being monitored and updated on a run to run basis by centering the \( \pi^+ \) mass peak on 0.135 GeV/c².

A simple clustering algorithm is used to identify showers and to reconstruct their energy and position. In each view and module clusters are first identified and then required to be associated in depth from module to module requiring that a line through the cluster centers extrapolate back to the target. Two clusters in the horizontal and vertical views are matched to form a shower if their energies are equal. Fig. 2 shows the mass distribution of all two cluster combinations with a total \( p_T > 2.5 \text{ GeV/c} \). There are clear peaks from the two-photon decays of the \( \pi^+ \) and \( \eta \) mesons. All two-cluster combinations with an invariant mass between zero and 0.3 GeV/c² were used for the \( \pi^+ \) cross section measurement. These candidates also were required to have a cluster separation at the face of the calorimeter greater than 3.0 cm and both showers were required to have an energy greater than 7 GeV.
The acceptance was calculated with a hybrid Monte Carlo using calorimeter test beam data to model the electromagnetic showers. Fig. 3 shows the invariant cross section for inclusive $\pi^+$ production as a function of $p_T$. The errors are statistical. For comparison the same cross section measured by the CCRS collaboration from pp collisions at $\sqrt{s} = 22.5$ GeV is plotted [3]. Within our present normalization uncertainty we see no difference between the $\pi^+$ invariant cross section in $\bar{p}p$ and pp collisions.

To measure the ratio of $\eta$ to $\pi^+$ production the background under the $\eta$ peak was subtracted by fitting a Gaussian plus a polynomial function to the data. The observed $\eta$ production rate was then corrected for the $\eta \rightarrow 2\gamma$ branching ratio, and the acceptance. The resultant $\eta/\pi^+$ cross section ratio is shown in Fig. 4. The ratio has a value of about 0.5. This is consistent with other measurements made in $\bar{p}p$ and pp collisions. Errors are statistical.

The fine transverse segmentation of the calorimeter makes it well suited for the study of direct photons where the background from $\pi^+$ and $\eta$ decays can dominate. For a 100 GeV/c $\pi^+$ the minimum separation of the two photons is 2.8 cm at the calorimeter face which produces two showers readily distinguishable from a single shower. Higher energy $\pi^+$'s have showers which begin to overlap but which can be separated from single showers by their width. An electromagnetic shower was considered to be a direct photon candidate if it could not be paired with another shower to form either a $\pi^+$ or an $\eta$ and if the shower radius was consistent with that produced by a single photon. Fig. 5 shows the ratio of these direct photon candidates to $\pi^+$'s corrected for acceptance. These data are very preliminary. In particular no systematic errors are included. For comparison we have plotted the theoretical expectations obtained from a computation of the expected $(\bar{p}p\rightarrow\gamma X)/(pp\rightarrow\gamma X)$ ratio by Contogouris and Mebarki [4] and direct photon measurements in proton-proton collisions of experiment R-806 [5].

REFERENCES


Fig. 1

Bottom calorimeter arm

$E_{T} = 25 \text{ GeV/c}$

Fig. 2

Fig. 3

Fig. 4

Fig. 5
This article surveys the results of experiment NA14 on the point-like scattering of photons by quarks. Both QED Compton scattering, (leading to a prompt $\gamma$) and QCD Compton scattering $\gamma q \to g q$ (seen as $\pi^0$ and $\pi^\pm$ production), are consistent with QCD calculations. This is a quantitative test of perturbative QCD up to order $a_s^2$. It is also shown here that QED Compton scattering events can be topologically isolated and measured.
The motivation behind the NA14 experiment has been the study of point-like interactions of photons with quarks. These reactions can be represented by simple first-order diagrams:

\[ \gamma q \rightarrow \gamma q \] QED Compton scattering

and

\[ \gamma q + g q \] QCD Compton scattering.

QCD corrections to both processes beyond leading order have been calculated \(^1\),\(^2\) and we have compared the predictions with experiment. The main background to be subtracted has been estimated using Vector Meson Dominance (VDM) to model the hadronic interaction of the photon; this is the dominant process at all but high transverse momenta.

The NA14 experiment (described at previous Moriond meetings \(^3\) and in publications \(^4\),\(^5\)) is designed to enable the comparison to be made for both QED and QCD Compton scattering. In addition to summarising these results we shall show here that the reaction \( \gamma q + \gamma q \) can be measured also through the characteristic topological properties of the events.

**Important features of the NA14 experiments:**
- high intensity, broad-band \( \gamma \) beam, mean energy 80 MeV;
- good angular coverage for both \( \gamma \) and charged particles;
- normalization (±20%) measured by several independent methods;
- \( \pi^- \) beams can be sent onto the target, allowing comparison with the \( \gamma \) beam and enabling reliable estimation of VDM effects and also of \( \pi^0 \) backgrounds.

**QED Compton scattering**

Evidence for the process \( \gamma q + \gamma q \) was obtained by using the \( \pi^- \) beam data to evaluate reliably not only the contribution from VDM but also to evaluate the fraction of single gammas which were misidentified \( \pi^0 \)'s. Thus we measured the cross-section for the process \(^4\). This cross-section agrees quantitatively with the evaluation of Born terms with the addition of QCD corrections \(^1\) - a confirmation that the QCD calculations are satisfactory. Moreover this result is not consistent with the gauge integer-quark model \(^6\), although there is the reservation that with this theory the calculations of corrections to the Born terms have not been made; see also Lipkin \(^7\).

**QCD Compton scattering**

The reaction \( \gamma g + g q \) is responsible for much of the production of pions with high \( p_T \). We have published our results for \( \pi^0 \) production and present here results for \( \pi^\pm \). The Born terms for this process have QCD corrections calculated by Aurenche et al. \(^2\). Fig.1 (a) and (b) show the \( p_T \) distributions of \( \pi^\pm \) for \( \pi^- \).
beam and $\gamma$ beam, respectively. The $\pi^-$ data has been used to estimate the VDM contribution. The data clearly need the QCD correction. Fig. 2 shows the increase in the QCD contribution with $p_T$ and that the fit is not sensitive to the choice of gluon fragmentation functions.

![Graph](https://example.com/graph1.png)

**Fig. 1** Inclusive production of $\pi^\pm$ plotted as a function of $p_T$, (a) for $\pi^-$ beam, (b) for $\gamma$ beam. The VDM contribution to photoproduction is derived from the $\pi^-$ beam data; the predictions for QCD Compton scattering include higher order corrections.

![Graph](https://example.com/graph2.png)

**Fig. 2** The number of observed $\gamma \to \pi^\pm$ events divided by the predicted number is plotted as a function of $p_T$ for three models: VDM alone (o), VDM + QCD with soft (•) and with hard (quark-like) (△) gluon fragmentation functions.

Topological isolation of deep QED Compton scattering

It is interesting to ask if the reaction $\gamma q + \gamma q$ can be identified by a characteristic final state: $\gamma + quark
t jet$. To study this possibility objectively we attempt to separate them from the principal background, $\gamma q + gq$. We take $\gamma$-trigger events with $p_T > 3$ GeV/c. If we can reconstruct a $\pi^0$ using the
Fig. 3 Examples of $\gamma q + \gamma q$ events; the tracks are plotted in the $\theta, \phi$ plane with the energy represented by the height of the bin. The trigger $\gamma$ is on the right, the bunch of tracks opposite correspond to the quark jet.

Fig. 4 Distributions of track-trigger distance $d$ (defined in text) for candidate $\gamma \to \gamma$ events and $\gamma + \pi^0$ events for two ranges of $p_T$. The graph on the right has been generated by the Lund Monte Carlo for pure QED Compton scattering $\gamma q + \gamma q$ with the fragmentation of the quark.
trigger and another $\gamma$ we call it a "$\pi^0$", otherwise a "$\gamma$". We find for these triggers:

<table>
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<th>Trigger</th>
<th>.clientX</th>
<th>.clientY</th>
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<th>Height</th>
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<tr>
<td>$\gamma$</td>
<td>53 ± 5</td>
<td>36 ± 4</td>
<td>32</td>
<td>5</td>
</tr>
<tr>
<td>$\pi^0$</td>
<td>43 ± 6</td>
<td>27 ± 5</td>
<td>32</td>
<td>5</td>
</tr>
</tbody>
</table>

% compatible with QED Compton: (43 ± 6)% (27 ± 5)%

Thus there is clearly a topological difference between the two classes of event. Some selected candidate "$\gamma$" events are shown in Fig.3. We can make the distinctions more quantitative by defining a track-trigger distance squared:

$$d^2 = (\Delta \phi)^2 + (\Delta \eta)^2$$

where $\Delta \phi$, $\Delta \eta$ are the differences in azimuth, pseudorapidity, respectively between any track and the trigger.

Fig.4 shows that the distance-squared distributions are similar at low $p_T$ but are different at high $p_T$. Using in addition a Monte Carlo generated distribution for pure QED Compton scattering (Lund with quark fragmentation) and fitting the distribution with polynomials, $f_{\pi^0}$ and $f_{\text{QED}}$ we are then able to fit the "$\gamma$" distribution to the form

$$f_\gamma = (1 - \epsilon) f_{\pi^0} + \epsilon f_{\text{QED}}$$

where $\epsilon$ is the fraction of triggers which are QED Compton scattering. The values obtained agree with what is predicted.

In summary we have been able to identify reactions due to point-like interactions of the photon and our results are in quantitative agreement with QCD calculations beyond leading order.

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PERTURBATIVE QCD AND JETS

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Abstract

A brief review of some of the recent progress in perturbative QCD is given.
1. Introduction

In this talk I shall give a brief summary of some of the topics which have received considerable attention in studies of perturbative QCD over the last year. This is by no means an exhaustive account of what has been done, but rather a sampling of the progress and new understanding which has been achieved.

2. Heavy Quark Production

I would like to separate the discussion of heavy quark production into two parts, (A) heavy quark production in gluon jets and (B) heavy quark production in hadronic collisions in general.

2.A. Heavy Quarks in Gluon Jets

In leading orders of QCD, the number of heavy quark pairs in a gluon jet of energy $Q$, is given by

$$
\rho = \frac{1}{3\pi} \int_4^{Q^2} \frac{dK^2}{K^2} \alpha(K) \left( 1 + \frac{2M^2}{K^2} \right) \left( 1 - \frac{M^2}{K^2} \right) \mathcal{N}_g(Q^2,K^2)
$$

(1)

where

$$
\mathcal{N}_g = \left( \frac{\ln Q^2/\Lambda^2}{\ln K^2/\Lambda^2} \right)^a \frac{\exp \left( \frac{2C_A}{\pi b} \ln \frac{Q^2/\Lambda^2}{K^2/\Lambda^2}^{1/2} \right)}{\exp \left( \frac{2C_F}{\pi b} \ln \frac{Q^2/\Lambda^2}{K^2/\Lambda^2}^{1/2} \right)}
$$

(2)

$M$ is the mass of the heavy quark and lowest order perturbation theory is obtained by setting $\mathcal{N}_g$, the gluon multiplicity at mass $K$, equal to 1.

One reason which makes a measurement of $\rho$ especially interesting is that one can do a reliable estimate of the non-perturbative corrections to (1). The leading non-perturbative corrections can be expressed in terms of $\langle \alpha F^2 \rangle$, the gluon condensate introduced by Shifman, Vainshtein and Zakharov. $\langle \alpha F^2 \rangle$ is one of the parameters which must be added to perturbative calculations to take into account the fact that the perturbation series is not Borel summable. Thus in the limit of zero mass light quarks one finds

$$
\delta \rho = \frac{\langle \alpha F^2 \rangle}{N^2-1} \frac{\alpha(M)}{M^4} \left( -\frac{1}{30} C_f + \frac{53}{3780} C_A \right)
$$

(3)

with $N$ the number of colors and $C_f$ and $C_A$ the Casimir operators in the fundamental and adjoint representations respectively. (3) gives a very small correction to (1).
Light quark condensates can contribute an additional correction to $\rho$ proportional to

$$\frac{m(\bar{\psi}\psi)}{M^2}.$$ 

Though an exact calculation of light quark condensates to $\rho$ has not been performed, they likely make an extremely small contribution also.

Thus Eq. (1) should give a very accurate prediction for the number of heavy quarks in a gluon jet with the main uncertainty coming from the values of $A$ and $M$ chosen. The $M$ appearing in (1) is the current algebra quark mass so that a good measurement of $\rho$ could give important information toward determining this bare mass. Of course, especially at present energies, charm production in gluon jets is the dominant flavor in such heavy quark production. I believe that it is very interesting and important to make a good measurement of $\rho$ for this process.

2.B. Heavy Quark Production in Hadronic Reactions

Significant progress in understanding heavy quark production has been made in the last year or so. In particular the concepts of gluon-gluon fusion, flavor excitation and intrinsic heavy quarks have undergone a significant revision in the past year. At present the experimental situation is still unsettled with regard to charm production in hadronic collisions though we may expect considerable clarification in the next year or so. At present there is no data on production of quarks heavier than charm in hadronic collisions.

There are three firm statements which can now be made about heavy quark production which were not completely apparent a year ago. (i) At leading order in $\alpha(M)$ the gluon-gluon fusion term, $gg \rightarrow Q\bar{Q}$, shown in Fig. 1, should give the dominant contribution to heavy quark production as

$$\sigma = \int dX_1 dX_2 G(X_1, M^2) G(X_2, M^2) \sigma_{gg \rightarrow Q\bar{Q}}.$$  (4)

$\sigma_{gg \rightarrow Q\bar{Q}}$ is the order $\alpha^2$ contribution to $Q\bar{Q}$ production by gluons having momenta $X_1 p_1$ and $X_2 p_2$ respectively, while $M$ is the mass of the heavy quark. The total cross section is proportional to $\alpha^2(M^2)/M^2$. (ii) Intrinsic heavy quark components in the wave function of the colliding hadrons give a contribution of size $\alpha^2/M^4$ to the cross section. (iii) The concept of flavor excitation is no longer operative.

The arguments for (i), recently given by Collins, Soper and Sterman are really quite straightforward. The essential idea is that heavy quark production involves momentum transfers at least of order $M$ and so constitutes a hard process taking place over short times. The process is then essentially no different than massive $\mu$-pair or $Z$-production, except that in heavy flavor production gluons play the dominant role. Eq. (4) can be systematically
improved in powers of $\alpha(M^2)$. We shall come back to the question of what energies are actually necessary to realize the dominant role played by the lowest orders of perturbation theory.

With respect to (ii) high energy hadrons undoubtedly have intrinsic heavy quark components in their wave functions with probability of size $\Lambda^2/M^2$. It requires a hard scattering to set free these heavy quarks, thus resulting in a production cross section of size $\Lambda^2/M^4$. Such a cross section is certainly negligible for quarks heavier than charm and is likely true for charm.

With regard to (iii) the idea of flavor excitation, illustrated in Fig. 2, was that a heavy quark pair might exist in the wave function of a fast hadron, however, in a perturbative rather than in an intrinsic (higher twist) sense. While undergoing a scattering with another hadron this heavy quark system would then be set free, generally in the fragmentation region of the hadron of which it had been a part. This view of large-$X$ charm production is not incorrect, however, this component is already included in the gluon-gluon fusion contribution.

Despite the progress made in understanding heavy flavor production in hadronic collisions there remain many questions which are still unanswered or for which only partial answers exist.

1. I have said that gluon-gluon fusion, as shown in Fig. 1, dominates heavy flavor production at high energies. Consider now graphs of the type shown in Fig. 3. Indeed, such graphs are order $\alpha(M^2)$ smaller than the leading contribution, however, the fact that

$$\frac{\sigma_{gg+Q\bar{Q}}}{\sigma_{gg}} \propto 100$$

at order $\alpha^2$ means that higher order contributions could be competitive with the leading term. Whether or not this is so is not known at present. It is very important to evaluate the $\alpha(M^2)$ corrections to gluon-gluon fusion before one can have any confidence in theoretical calculations of heavy flavor production.

2. Recently Ellis and Leon have done the order $\alpha(M^2)$ correction to gluon-gluon fusion in the large-$X$ region. They find a very small correction, but in the process of their work they noticed that the charm cross section predicted from gluon-gluon fusion depends strongly on the choice of the charm mass, much more strongly than the dimensional factor $1/M^2$. I think the origin of this dependence is clear. To produce a charm-anti-charm system at mass $M$ in a collision of center of mass energy $\sqrt{s}$ means that the $X_1$ and $X_2$ of the colliding gluons satisfy $X_1 X_2 s = M^2$. Especially at fixed target energies this requires $X$ values which are not too small and the rapidly falling gluon distribution strongly favors $X$ values as close to threshold as possible. Forcing $X$ close to threshold can vitiate the use of perturbation theory. The application of perturbation theory
requires that one be able to produce a quark-anti-quark pair in a mass region 2M ≤ M ≤ 2CM, with C a constant not close to 1, (C = 2 might be a reasonable choice), without encountering a strong X-dependence of the gluon distributions in that range of M. As an example take M = 1.5 GeV then perturbation theory should be valid so long as one does not encounter strong gluon structure function dependence in the X-region up to X ≈ \( \frac{4M}{\sqrt{s}} \), at least for perfectly central production. Now if the gluon distribution varies like \( XG(X) \sim (1-X)^n \), perturbation theory should be valid when \( \left( 1 - \frac{6}{\sqrt{s}} \right)^{10} \) is close to 1. (Recall that the gluon distribution comes in quadratically in gluon-gluon fusion.) Such a criterion cannot be met at fixed target or even ISR energies. Thus we would expect a strong dependence on the choice of the charm mass and in addition, we might expect a strong energy dependence as one goes from fixed target and ISR energies. However, at collider energies the mass dependence should again become weak with the cross section depending on \( M^2 \) only as \( 1/M^2 \). Here again, as in the case of heavy quarks coming from gluon jets considered in the previous section, a good measurement of the cross section for charm production when combined with a higher order calculation and a knowledge of the small-X gluon distribution could provide a determination of the bare charm mass.

3. Small-X Physics

Interest in small-X physics was stimulated a few years ago when the reliability of the Altarelli-Parisi equation at SSC energies was questioned because of the rapid growth in gluon and sea-quark distributions predicted in the small-X region. As we shall see in a moment there are indeed important corrections to the Altarelli-Parisi equation at very small values of X, but such corrections do not modify predictions for new article production at SSC energies since these modifications are only effective at mass values of a few GeV.

To see that the gluon distribution must increase rapidly at small values of X one may take the Altarelli-Parisi equation as a starting point. However, the Altarelli-Parisi equation requires input X-distributions at some \( Q_0^2 \), so we assume that \( XG(X,Q_0^2) \) is given. (Quark interactions are not very important in small-X evolutions and we shall only keep the purely gluonic part of the theory in our present discussion.) Then

\[
XG(X,Q^2) = \int_X^1 \frac{dX'}{X'} K\left( X/X', Q^2, Q_0^2 \right) X' G\left( X', Q_0^2 \right).
\]  \( \text{(6)} \)

now,

\[
K\left( X/X', Q^2, Q_0^2 \right) = \exp \left[ \frac{C}{\pi b} \ln \left( \frac{\ln Q^2/\Lambda^2}{\ln Q_0^2/\Lambda^2} \right) \ln X/X' \right].
\]  \( \text{(7)} \)
with $k$ a known and slowly varying function when $X'/X$ is large. If $X'G(X', Q^2)$ does not grow strongly with $X'$, the small-$X$ behavior of $XG(X, Q^2)$ will be determined by the growth of $K$ when $X/X'$ is small. In general, however, we have no reason to expect $X'G(X', Q^2)$ to be slowly varying. Indeed it likely grows\textsuperscript{7} something like $(X')^{-1/2}$ so long as $X'$ is not too small, so that the small $X$ behavior of $XG(X, Q^2)$ is determined both by the initial distribution and by $K$. Since the initial distribution is not calculable within perturbative QCD the small $X$ behaviour of $XG(X, Q^2)$ cannot be accurately calculated. Nevertheless the growth of $K$ for $X'$ fixed as $X$ becomes small gives a lower bound for the $X$ dependence of $XG(X, Q^2)$ and we may conclude that $XG(X, Q^2)$ grows rapidly for small $X$, at least in that domain where the Altarelli-Parisi equation is valid.

Now I would like to review the physical picture\textsuperscript{8} of this growth in $XG(X, Q^2)$ at small $X$. Suppose we consider a proton of momentum $p$, in a frame where $p = l_2 p$ and with $p$ chosen very large. In this frame $XG(X, Q^2)$ represents the number of gluons in the proton, per unit rapidity interval, with longitudinal momentum centered about $Xp$ and having transverse size $|\Delta b| \approx 1/Q$. Now when $XG(X, Q^2) \propto Q^2 R^2$, with $R$ the proton radius, the gluons within this unit of rapidity begin to spatially overlap in the longitudinally thin disc which they occupy. When $X$ is further decreased we may expect the gluons in this rather dense system to scatter and annihilate with one another, eventually reaching a saturation limit as $X \to 0$. When the gluon density becomes large enough that gluon interactions and annihilations are no longer negligible, the Altarelli-Parisi equation ceases to be valid and (6) does not represent the actual gluon density\textsuperscript{8}.

In the domain of $X$ and $Q^2$ where gluon interactions due to gluon crowding are small one can in fact give a modified Altarelli-Parisi equation which takes these interactions into account. In this case one has\textsuperscript{8,9}

$$Q^2 \frac{\partial}{\partial Q^2} XG(X, Q^2) = \frac{\alpha C_A}{\pi} \int_{X'}^{1} \frac{dx'}{x'} \frac{X}{X'} \gamma(X/X') X' G(X', Q^2)$$

$$- \frac{4 \pi^2}{N^2 - 1} \left( \frac{\alpha C_A}{\pi} \right)^2 \frac{1}{Q^2} \int_{X}^{1} \frac{dx}{x} \gamma(X^2) G^{(2)}(X', Q^2)$$

as the modified equation. $G^{(2)}$ is a two gluon correlation. The second term in (8) represents gluon recombination or, equivalently, gluon shadowing. (8) should be a valid equation for small $X$ so long as the second term is small compared to the first term. At $X$ and $Q^2$ values for which both terms on the right-hand-side of (8) are comparable this equation can no longer be trusted as higher gluon correlations also become important.

Nevertheless one should be able to use (8) to get an indication of where the non-linear effects become important. We may estimate the regions where gluon saturation occurs by setting the two terms on the right-hand-side of (8) to be equal and by using $\gamma(X) \to 1$ the small-$X$ limit of this anomalous
dimension. Let us do this estimate first for a large nucleus. Our normalization is such that

\[ \frac{G^{(2)}(X, Q^2)}{G(X, Q^2)} = \left[ \frac{G(X, Q^2)}{\frac{8}{9} \pi R^2} \right]^2 \]

where \( G^{(2)} \) is the two gluon correlation in a large nucleus of size \( R \) and \( G \) is the proton's gluon distribution. Anticipating \( Q^2 \approx 1 \text{ GeV}^2 \) we set \( \frac{\alpha C}{\pi} \approx \frac{1}{3} \) and obtain

\[ X G(X, Q^2) \approx 15Q^2 A^{-1/3} \]

as the boundary of saturation. For \( A^{1/3} \approx 5-6 \) it should be possible to achieve saturation of gluon densities in the region \( X \approx 0.02-0.03 \) with \( Q^2 \approx 1-2\text{GeV}^2 \). Such regions would be available in heavy ion collisions at RHIC and it is exactly these gluons which should eventually thermalize to give the central region of the quark-gluon plasma expected at such a collider.

An analogous estimate for proton-proton or proton-anti-proton collisions is a little more difficult because the proton is not really a loosely bound system. Nevertheless, in order to get an estimate of the \( Q^2 \) and \( X \)-values for which gluon saturation might occur in a high energy hadron-hadron collision we shall take as a crude model the picture in which a proton is a loosely bound system of three valence constituent quarks. Then, proceeding as in the nuclear case, one needs

\[ X G(X, Q^2) \approx \frac{25}{\pi a} Q^2 R^2 \]

for saturation, with \( R \) now the proton radius. For \( Q^2 \approx 1\text{GeV}^2 \) and \( R = 1\text{fm} \) one needs extremely small values of \( X \), probably near \( 10^{-4} \) in order to have a chance of reaching saturation. Clearly heavy ions appear more efficient in producing such dense systems of gluons, although we shall see in Section 4 that it is possible to produce a high gluon density at much lower energies and at much larger \( Q^2 \) over a severely restricted part of the proton.

The dense system of gluons which we have been talking about is part of the wave function of a high momentum heavy ion or proton. It will not be in either a kinetic or flavor equilibrium, although such an equilibrium may indeed take place after a collision. Are such high density systems of interest and if so, why are they of interest? What I find fascinating about such high density gluon systems is the fact that one is producing very high chromodynamic field strengths. One can see the size of the potential and field strengths involved by the following argument: in light-cone gauge \( X G(X, Q^2) \) represents the number of gluons of size \( |\Delta p| \approx 2/Q \) in a unit of rapidity centered about \( X \). \( X G \) also represents the size of \( A^2 \) due to quanta of this type. Now the linear and non linear terms in (8) balance when \( X G \approx Q^2/\alpha \). This means that saturation occurs, when making field measurements coherently.
over a transverse size $r \approx 2/Q$ at values $E^2 \sim B^2 \sim \frac{1}{r^2} \alpha$. These are very large field strengths, of such size that the individual quantum description of the field is no longer appropriate. If there are interesting non perturbative effects in QCD, at distances small compared to $1/\Lambda$, they would likely show up most clearly in such high field configurations. In any case, the physics of the regime of high field strength gluon saturation is that of non perturbative QCD though a non perturbative regime where the normal condensates reflecting the non perturbative vacuum are irrelevant.

4 - Minijets and Related Topics

Minijets, jets of transverse momentum greater than, say, 5GeV are strongly produced at the CERN collider having an inclusive cross section of order 10mb at the highest collider energies. Because $\alpha$ is small at $p_\perp = 5$GeV one has the possibility of explaining at least portions of this data in terms of perturbative QCD. To see that the problem of relating minijet data to perturbative QCD is in fact non trivial we begin by reminding the reader that the cross section for two jet production in the minijet regime is given by

$$\frac{d\sigma}{dy_1 dy_2 dp_\perp^2} = \prod_{i=1}^2 G(x_1, p_\perp^2) G(x_2, p_\perp^2) \hat{\sigma} \to gg$$

(12)

where $\hat{\sigma}$ is the elementary cross section as given, for example, in Ref. 10. (We shall see later that (12) is probably not a reliable starting point but this is the usual formula and will suffice for our present purposes). For simplicity we have limited ourselved to gluon jets, which should dominate minijet production. Then the total inclusive production for a minijet having $p_\perp > M$ is given by

$$\sigma(M, s) = \int \frac{dX_1 dX_2}{X_1 X_2} \frac{d\sigma}{dy_1 dy_2 dp_\perp^2}.$$ 

(13)

Now, is it possible to calculate the large $s$ behavior of $\sigma(M, s)$ for $M$ fixed? Clearly large $s$ values require small $X$ values of $G(X, M^2)$ and we have seen in the last section that one is not in a position to do a reliable calculation of the small $X$ behavior of $G$. Thus we feel that a reliable calculation of the energy dependence of minijet production is not possible purely within perturbative QCD. However, if accurate values of gluon distributions were available from deeply inelastic lepton measurements at small values of $X$ it should be possible to predict $\sigma$, and we shall indicate how this can be done a little later in this section.

In the remainder of this section I would like to describe a minijet measurement which can be carefully discussed within perturbative QCD. The discussion given below summarizes part of a piece of work I have done in collaboration with H. Navelet and which should be available in a more
complete version soon. The measurement is most easily described in terms of
the two jet inclusive cross section, $d\sigma / dy_1 dy_2 d^2 p_{1\perp} d^2 p_{2\perp}$.

$y_1 - y_{\text{max}} = \Delta_1$, with $y_{\text{max}} - y_2 = \Delta_2$, where $y_{\text{max}}$ and $y_{\text{min}}$ are the maximum and
minimum values of $y_2$ and $y_1$. Then define

$$\sigma(M, s) = \int d^2 p_{1\perp} d^2 p_{2\perp} \Theta (p_{1\perp} - M) \Theta (p_{2\perp} - M) \frac{d\sigma}{dy_1 dy_2 d^2 p_{1\perp} d^2 p_{2\perp}}$$  \hspace{1cm} (14)$$

where the variables $\Delta_1$ and $\Delta_2$ have been suppressed. We may use factorization to write

$$\sigma(M, s) = \left( \frac{\alpha C_A}{2\pi^2} \right)^2 \frac{\pi^2}{2 M^2} X_1 \left( G \left( X_1, M^2 \right) + \frac{4}{3} Q \left( X_1, M^2 \right) \right) X_2 \left( G \left( X_2, M^2 \right) + \frac{4}{3} Q \left( X_2, M^2 \right) \right)$$  \hspace{1cm} (15)$$

where $G + \frac{4}{3} Q$ is the usual gluon plus quark factor which appears in jet
physics. In lowest order perturbation theory $f = 1$ while in a leading
logarithmic approximation $f = f(\alpha y)$ with $y = \ln \left( \frac{X_1 X_2 s}{M^2} \right)$. We shall be mainly
concerned here with a discussion of the behavior of $\sigma(M, s)$, or $f$, in the
leading logarithmic approximation. The $\alpha$ which appears in (15) and throughout
the rest of this discussion is $\alpha(M)$ which is already small when $M$ is in the
range of 5-10GeV.

In the leading logarithmic approximation one finds$^{11,12}$

$$f = \frac{1}{2\pi} \int_{-\infty}^{0} \frac{d\nu}{\nu^2 + 1/4} e^{\frac{2\alpha C_A}{\pi} y x(\nu)}$$  \hspace{1cm} (16)$$

where $x(\nu) = -\gamma - \text{Re} \Phi(\frac{1}{2} + i\nu)$ with $\gamma$ the Euler constant and $\Phi$ the logarithmic
derivative of the $\Gamma$-function. For real values of $\nu$, $-x(\nu)$ is monotonically
decreasing function of $|\nu|$ away from $\nu = 0$. At large values of $\alpha y$ the
asymptotic behavior of $f$ is determined by the saddle point at $\nu = 0$ in (16).
In fact we are interested in intermediate values of $\nu$ since, as we shall see
later, the leading logarithmic approximation cannot be reliable for too large
values of $y$. In order to determine $f$ in an intermediate region of $y$ we have
evaluated the perturbation expansion for $f$ in some detail. For $y \ll 2$ the
leading term in the large $y$ expansion give a good representation of $f$ as

$$f \sim e^{\frac{2\alpha C_A}{\pi} y 4 \ln 2}$$  \hspace{1cm} (17)$$

while for $y \gtrsim 2$ perturbation theory in $\alpha y$ converges rapidly. Thus the
effective rate of growth for $y \gtrsim 3$ of $\sigma(M,s)$ is close to $\sqrt{s}$.

The growth, in $s$, indicated by (17) is the growth due to the bare perturbative Pomeron in QCD. What is new about our discussion is that for a limited range in $s$ this behavior should in fact dominate the cross section. Let us in fact determine the range in $y$ over which the bare perturbative Pomeron can be expected to correctly give $f$. In general we may write $f$ as

$$f = f_0(\alpha y) + \alpha f_1(\alpha y) + \alpha^2 f_2(\alpha y) + \ldots$$

where, now, $f_0$ contains the leading logarithms. Now the leading logarithmic approximation breaks down when $\alpha f_1$ and $\alpha^2 f_2$ etc are comparable to $f_0$. This happens when

$$y = \frac{2}{\ln(C/\alpha)} \ln(C/\alpha)$$

with the constant in (19) undetermined at present. When $y$ attains the value indicated in (19) the higher Pomeron exchanges become important and presumably slow the rate of growth indicated in (17) to a $\ln^2 s$ growth. The picture should be something like that shown in Fig. 4.

The cross section given by (15) is not small and is certainly in the 1-10mb region through the energy range of the CERN collider. The growth produced by (15) as one increases $s$ is due to multijet production. That is, given the two jet trigger we have been discussing it is very likely that additional jets will be present, at least when $y$ is not too small. The trigger for the two jet inclusive process picks out a parton-parton cross section which, in terms of the quantities appearing in (15), is given by

$$\left(\frac{\alpha C}{\pi}\right)^2 \pi^2 \frac{f}{2M^2}$$

What we have tried to show here is that the high energy behavior of this parton-parton cross section involves $j$-plane singularities, as usual, but that now one can justifiably use weak coupling techniques although ultimately, when (19) is attained this weak coupling regime will correspond to a non perturbative high field strength regime; a high field strength regime now, however, only over a small part of the proton.

5. Classical Simulations in High Energy Reactions

Recently a number of very ambitious schemes have been proposed for studying complete events in $e^+e^-$ and in hadronic collisions$^{13-19}$. The practical motivation is clear in that such information is needed in order to estimate backgrounds and uncertainties in a given detector. Theoretically,
these programs are quite interesting because they attempts to represent a complicated quantum mechanical process in terms of classical time evolution. In fact at first sight it is not at all clear that Monte Carlo simulations of high energy events have anything but rough empirical validity. The situation in e^+e^- collisions is simpler so let me begin there.

5.A. Classical Simulations in e^+e^- Annihilation

The models developed to describe e^+e^- annihilation events are of types, branching models and string models. Both types of models are quite successful in explaining quite detailed properties of e^+e^- events. The models are, however quite different in spirit.

Branching models can be schematically viewed as shown in Fig. 5. Immediately after the virtual photon has converted into an energetic quark-anti-quark pair, the quark may emit gluons which then emit more gluons. The probability of a quark or gluon emitting another gluon is taken to be proportional to perturbative QCD matrix elements. When the quarks and gluons go below a mass Q_0, they are converted into hadrons by some non-perturbative hadronization model, usually a cluster or string model. The evolution of individual quanta between mass Q, the mass of the virtual photon and Q_0 is thus given by the branching note that the procedure is classical in that there is no addition of amplitudes and no interference in the branching. The Q^2-dependence of events is given only by the branching part of the model and does not depend on the method of hadronization.

How can one possibly hope to represent a complicated quantum system without taking interference into account? That this is not completely impossible comes about because of two key facts. (i) In large regions of phase space there is effectively only one non-negligible Feynman graph. (ii) In another large region of phase space, where in fact many Feynman graphs interfere, interference is complete and the quantum amplitude is exactly zero. In the classical simulation one may simply suppress these regions of phase space. It is a remarkable fact that the Marchesini-Webber branching algorithm gives results which agree with perturbative QCD in a rather systematic way. To make this statement more precise, define

\[ \sigma_{n_1 n_2 \ldots n_r} = \frac{1}{\sigma} \int \frac{d\sigma}{dX_1 \ldots dX_r} X_1^{n_1} \ldots X_r^{n_r} dX_1 \ldots dX_r \]  

Then

\[ \sigma_{n_1 \ldots n_r} / (n)^{\gamma} = C_{n_1 \ldots n_r} \left( 1 + \sum \sqrt\alpha + O(\alpha) \right) \]  

with the constants C and d correctly given by the Marchesini-Webber algorithm.

Nevertheless a number of questions remain with respect to branching. (i) Is it possible to include, systematically, order \( \alpha \) and higher corrections? (ii) Perhaps even more importantly, is it possible to get the right \( Q^2 \)
dependences at the constant and $\sqrt{x}$ levels for observables which involve angular correlations? (Note that all angles have been integrated out in the expressions in (21) and (22)). As we shall see a little later the Marchesini-Webber model has angular correlations only partially built in. At present it is not at all clear that any classical branching model can give angular correlations in a systematic way.

- String models were originally introduced as an attempt to follow the hadronic time evolution of a high energy collision by Preparata and his collaborators\textsuperscript{20}. They have been most completely developed and successfully utilized by the Lund collaboration. In Fig. 6 I have illustrated the physical picture where a virtual proton converts into a quark-anti-quark pair which then separates and form a thin flux tube whose ends connect to the pair. After a certain amount of stretching the flux tube breaks when a new quark-anti-quark pair is created out of the vacuum. The process continues until the final strings become physical hadrons. The string itself does not have vibrational degrees of freedom, but it does have "kink" degrees of freedom corresponding to hard gluon emission.

The Lund model has been very successful phenomenologically. The relationship of the model to QCD is, however, not yet clear. In particular the early stages of evolution which should be governed by perturbation theory do not appear to have any connection with the Feynman diagrams one usually associates with perturbative QCD. At the moment the Lund model is very useful, however, to become a truly interesting model theoretically it is quite important that a firm connection with the fundamental theory be made.

5.B. Monte Carlo Simulations in Hadronic Interactions

The grandfather of those models whose object is to generate realistic events in a hadron-hadron collision is ISAJET. Until recently such models treated hard scattering events well as far as, say, jet cross sections and the evolution of jets in the final state are concerned. There was also a phenomenological treatment of the soft particles present in such events. The progress made recently has been a proper inclusion of initial state radiation off the lines involved in the hard scattering. The various elements are schematically illustrated in Fig. 7.

Nevertheless, such simulations are, perhaps, still at the stage of being an art rather than a science. For example it is not known whether there is important interference between initial state radiation and beam fragments or between either of these with final state radiation. A realistic way of putting in multiple minijets is far from clear. The problems here are very important from a practical point of view, but at the moment they seem very hard as far as keeping systematic contact with QCD, except for those features depending only on large transverse momentum.
Coherence and the string effect

One of the important successes of the Lund model as opposed to independent fragmentation models was its prediction of the so-called "string-effect". The experimental situation is the following: consider a three jet event in a e^+e^- collision, viewed in the plane of three jets. Project the momenta of the associated particles onto that plane. Define the gluon jet to be the jet of smallest momentum. Then, experimentally, one observes an excess of associated particles between the gluon jet and either of the quark jets and a depletion of associated particles between the two quark jets. The excess or depletion is, for example, as compared to a prediction using independent fragmentation of the three jets. In the string model \(^1\) this comes about because the two fragmenting strings, each connecting one of the quarks to the gluon are moving away from the straight line connecting the quark and anti-quark.

Recently an incisive analysis of this effect has been given within QCD\(^2\). This analysis points out, at the same time, the great virtues of the Marchesini-Webber model and some of its limitations. To simplify the analysis, let us work at the level of the leading term in the \(1/N_c\) expansion\(^2\). Referring to Fig. 8, we may write the cross section for emission of a soft gluon of momentum \(k\) as

\[
\frac{d\sigma}{d\Omega_k} = \frac{\alpha_c}{4\pi^2} k \langle \gamma_j \rangle \cdot W
\]

with

\[
W = \frac{1 - \cos\theta_{1g}}{(1 - \cos\theta_1)(1 - \cos\theta_g)} + 1 \leftrightarrow 2.
\]

Now imagine fixing \(\theta\) and varying the azimuthal angle which the gluon, \(k\), makes with the jet plane. \(W\) achieves a maximum when \(k\) is in the plane of the jets and between the gluon and quark-1. \(W\) is at a minimum in the plane of the jets and between the two quarks. This is the string effect.

Now one may write (24) as

\[
2W = \left( \frac{1}{1 - \cos\theta_1} + \frac{\cos\theta_{1g} - \cos\theta_{1g}}{(1 - \cos\theta_1)(1 - \cos\theta_g)} \right) + \left( \frac{1}{1 - \cos\theta_g} + \frac{\cos\theta_{1g} - \cos\theta_{1g}}{(1 - \cos\theta_1)(1 - \cos\theta_g)} \right) + 1 \leftrightarrow 2.
\]
\[ 2W = \frac{1}{1-\cos \theta_1} \cos (\theta_1 - \theta_{1\sigma}) + \frac{1}{1-\cos \theta_2} \cos (\theta_2 - \theta_{2\sigma}) + \]
\[ + \frac{1}{1-\cos \theta_{g\sigma}} \cos (\theta_{1\sigma} - \theta_{g\sigma}) + \cos (\theta_{2\sigma} - \theta_{g\sigma}) \]  \tag{26}

(26) is not exactly the same as (25) but it is a reasonable average approximation. (26) is the expression used in the Webber Monte Carlo and gives an excellent fit to the data.

Thus the Marchesini-Webber model has built in enough of the coherence of QCD to give a good fit even to an effect where coherence is of the utmost importance. This is very nice. However, the model does not have the leading order QCD effects exactly built in for such angular correlations and that is a shame. Whether or not it is in principle possible to exactly build in such angular correlation is not known at present.
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Fig. 1

Fig. 2

Fig. 3

Fig. 4
Fig. 5

Fig. 6
Fig. 7

Fig. 8
ABSTRACT

We present results from a high statistics study of the nucleon structure function $F_2(x,Q^2)$ in the kinematic range $0.25 \leq x \leq 0.7$ and $50 \text{ GeV}^2 \leq Q^2 \leq 150 \text{ GeV}^2$. The analysis is based on $1.3 \times 10^6$ reconstructed events recorded at beam energies of 120, 200 and 280 GeV. By comparing data taken at different beam energies, we find $R = \sigma_L/\sigma_T = (8 \pm 14 \text{ (stat.)} \pm 3 \text{ (syst.)}) \times 10^{-3}$ independent of $x$ in the range $0.25 \leq x \leq 0.7$ and $50 \text{ GeV}^2 \leq Q^2 \leq 150 \text{ GeV}^2$. The kinematic range of these data makes them well suited for quantitative tests of Quantum Chromodynamics (QCD). From a next-to-leading order nonsinglet fit, we find a QCD mass scale parameter $\Lambda_{QCD} = 220 \pm 20 \text{ (stat.)}^{+90}_{-70} \text{ (syst.)} \text{ MeV.}$
We present new results on the nucleon structure function $F_2(x,Q^2)$ and $R = \sigma_L/\sigma_T$ measured in deep inelastic scattering of muons on an isoscalar carbon target. The data were collected with a high-luminosity spectrometer at the CERN SPS muon beam. The experimental apparatus is described in detail elsewhere [1]. Preliminary results obtained with the same set-up have been reported earlier [2].

The analysis presented here is based on $1.3 \times 10^6$ reconstructed events after kinematic cuts. Muon beams of 120, 200 and 280 GeV energy were used for this measurement. Beam polarities, kinematic ranges and data samples are summarized in Table 1. At present, the 280 GeV data represent only ⅓ of the total statistics recorded at this energy. In view of the high statistical accuracy of these data, a large effort was invested in calibrating the apparatus, and in monitoring its performance, in order to reduce systematic errors to a similar level. As the most important systematic limitation of the experiment is the energy calibration of the incident and scattered lepton, special emphasis was put on calibrating the magnetic field in the iron toroids, where it is not measurable directly. A map of the magnetic excitation $H$ was measured in the thin air gaps between individual discs of the iron toroids [1]. Inside the iron, it was converted into magnetic induction $B = \mu(H) \cdot H$ using accurately measured permeability curves for a large number of iron samples. The magnetic flux through the iron toroids and its dependence on the azimuth angle $\phi$ were verified with induction loops wound around various segments of the magnet. We estimate the uncertainty of the resulting field map to be smaller than $2 \times 10^{-3}$ over the entire magnet volume. The air gap magnet of the beam momentum spectrometer [1] was calibrated to an accuracy ranging from $1.5 \times 10^{-3}$ at 120 GeV beam energy to better than $1 \times 10^{-3}$ at 280 GeV.

**TABLE 1 : The Data Sample**

<table>
<thead>
<tr>
<th>Beam energy (GeV)</th>
<th>Beam signs</th>
<th>$Q^2$ range (GeV)</th>
<th>$x$ range</th>
<th>Number of events</th>
</tr>
</thead>
<tbody>
<tr>
<td>120</td>
<td>$\mu^+/$$\mu^-$</td>
<td>25-115</td>
<td>0.25-0.8</td>
<td>600 000</td>
</tr>
<tr>
<td>200</td>
<td>$\mu^+/$$\mu^-$</td>
<td>42-200</td>
<td>0.25-0.8</td>
<td>600 000</td>
</tr>
<tr>
<td>280</td>
<td>$\mu^+$ only</td>
<td>60-280</td>
<td>0.25-0.8</td>
<td>115 000</td>
</tr>
</tbody>
</table>
To calibrate the luminosity of the experiment, the incident beam of \( \approx 2.10^7 \) \( \mu \) sec intensity was counted with a fast plastic scintillator hodoscope using two different methods [1] which normally agree to \( \approx 0.5\% \).

The data were analyzed using a detailed Monte Carlo simulation of the experiment which takes into account

- the phase space of the incoming beam, including all correlation effects;
- efficiencies and resolution properties of all detectors in the apparatus;
- multiple scattering and energy loss of both incident and scattered muons [3], simulating the stochastic nature of energy losses due to ionization, bremsstrahlung, pair production and photonuclear effects;
- additional detector hits from hadronic shower punch-through close to the interaction vertex.

For each of the three beam energies, \( 3.5 \cdot 10^6 \) events were generated and processed through exactly the same chain of reconstruction programs as the experimental data. From the reconstructed Monte Carlo events, fine-grain acceptance matrices in \( Q^2 \) and \( x \) which include all effects from resolution smearing are calculated to convert the experimental distributions into deep inelastic cross sections. The acceptance is typically 75\% and is rather flat in the kinematic region \( Q^2/Q_{\text{max}}^2 > 0.2 \), \( x > 0.3 \).

To extract the one-photon exchange cross section from the measured data, corrections must be applied for higher order processes. The radiative corrections used in this analysis are described in detail in refs. [5] and include

- lepton current processes up to order \( \alpha^4 \),
- vacuum polarization by leptons and hadrons,
- hadron current processes up to order \( \alpha^3 \),
- effects of weak-electromagnetic (\( \gamma - Z^0 \)) interference [4].

They amount to at most 10\% over the kinematic range of this measurement. The error on \( F_2(x,Q^2) \) from uncertainties on these corrections is estimated to be smaller than 1\%.
The structure functions at the three beam energies, assuming $R = 0$, are shown in Fig. 1. The comparison of $F_2$'s measured at different beam energies allows to determine $R$ and, in addition, provides a powerful cross-check of systematic errors. A variation of $R$ affects mainly the region of large $y = \omega/E$, i.e. of low $x$ and high $Q^2$. In contrast, errors on the relative normalization of the data sets are independent of any kinematic variable and a scale error on the momentum measurement of incident or scattered muons affects mainly the region of small $y$ i.e. large $x$ and small $Q^2$. Effects from these three different sources are therefore only weakly correlated and can be studied separately.

Fig. 1: The nucleon structure function $F_2(x,Q^2)$ measured at the three beam energies 120 (circles), 200 (squares) and 280 GeV (triangles). The 120 GeV data were multiplied by a factor 1.025 to adjust the relative normalization of the three data sets. Only statistical errors are shown.
While 200 and 280 GeV data are found to be in very good agreement, the data points at 120 GeV are 2.5% lower everywhere in the kinematic region of overlap. Therefore, they were multiplied by 1.025 for the further analysis. The systematic errors on all results presented below do not depend on the absolute normalization of the data but only on the relative normalization of the three data sets with respect to each other. We assume normalization uncertainties of ± 1.5% of both the 120 and 280 GeV data relative to the 200 GeV data. We then study the mutual agreement of the three data sets under variation of the overall calibration of the magnetic field. This is shown in the form of a $\chi^2$ curve in Fig. 2 which exhibits a clear minimum at a recalibration factor $f_B = 1.0007 \pm 0.0008$, indicating that the calibration described above is indeed correct to $10^{-3}$. We chose to retain the original calibration but assign to the magnetic field an asymmetric error $\Delta B/B = +2 \cdot 10^{-3}$.  

![Fig. 2: The $\chi^2$ of the three $F_2$ data sets with respect to each other as a function of a change in the overall calibration of the spectrometer magnetic field. The abscissa $f_B$ is a global recalibration factor applied to the field map described in the text.](image)

$R = \sigma_L/\sigma_T$ is also determined by minimizing the $\chi^2$ of the three data sets with respect to each other. This is done separately in each bin of $x$ but assuming $R$ to be independent of $Q^2$ in the kinematic range of this analysis ($50 \text{ GeV}^2 \leq Q^2 \leq 150 \text{ GeV}^2$), as suggested by QCD calculations which predict only a weak (logarithmic) variation of $R$ with $Q^2$ [6]. The result is shown in Fig. 3 and is clearly compatible with $R = 0$. Good agreement is also observed with the measurement by the European Muon Collaboration (EMC) on an iron target [7]. From our data, we find a mean value of $R = [8\pm14(\text{stat.})^{+38}_{-35}(\text{syst.})] \cdot 10^3$ in the range $0.25 \leq x \leq 0.7$. We use $R = 0$ independent of $Q^2$ and $x$ for the further analysis.
Fig. 3: \( R = \sigma_L/\sigma_T \) as a function of \( x \). Also shown is the measurement by the EMC collaboration on an iron target [7]. Inner error bars are statistical errors only, outer error bars are statistical and systematic errors added linearly.

Fig. 4 shows the final \( F_2(x,Q^2) \) combined from the three data sets. To compare the observed deviations from Bjorken scaling to QCD predictions we show in Fig. 5 the logarithmic slopes \( d\ln F_2/d\ln Q^2 \), assumed to be independent of \( Q^2 \) and fitted to the data in bins of \( x \). Data points with \( y = v/E < 0.2 \) were not used in these fits to reduce the sensitivity to the magnetic field uncertainty. Also shown in Fig. 5 are QCD predictions for different values of the QCD mass scale parameter \( \Lambda \). They were obtained by a next-to-leading order computation in the \( \overline{\text{MS}} \) renormalization scheme, using the Altarelli-Parisi evolution equations [8] and the \( x \) dependence of \( F_2 \) at fixed values of \( Q^2 \) from the data shown in Fig. 4. Within errors, the data are compatible with the QCD predictions for \( \Lambda_{\overline{\text{MS}}} = 230 \text{ MeV} \). We consider this comparison the most stringent test of QCD in deep inelastic scattering because it does not require any parametrization of the structure function in \( x \) and \( Q^2 \). In the kinematic range of our data, it is also insensitive to assumptions on the gluon distribution.

We also did a flavour non-singlet, next-to-leading order QCD fit to our data, based on the method developed by Gonzales-Arroyo et al. [9]. This fit yields \( \Lambda_{\overline{\text{MS}}} = 220 \pm 20(\text{stat.})^{+90}_{-70}(\text{syst.}) \text{ MeV} \) for a \( \chi^2/\text{DOF} = 190/165 \) and is superimposed to the data in Fig. 4. Using other QCD fit programs or applying
The structure function $F_2(x,Q^2)$ combined for all beam energies assuming $R = 0$. Only statistical errors are shown. The solid lines represent the QCD fit described in the text.

more restrictive cuts in $x$ or $y$ changes $A$ by less than 20 MeV; such effects are included in the estimate of the systematic error. Singlet fits under reasonable assumptions on the gluon distributions also give very similar values for $A$. We stress that these results are obtained in a kinematic region ($Q^2 \geq 25$ GeV$^2$) which is generally believed to be well described by perturbative QCD predictions and is not obscured by collective ("higher twist") effects from quark-quark interactions [10].

In conclusion, we have presented a new high statistics measurement of the nucleon structure function $F_2(x,Q^2)$ from deep inelastic muon-carbon scattering in the high $Q^2$ ($Q^2 \geq 25$ GeV$^2$) regime. Careful calibration of the experimental apparatus has allowed to reduce systematic uncertainties to a level close to the statistical accuracy of the data. $R = \sigma_L/\sigma_T$ is found to be compatible, within small
Fig. 5: Logarithmic slopes of the structure function $d\ln F_2 / d\ln Q^2$ as a function of $x$, compared to QCD predictions for different values of $\Lambda$. Inner error bars are statistical errors only, outer error bars are statistical and systematic errors added linearly.

errors, with $0$ in the kinematic range $0.25 \leq x \leq 0.7$. The pattern of scaling violations observed in the data is in good agreement with predictions from perturbative QCD for a mass scale parameter $\Lambda_{\text{MS}} = 230$ MeV.

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[9] A. Gonzales-Arroyo, C. Lopez and F.J. Yndurain,
Abstract:
A trustworthy determination of the nucleon structure functions is required in order to estimate production rates in future hadron colliders at very high energies but also to make predictions for present experiments.
The predictions of QCD for the evolution of the unpolarized nucleon structure functions have been exhaustively studied. Although a great deal of theoretical work has been devoted to spin physics at short distances, the polarized nucleon parton distributions are not well understood, owing in part to the scarcity of experimental data in a region of phase space where perturbative QCD can be confidently applied. In a first part, giving these structure functions at low energies, we show perturbative QCD predictions for their $Q^2$ behaviour up to high energies. In a second part, we calculate some spin-dependent observables which are dominated by short distance physics. They represent valuable tests which should be confronted with experimental data. The EMC collaboration will soon allow us to compare their data with theoretical expectations for muoproduction using a polarized target.

1. Perturbative QCD predictions for the $Q^2$ evolution up to super high energy.

We present a new set of spin-dependent structure functions [1] for quarks and gluons inside a polarized nucleon. This is the only available set which satisfies all the sum rules and analytic requirements and is valid in the kinematic range: $10^{-5} \leq x \leq 0.9$ and $5 \leq Q^2 \leq 5 \times 10^8$ GeV$^2$, the effects of heavier flavours being taken into account. The method, for solving Altarelli–Parisi equations, we are using here is based on the moments of the polarized parton distributions [2]. The choice of the low energy input is not unique [3] due to some uncertainty in the experimental data, but this does not seriously affect the predictions at high energies for the gluon and valence quark polarizations. This is not the case for the sea quark polarization, which can be either positive or negative. Anyway, this polarization is of the order of $10\%$, and the sea contribution is small and remains largely below the other parton distributions. We show in fig.1 the
gluon polarization $\Delta G(x)/G(x)$ for different $Q^2$ values. Here $\Delta G(x) = G^+(x) - G^-(x)$ where $G^\pm (x)$ denote the parton distributions in a polarized nucleon either with helicity parallel (+) or antiparallel (−) to the nucleon helicity. We see that at low energy this polarization is less than 20% but it rises for all $x$ with energy. Such a spectacular effect which was observed previously and widely discussed before [3,4] is even more striking with this new determination because of a very rapid rise, mainly in the low $x$ region. Note that the valence quark polarization also grows with $Q^2$ but does not change drastically from the low energy input. (see ref.1)

![Fig. 1 - Gluon polarization $\Delta G(x)/G(x)$ as a function of $x$ for $Q^2 = 5$ GeV$^2$ (dashed -dotted curve), $Q^2 = 5 \times 10^4$ GeV$^2$ (dashed curve) and $Q^2 = 5 \times 10^8$ GeV$^2$ (solid curve).]

2 – Predictions for EMC experiment

Spin asymmetries in deep inelastic scattering and semi-inclusive deep inelastic scattering with longitudinally polarized electron and proton beams are examined and will be compared soon with the EMC collaboration data. We will consider successively totally inclusive asymmetry and semi-inclusive asymmetry.
- For totally inclusive asymmetry, we define:

\[ A_1(x) = \frac{\sum_q e_q^2 \Delta q(x)}{\sum_q e_q^2 q(x)}. \]

On fig. 2 we compare the different predictions for \( A_1^p \). A detailed study of scaling violations for spin-dependent structure functions [3] has been worked out using partonic structure functions, constrained by the two Bjorken sum rules. That leads to some differences in the asymmetry as explained in the caption.

- For semi-inclusive asymmetry, we define:

\[ A_1^{H/}(x_H, y, Z_H) = \frac{\sum_i e_i^2 \Delta q_i^H(x_H) D_i^{H/}(Z_H \gamma[1-(1-y)^2])}{\sum_i e_i^2 q_i^H(x_H) D_i^{H/}(Z_H \gamma[1+(1-y)^2])} \]

where \( H \) and \( H' \) stand for a proton and a positive or negative pion. The next to leading logarithmic corrections to cross sections for semi-inclusive deep inelastic scattering from a polarized target have been worked out [5], this effect is small in magnitude and therefore will be difficult to establish. We give in fig. 3 the asymmetries \( A_1^x \) for positive and negative pions from a proton target.

Fig. 2 - Various predictions at \( Q^2 = 5 \text{ GeV}^2 \) for the asymmetry parameter \( A_1^p \) on a proton target as a function of \( x \). Full curve: classical Kaur-Carlitz model. Dashed curve: model with negative sea polarization. Dotted dashed curve: model assuming that spin distribution for \( u \) and \( d \) valence quarks are differently modified by the presence of the gluon.
Fig. 3 - $x_H$ dependence of the pion asymmetry at $Q^2 = 20$ GeV$^2$ from a proton target for several $z_T$ values. Full curves: leading predictions for positive pions. Dashed curves: next to leading corrections. Dotted-dashed curves: leading predictions for negative pions. The shaded area gives the uncertainties coming from the parton densities.

References:
HIGH-MASS DIMUON PRODUCTION IN $\pi^- W$ AND $\pi^- D$ INTERACTIONS

(NA10 Collaboration)
presented by C. Vallée and M. Winter

ABSTRACT: We present recent results on the production of continuum $\mu^+\mu^-$ pairs with masses higher than 4 GeV/c$^2$ in $\pi^- W$ and $\pi^- D$ interactions, obtained with the NA10 spectrometer. The valence quark distribution for the $\pi^-$ is determined with the most up to date theoretical tools, a direct test of scaling violation is provided, and differential nuclear effects in the Drell-Yan process are investigated.
INTRODUCTION

The high statistics (see Table 1) accumulated with the NA10 spectrometer\textsuperscript{1}) from 1981 through 1985 provide the basis for a significant improvement in our understanding of the Drell-Yan process.

<table>
<thead>
<tr>
<th>Period</th>
<th>$\pi^-$ momentum (GeV/c)</th>
<th>Number of W events</th>
<th>Number of D events</th>
</tr>
</thead>
<tbody>
<tr>
<td>1981-1982</td>
<td>194</td>
<td>150,000</td>
<td>0</td>
</tr>
<tr>
<td>1983+1985</td>
<td>286</td>
<td>84,400</td>
<td>9,200</td>
</tr>
<tr>
<td>1984</td>
<td>140</td>
<td>45,900</td>
<td>4,700</td>
</tr>
</tbody>
</table>

Data were taken at 140, 194 and 286 GeV/c, on two different targets (D and W) exposed simultaneously.

In section 1 of this report we compare the scale-invariant differential cross sections $\pi^- 3/2 d\sigma/d\sqrt{\tau}$ measured on W at 194 GeV/c and 286 GeV/c with the aim of testing scaling violation in a theory-independent fashion. We also compare the corresponding doubly-differential cross sections $d^2\sigma/d\sqrt{\tau} dx_F$ with the predictions of QCD in the Soft-Gluon-Emission (SGE) approximation\textsuperscript{2}), and extract within this formalism the distribution of the $\pi^-$ valence quarks.

In section 2 we compare the cross sections measured at 140 GeV/c and 286 GeV/c on D and W targets in order to explore possible differential nuclear effects.

We note that our data on W at 194 GeV/c have already been published\textsuperscript{3)} and analysed\textsuperscript{4}), although using different theoretical approximations.

SECTION 1

a) A very direct way of testing scaling violation in lepton-pair production is to measure the invariant cross section $M^2 d\sigma/dM (= \pi^- 3/2 d\sigma/d\sqrt{\tau})$ as a function of the variable $\sqrt{\tau}$ (where $\tau = M^2/s$) at different values of the center-of-mass energy squared $s$. This cross section should be independent of $s$ only if scaling holds.

While earlier experiments\textsuperscript{5,6,7)} had failed, mostly because of poor statistics, to exhibit scaling violation in this way, our data allow a more severe test. Figure 1 represents the ratio:

$$R_{12} = \frac{s_1 d\sigma/d\sqrt{\tau}(s_1)}{s_2 d\sigma/d\sqrt{\tau}(s_2)}$$

$s_1 \approx 539$ Gev² ; $s_2 \approx 366$ Gev² ;
with $x_F$ integrated at both energies from -0.1 to 0.5. The full horizontal line is the one about which the ratio of data points should scatter if scaling held (i.e. $R_{12}^{1.2} = 1.0$) and if there were no uncertainties in the absolute normalizations at the two energies; the latter introduce into the ratios an uncertainty of ±8%. The dashed horizontal line is fitted to the data points with $0.24 < \sqrt{\tau} < 0.36$. The dashed curves represent the ratios predicted theoretically in the SGE approximation \(^2\), for two different values of the QCD parameter $\Lambda$ (i.e. 100 and 400 MeV/c).

We have estimated the compatibility of the observed ratio with the scaling hypothesis by computing the probability that the four points with $\sqrt{\tau} > 0.54$ represent statistical deviations from the dashed horizontal line; we find approximately 4%. On the other hand, the experimental ratio is clearly compatible with the QCD prediction for either value of $\Lambda$, especially when allowing for the normalization uncertainties mentioned before (~8%). Thus our data lend strong support to the hypothesis of scaling violation.

b) We next proceed to comparing the measured doubly-differential cross section $d^2\sigma/d\tau dxF$ with the one predicted in QCD, and to determining the distribution of valence quarks in the pion.

Similar analyses have already been performed (see for instance Ref.4 and references quoted therein), but the theoretical cross sections used were at most computed to the first perturbative order in $\alpha_s$, namely in the Next-To-Leading-Logarithm approximation (NLLA). Higher-order corrections were neglected, which are by now known to be significant\(^2\); it is hence not surprising that those analyses found the theoretical cross sections significantly lower than the observed ones. The two cross sections differed mainly by a normalization factor (generally denoted by $K$) which varied between 1.3 and 1.6 in the NLLA, depending on the nucleon structure functions adopted\(^4\). A further uncertainty in $K$, of ±25%, came from the lack of precision of the pion sea and gluon parameters used as inputs. As was already quoted in Ref.4, one can, by adopting a certain nucleon parametrization and by choosing appropriately (but well within their quoted errors) the pion parameters just mentioned, obtain $K = 1.03 \pm 0.03$ (stat.).

The theoretical cross section used here goes beyond the NLLA by including the the contribution of the infrared singularities of most of the higher-order diagrams, i.e the so-called "Soft-Gluon-Emission"\(^2\). In comparison with the NLLA, the SGE approximation modifies both the partonic cross section of the elementary process, as well as the evolution of the quark distributions. In our kinematical domain, the largest difference between the cross section computed in the SGE approximation and the one computed in the NLLA occurs at small values of $\sqrt{\tau}$, where it may reach 15 to 30%, depending on the value of $\Lambda$. As an another improvement we now parametrise the distributions of the valence quarks according to Ref.11 instead of the one used in all previous analyses\(^{10}\); it reads as follows:

$$q_\nu(x,M_\perp^2) = \left[r(x^-x^+) + A \alpha_s^{-d\nu}(M_\perp^2) x^\mu \right](1-x)^{\beta(M_\perp^2)} ,$$
where $A$ and $\beta(M^2_o)$ have to be determined from the data. All other parameters are fixed by Ref. 11, while $M^2_o$ is arbitrary; we have chosen $M^2_o = 25 \text{ GeV}^2/c^4$. This parametrization is presumed to be more accurate when $x$ tends to zero. In addition the $Q^2$ dependence is more easily introduced.

In fitting the theory to the data, the parameters $A$ and $\beta(M^2_o)$, as well as the overall normalization $(K)$ were left free to vary. We have derived these parameters, for the time being, exclusively from a fit to the 194 GeV/c data.

Allowing for the stated uncertainties on the nucleon valence and pion sea quark as well as gluon distributions, we find the following ranges of values for these parameters:

$$0.4 \leq A \leq 0.8 ; 0.9 \leq \beta(M^2_o) \leq 1.2 ; 0.9 \leq K \leq 1.2$$

The central value of $K$ is now quite close to 1.0 and this – in contrast to our earlier analysis – without having "tuned" any input parameters; this is mainly due to the inclusion of the SGE corrections.

In Figure 2 we compare the experimental cross section $d\sigma/dx_F^N$ in various $\sqrt{s}$ bins with theory. The agreement is good for $0.24 < \sqrt{s} < 0.42$ (i.e. $x^2/d.o.f. = 1$); at $\sqrt{s} > 0.54$ the agreement is poor (i.e. $x^2/d.o.f. = 2$), i.e. the anomaly reported in Ref.4 persists.

Figure 3 shows an analogous comparison for the 286 GeV/c data using for the theory the parameters extracted from the 194 GeV/c fit. The agreement is good for $0.27 < \sqrt{s} < 0.36$ (i.e. $x^2/d.o.f. = 1$); the disagreement at high masses (i.e. for $\sqrt{s} > 0.45$) can be taken as a confirmation of the "anomaly" at a different c.m. energy, while the discrepancy for $\sqrt{s} < 0.27$ might originate from reinteractions or heavy flavour decay.

In summary, fits with the improved theoretical tools yield good agreement in shape and in magnitude (i.e. $K \approx 1$) over most of the kinematical range, but still fail to eliminate the "anomaly" reported by us. The cause of this discrepancy at high masses may lie in nuclear effects, in as yet unavailable QCD corrections, in the nucleon valence distributions provided by D.I.S. or in all these sources jointly.

**SECTION 2 : DEUTERIUM - TUNGSTEN COMPARISON**

In this section, we present preliminary results on the comparison of the differential cross sections on W and D.

Since the linear dependence of the total cross section on the atomic mass $A$ has been established to quite high accuracy, we shall compare only the shapes of the two distributions, normalizing the two samples to the same number of events.

In order to investigate possible nuclear effects in the structure functions, we compare the distributions from the two targets as functions of $x_1$, $x_2$ on the one hand, and of $\sqrt{s}$ and $x_F$ on the other (Figure 4), where $x_1$ (resp. $x_2$) stands for the momentum fraction of the projectile quark (resp. target quark). Since we obtained similar results at 140 GeV/c
as well as at 286 GeV/c, we show in Figure 4 the comparison for the combined data. There is an indication of a depletion of the W data at $\sqrt{\tau} > 0.45$, with a corresponding trend in $x_2$. The correlation of this effect with the ones observed by the EMC and SLAC collaborations\textsuperscript{14, 15} as well as with the "anomaly" reported above is under study.

We have also investigated, at both c.m. energies, the ratio of the dimuon transverse momentum ($p_T$) distributions for D and W events; Figure 5 shows a larger mean transverse momentum squared in W, which can be expressed as:

$$<p_T^2>_W - <p_T^2>_D = \begin{cases} 
0.19 \pm 0.03 \text{ GeV}^2/c^2 \text{ (286 GeV/c)} \\
0.18 \pm 0.03 \text{ GeV}^2/c^2 \text{ (140 GeV/c)}
\end{cases}$$

This difference is probably due to the scattering of the incoming quark within the nucleus in which the dimuon is produced; our data are in qualitative agreement with some predictions made under this assumption\textsuperscript{16, 17}. It is certainly not an experimental effect, due to the scattering of the incident pion in the 12 cm long Tungsten target, since data taken with a shorter target (5.6 cm) showed the same effect.

CONCLUSION

We have measured the differential cross section for the production of high-mass $\mu^+\mu^-$ pairs in $\pi^-W$ and $\pi^-D$ interactions at 140 GeV/c, 194 GeV/c and 286 GeV/c.

We observe that the dependence of the "invariant" cross section $s \tau^{3/2} d\sigma/d\sqrt{\tau}$ with respect to $\sqrt{\tau}$ changes with the cm energy. The observed change disfavors the scale invariance hypothesis at the 95% level. In contrast to all previous tests of scaling violation in dilepton production, the present one is independent of any specific model.

We have furthermore compared the differential cross sections $d^2\sigma/d\sqrt{\tau}dx_F$ measured on W at 194 GeV/c and 286 GeV/c with the one predicted by QCD in the SGE approximation, the most complete theoretical cross section presently available. We have thus determined the $\pi^-$ valence quark distribution, parametrized according to Ref. 11. Except for some discrepancies (in particular the "anomaly" already reported\textsuperscript{4}) at 194 GeV/c), the agreement between theory and experiment is good, in shape as well as in magnitude (i.e $K \approx 1$).

Finally, we have performed a preliminary comparison of the differential cross sections measured on D and W, and this at 140 GeV/c as well as at 286 GeV/c. This comparison shows that the mean transverse momentum squared observed with W is by about $0.19\pm0.03 \text{ GeV}^2/c^2$ higher than the one obtained with D; it also indicates a nuclear effect, similar to the "EMC effect", in the $x_2$ and $\sqrt{\tau}$ distributions. The difference in the mean transverse momentum may be due to multiple scattering of the incoming parton in the nucleus where the dimuon is produced.

We also presented some results on the angular distributions of the muons. These will appear in print shortly\textsuperscript{18}. 
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18) S. Falciano et al. (NA10 Collaboration), to appear in Z. Phys. C.
\( \sqrt{s_1} = 23.2 \text{ GeV}, \; \sqrt{s_2} = 19.1 \text{ GeV} \)

\( \Lambda = 100 \text{ Mev} \)

\( \Lambda = 400 \text{ Mev} \)

Fig 1.
Fig. 2
NA10

Fig. 3
FIGURE 4

\[ \frac{\sigma (\pi^- \rightarrow \mu^+ \mu^-)}{\sigma (\pi^- \rightarrow D^- \mu^+ \mu^-)} \text{ (SHAPE ONLY)} \]

DATA 286 GEV + 140 GEV

FIGURE 5

\[ \frac{\sigma (\pi^- \rightarrow \mu^+ \mu^-)}{\sigma (\pi^- \rightarrow D^- \mu^+ \mu^-)} \text{ as a function of the Dimuon transverse momentum (shape only)} \]
We have analysed the QCD evolution of the structure function $F_2$ of iron and hydrogen and we have obtained that the momentum distribution of valence quarks is lower in iron than in hydrogen and that the opposite is true for the glue and sea quark distribution in the region of $0.1 < x < 0.65$. 
We present the results of the QCD analysis of the evolution of the structure function $F_2$ of hydrogen and iron measured in the deep inelastic muon scattering by the European Muon Collaboration EMC. We have previously observed that the structure function $F_2$ in iron and deuterium are different\(^1\). Main features of this observation, called the EMC effect, have been confirmed by different experiments\(^2\). In this paper we will extract the difference in distribution of valence quarks, sea quarks and glue in iron and in nucleon. However one should keep in mind that in deep inelastic scattering only structure function $F_2$ is measured directly so the results of measurements are insensitive to the flavour of quarks, contrary to neutrino scattering. We will show in this paper that the difference in measured structure functions $F_2$ of iron and nucleon analysed in terms of QCD implies a different distributions of valence and sea quarks and glue in heavy nuclei and in nucleon. Assumptions used in this analysis will be listed below.

In the quark-parton model $F_2$ is the sum of quark distributions weighted by its charge squared. Assuming that the nucleon is composed of $u$ and $d$ valence quarks $u_v$ and $d_v$ and the sea quarks $u, d, s, c, u, d, s, c$ we can write the following expression for the structure function $F_2$ of hydrogen:

$$F_2^{H}(x,Q^2) = \frac{5}{18}q_s + \frac{3}{18}q_{ns}$$

where the singlet part of the structure function is given by the expression $q_s = u_v + d_v + 2(u - d + s + c) = u_v + d_v + \text{sea}$ and the nonsinglet one by $q_{ns} = u_v - d_v$. Writing these expressions we have neglected the differences $u(x,Q^2) - g(x,Q^2)$ and $g(x,Q^2) - s(x,Q^2)$.

For iron we have assumed that the structure function $F_2$ of nucleus is the sum of the structure functions of all the nucleons in the nucleus normalised to one nucleon. This is a strong assumption since in some models of the EMC effect a nucleus is composed of deltas, pions .... etc. in addition to nucleons. We have also assumed that the $u_v$ and $d_v$ distributions in proton are equal to $d_v$ and $u_v$ distributions in neutron, respectively. The following formulae can thus be obtained for the structure function $F_2$ of iron:

$$F_2^{Fe}(x,Q^2) = \frac{5}{18}q_s - \frac{1}{14}(3/18)q_{ns}$$

Here the singlet and nonsinglet part of the structure function $q_s$ and $q_{ns}$ have exactly the same definition for the nucleons in iron as $q_s$ and $q_{ns}$ for the proton in hydrogen.

It can be seen from these expressions that the singlet part of the structure function depends on the valence and sea quark distribution while the nonsinglet part depends on the difference between valence $u$ and valence $d$ quark distributions. The evolution of the singlet part of the structure
function with $Q^2$ is coupled to the evolution of glue. Hence the glue
distribution can be extracted from the variation of the singlet part with $Q^2$
Here $Q^2$ is the negative four momentum squared of virtual photon transfered
between scattered muon and a quark.

The charged lepton scattering is insensitive to the flavour of quarks
and it is not straightforward to distinguish valence and see quarks. We have
been able to fit separately valence and see quark distributions by aplying
the sum rules for the valence quarks and thanks to very different functional
form of valence and see quark distributions.

There is a strong correlation between the determination of the $\Lambda_{\text{QCD}}$
and the glue distribution. In this analysis the $\Lambda_{\text{QCD}}$ was fixed. It has been
obtained previously by fitting the evolution of nonsinglet part of the
structure function of hydrogen for the fractional quark momenta greater
than 0.35. Therefore such determination of $\Lambda_{\text{QCD}}$ was independent of the
see quark and glue distributions and we have verified that the glue and the
sea quark distributions are steep enought to justify this determination.

The data have been obtained by the European Muon Collaboration at
CERN between 1978 and 1980. The hydrogen target data sample has been
obtained with the incoming muon momenta of 120, 200, 240 and 280 GeV
and the iron data with the muon momenta of 120, 200, 250 and 280 GeV\(^3\).
The total statistics of the hydrogen and iron data was 410 000 and 984 000
events, respectively.

Assuming that the QCD evolution of the structure function $F_2$ is
described by the Altarelli-Parisi equations\(^4\) we have used its particular
solution by Gonzales\(^5\) in which the singlet and nonsinglet part of the
structure function was calculated to the second order in QCD. With fixed
$\Lambda_{\text{QCD}}$ we extract the distribution of quarks and gluons at a given $Q_0^2$ by
fitting, with the program MINUIT, the measured $F_2$ for all values of $x$ and $Q^2$
accessible in our experiment. The following parametrisation has been used
in the fit:

$$
\begin{align*}
x(u_v+d_v) &= A_1 x^{a_1} (1-x)^{b_1} \\
xq_{\text{sea}} &= A_2 (1-x)^{b_2} \\
x(u_v-d_v) &= A_3 x^{a_3} (1-x)^{b_3} \\
xG_{\text{glue}} &= A_4 (1-cx)(1/x)^{a_4}(1-x)^{b_4}
\end{align*}
$$

The integral of $(u_v+d_v)$ and $(u_v-d_v)$ distributions have been
normalised to 3 and 1 respectively accordind to quark sum rules. We have
checked that we could released this constraint in the final stage of fitting
without a significant change of results but the curves presented here have
been obtained with this constrain. In addition the sum of the fractional 
momenta of valence and sea quarks and glue was normalized to 1.

The data of hydrogen and iron have been fitted with the fixed value 
of $\Lambda_{\text{ms}} = 110 \text{ MeV}^3$ corresponding to our best fit value of $\Lambda$ obtained from 
the analysis of nonsinglet part of the structure function of hydrogen. We 
have estimated the systematical errors in two ways. First, we have 
repeated our fit with two extreme values of $\Lambda_{\text{ms}} = 50$ and 200 MeV allowed 
by systematical errors. Secondly, we have varied our data to the extreme 
limits allowed by their systematical errors and fitted them with the $\Lambda_{\text{ms}} = 
110 \text{ MeV}$. We do not consider the systematical errors of data and $\Lambda$ as 
independent since the quoted systematical errors of $\Lambda$ reflect the 
systematical errors of data used in this analysis.

The $R = \sigma_L / \sigma_T$ has been calculated according to QCD and has 
been taken into account in the fits. We have tested that the systematical 
errors of $R$ are much smaller than those of $\Lambda$.

The typical $\chi^2$ of the fits is 500 for 205 degrees of freedom for hydrogen and 300 for 180 degrees of freedom for iron. We do not worry 
about this large $\chi^2$ because the errors used in the fit do not include 
systematical errors and our statistical errors are very small thanks to 
large statistics of our data.

In the Gonzales program all the distributions are parametrised for 
the low $Q^2$ and then they are evoluted for higher values of $Q^2$. Therefore we 
have obtained the parametrisation of valence and sea quarks and glue for $Q^2 = 4 \text{ GeV}^2$ as the result of fit. However the data were fitted for the all $Q^2$ 
range accessible by our experiment. Knowing exactly the QCD evolution of 
the fitted distributions of valence and sea quarks with $Q^2$ we will present 
in the following their distributions for the mean $Q^2$ value of our experiment 
of 22.5 $\text{ GeV}^2$, however the distribution of gluons is presented for $Q^2 = 4 
\text{ GeV}^2$.

The momentum distributions for the valence quarks $u_v + d_v$, the sea 
quarks and the glue in iron and hydrogen are presented in Fig. 1. The solid 
line presents the results of fit with $\Lambda_{\text{ms}} = 110 \text{ MeV}$ while the dashed line 
presents the results with values of $\Lambda_{\text{ms}} = 50$ and 200 MeV. The dotted lines 
show the variation of results with the systematical errors of data. The 
results below $x=0.1$ and above $x=0.65$ are not plotted since the experimental 
results in this region do not give significant enough constrains for the fit.

It can be seen from the Fig. 1 that the momentum distribution of 
valence quarks are lower for iron than for hydrogen. The opposite is true for 
the glue and sea distributions particularly at large $x$. Thus the gluons and
Fig. 1.
sea quarks are more energetic in iron than in hydrogen contrary to the valence quarks. It can be noted however that the results for the valence and sea quarks are much more significant than for glue.

We have verified that the fit is stable against reasonable changes in the functional form of the fitted distributions. In particular we have fitted our data with the functional form of the sum and the difference of the $u_V$ and the $d_V$ distributions multiplied by $(1-xy)$. We have also fitted separately $u_V$ and $d_V$ distributions. Finally we have modified the singlet and sea distribution to allow them to extend above $x = 1$. Results of these test fits have not changed our conclusions.

We have tried to check if a possible strong correlation between parameters cannot create observed differences in our results for iron and hydrogen. Also we wanted to verify if another minimum of fit cannot be found which could change our conclusions. We have fixed the hydrogen glue in the fit of iron data. The $\chi^2$ divided by the number of degrees of freedom ($\Delta\chi^2/ndf$) has changed by 79% ($\Delta\chi^2/\chi^2 = 47\%$). Fixing the hydrogen sea in the fit of iron data has changed $\Delta\chi^2/ndf$ by 59% ($\Delta\chi^2/\chi^2 = 33\%$). When we have fixed the iron glue and sea in the fit of hydrogen data the $\Delta\chi^2/ndf$ has changed by 27% and 18% respectively, while the $\Delta\chi^2/\chi^2$ has changed by 11% and 7%.

In order to test further dependence of our results on different presentation of data we have made our fits with data in which different incoming muon momenta have been merged for $R_{QCD}=0.3$. The conclusions have not changed. The $\chi^2/ndf$ of the fit for hydrogen data was about 1 while for the iron target it was about 1.9.

We have presented above a new analysis of old measurement of the structure function $F_2$. We have analysed the QCD evolution of $F_2$ and we have obtained different distributions of valence and sea quarks and gluons in iron and hydrogen in the region of $0.1 < x < 0.65$. In order to obtain these results we have assumed that the QCD describes the $Q^2$ variation of the structure function $F_2$ in hydrogen and iron and that the $\Lambda_{QCD}$ is the same in iron and hydrogen. We have also assumed that the structure function $F_2$ of iron is the sum of structure functions $F_2$ of all nucleons in the nucleus. Finally we have used the sum rules for the valence quarks in the fit.

The structure functions $F_2$ of hydrogen and iron were the only experimental input to the fit. They were fitted independently and exactly in the same way. The different distribution of valence and sea quarks and glue in iron and hydrogen obtained in this paper gives the QCD interpretation of
the differences in the distribution of $F_2$. A reader should use with caution results presented in this paper if his favored interpretation of the EMC effect disagrees with the assumptions used in this paper.

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In an experiment performed at Fermilab we explore the structure of the pion via the Drell-Yan process, in the reactions $\pi^\pm N \rightarrow \mu^+\mu^- X$ at 80 and 255 GeV/c. The emphasis is on the kinematic region in which the momentum of the pion is largely carried by a single quark, which proves to be rather accessible to experimental study. The data indicate substantial departures from the expectations of asymptotic QCD, but agree well with a calculation which includes the effect of gluon exchange among the quarks inside the pion. With our large sample of muon pairs, we also set the current-best limits on the decay $D^o \rightarrow \mu^+\mu^-$ and on $D^o-\bar{D}^o$ mixing.
Our knowledge of the structure of hadrons is largely empirical, and is derived in most part from deep-inelastic lepton-hadron scattering experiments conducted at SLAC, Fermilab, and CERN. An interpretation of these results has been provided by theoreticians, beginning with the parton model of Bjorken and Feynman, and later the theory of QCD. However, these theories make no predictions of the detailed structure of matter, as they can give (at present) little understanding into the confinement of quarks and gluons in observable hadrons, and provide instead only a sense of the behavior when the fundamental constituents are 'asymptotically free.' Fermilab experiment E 615 extends the exploration of hadron structure by pursuing the limit in which a single quark carries nearly all of the momentum of a π meson. Significant departures are observed from the more 'standard' QCD expectations for this experiment, but which appear to agree well with a QCD calculation of pion structure that includes non-asymptotic effects of the binding of the quarks inside the pion.

Experiment E 615 is the successor to Fermilab experiments E 331¹ and E 444², all of which studied the reaction πN → μ⁺μ⁻X. The basic view of this reaction is that an antiquark from the pion annihilates with a quark from the nucleon to produce a virtual photon which materializes as the muon pair in the laboratory. The method for relating measurements of this process to the structure of the incident hadrons was developed by Drell and Yan³ following the pioneering experiment on pN → μ⁺μ⁻X by a group headed by Lederman⁴ in the late 1960's. The virtual photon 'x-rays' the beam- as well as the target-particle, so that any long-lived hadron may be studied.

Fermilab experiment E 444² was the first to use this technique to probe the pion, and to extract the pion structure function. The momentum distribution of valence quarks in the pion was found to be roughly \( \sqrt{x(1-x)} \), where \( x \) is the fraction of the pion's momentum carried by the quark. This distribution indicates it is much more probable for a single quark to carry most of the momentum of a pion than of a proton. Hence the use of a pion beam permits the study of elementary processes in the interesting limit that a laboratory particle is nearly equivalent to a quark. Since the quantum numbers of quarks and pions are different the equivalence cannot be complete, and there may emerge unusual features not encountered in other investigations of the strong interaction.

Fermilab experiment E 615 was designed to emphasize the forward production of muon pairs by pions. The detector, shown in Figure 1, was built by a collaboration of physicists from the University of Chicago, Iowa State University, and Princeton University.⁵ A large-aperture magnetic spectrometer was preceded by the 'selection' magnet which served to focus
high-mass muon pairs into the spectrometer, and to absorb the unscattered beam and secondary hadrons in low-Z material which filled its gap. The selection magnet was constructed out of steel and copper recycled from the main-ring magnets of the Argonne Zero Gradient Synchrotron. The experiment took data during the first running period of the Tevatron in 1983-84, and utilized an 80-GeV pion beam derived from 400-GeV protons, and a 255-GeV beam from 800-GeV protons. For a detailed discussion of the apparatus, see ref. 6.

Results from the 80-GeV run\(^7\) (and from the 255-GeV test run\(^8\)) confirm the expectation of interesting physics when a pion is almost a quark. An important piece of evidence is related to the angular distribution of the muon pairs, as shown in Figure 2. Muon pairs produced by the one-photon annihilation of free quarks should follow the familiar \(1 + \cos^2 \theta\) distribution in the pair rest frame. However the experimental result is that the angular distribution approaches \(\sin^2 \theta\) when the antiquark in the pion has a large momentum fraction. This is a model-independent indication that the virtual photon had longitudinal polarization, and suggests that the antiquark is off mass-shell due to its containment inside the pion. The new evidence confirms a preliminary indication of the same effect in experiment E 444\(^9\), but which was not found in the CERN experiment NA 3\(^10\).

Another striking result is the behavior of the average transverse momentum of muon pairs produced by large-x antiquarks in the pion. While the average squared transverse momentum is about \(1 \text{ GeV}^2/c^2\) at low and moderate \(x\), at large \(x\) the average dips sharply to \(0.6 \text{ GeV}^2/c^2\). The relatively large average transverse momentum observed in muon-pair production by hadrons is often taken as a sign of the effects of gluon bremsstrahlung. Now there is evidence that the 'standard' gluon corrections become less relevant in the same kinematic region that the muon-pair angular distribution changes markedly.

A third noteworthy feature concerns a detail of the pion structure function. An intriguing question which requires considerable experimental sensitivity is whether the probability distribution of the valence quark extrapolates to a finite intercept at \(x = 1\). The analysis of the 80-GeV run of E 615 indicates that this is so, although the statistical and systematic significance of this result is only three standard deviations.

All three of the above experimental results are consistent with the expectations of a QCD calculation by Berger and Brodsky\(^11\). This calculation is somewhat non-standard in that it takes explicit account of the fact that the annihilating quark and antiquark are contained in a nucleon and a pion. As a bonus they are led to an actual prediction of the form of the pion structure function applicable to the limit that the antiquark carries most of the momentum.
Figure 1. View of the apparatus of E 615.

Figure 2. The parameter $\lambda$ as a function of $x_\pi$ obtained from fits to the angular distribution of muons in the muon-pair rest frame, using the form $\frac{d\sigma}{d\cos \theta} \propto 1 + \lambda \cos^2 \theta$. The solid curve is based of the QCD model of Berger and Brodsky\textsuperscript{11} using the value $\langle P_T^2 \rangle = 0.62 \text{ GeV}^2/\text{c}^2$ deduced from the observed pion structure function.
of the pion. In addition to a piece of the pion structure function associated with a $1 + \cos^2 \theta$ angular distribution of the muon pair, they find a piece which has a finite intercept at $x = 1$ and is associated with a $\sin^2 \theta$-angular distribution.

The intercept of the pion structure function at $x = 1$ measured in E 615, when interpreted according to the model of Berger and Brodsky, leads to a prediction that the average squared transverse momentum of the muon pair should be $0.6 \pm 0.16 \text{ GeV}^2/c^2$, in close agreement with the observed value. The model may then be used to predict the variation of the muon-pair angular distribution with $x$, yielding the solid curve shown in Figure 2.

Additional and much more accurate studies of the structure of the pion will become available when analysis of the 255-GeV run of E 615 is complete. Figure 3 shows the raw mass spectrum of the muon pairs collected in the 255-GeV run, and Figure 4 shows a preliminary version of the pion structure function obtained from that data.

A confirmation that quark-antiquark annihilation is responsible for muon-pair production even at large $x_F$ is provided in Figure 5, which plots the ratio of the cross sections for muon-pair production by $\pi^+$ and $\pi^-$ beams. This ratio would be $1/4$ if only valence quarks contribute, while the solid curve in Figure 5 shows the expected ratio taking into account the sea-quark distribution of the nucleon. Note that for moderate $x$ the data points lie above the curve, but for large $x$ the data drop below the curve. This suggests that there are charge-symmetric QCD corrections to the cross section at moderate $x$ which die out at large $x$.

Another result from the analysis of the 255-GeV data does not concern pion structure, but rather properties of the decays of charmed mesons[12]. We have used our large sample of muon pairs to search for the decay $D^0 \rightarrow \mu^+\mu^-$, and set an upper limit on the branching fraction to this channel of $< 10^{-5}$. We also examined our sample of $\mu^+\mu^+$ pairs for evidence of $D^0-\bar{D}^0$ pair production followed by a 'charm oscillation' in which the $\bar{D}^0$ is converted into a $D^0$. Then both charmed mesons could decay semileptonically: $D \rightarrow \mu^+ X$. We are able to set an upper limit on the $D^0-\bar{D}^0$ mixing parameter of $6 \times 10^{-3}$, which is the best limit on this quantity obtained in any experiment to date.
Figure 3. The invariant-mass spectrum of muon pairs collected in the 255-GeV run, uncorrected for acceptance.

Figure 4. The pion structure function from a preliminary analysis of the 255-GeV data sample.
Figure 5. The ratio of the cross sections for muon-pair production by $\pi^+$ and $\pi^-$ beams. The solid curve is a calculation which assumes that quark-antiquark annihilation is the production mechanism.

References

HIGH \( p_T \) PHYSICS FROM THE OMEGA PHOTON COLLABORATION

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ABSTRACT

Preliminary data on \( p_T \) distributions of single particles from the photoproduction of multi-hadron final states by photons in the range 60 - 180 GeV are presented. Results indicate an excess of particles with \( p_T > 2 \) GeV at the higher photon energies in good agreement with Q.C.D. predictions.
Introduction

A study is in progress of \( Y\gamma \rightarrow x\gamma \)

With \( E\gamma = 60 \rightarrow 170 \text{ GeV}, \Delta \gamma \pm 20\% \) events including most of the hadronic cross-section.

The status is that we have preliminary results on single particle \( p_T \) distributions on \( \sim 15\% \) of data.

Data Analysis

The data has been analysed with a preliminary version of TRIDENT, the Omega pattern recognition and vertex reconstruction programme. This version is not fully tuned to avoid picking up incorrect track points and breaking up tracks. If these faults occur then track parameters have big errors. This occurs at few per 1000 level of all tracks at the present time and is potentially serious. The reason is that big errors give big \( p_T \) error and migration of tracks which were, in truth, low \( p_T \) to high apparent \( p_T \).

In order to avoid these effects to a reasonable extent we have put geometrical cuts (which should not affect \( p_T \) distributions) to clean up the data by only accepting those tracks which are detected over most of their traverse of the detector. Some aspects of these cuts (which will not be used on our final analysis) are not fully understood.

Data satisfying the above cuts have been plotted as a function of \( p_T \) for single particles for various \( E\gamma \) bins. The data at low \( p_T \) is clearly dominated by the VMD contribution and the aim is to extract any non-VMD component. At present we have used data from other experiments in the closest kinematic range where \( \gamma \) induced direct processes are surely absent, i.e. from \( \pi \rho \) interactions. The problems with this approach are

- statistics of these other experiments are inadequate
- range of energy not ideal
- may have bias in our or other data which is different

We have used the data of Donaldson et al. (Phys. Lett. 73B, 375 (1980)) \( x_f > 0 \), \( 0.5 < p_T < 3.5, E_x = 100, 200 \text{ GeV} \). The NA3 data is not in quite the same kinematic range - \( 0 < y < 0.4, 3.0 < p_T < 5.5, E_\pi = 200 \text{ GeV} \). The NA3 data probably shows a faster dependence on \( p_T \) and so would indicate a larger photon induced high \( p_T \) spectrum. The normalisations have been done at \( p_T \sim 1 \) and are consistent with the expected \( \pi \rho /\gamma \rho \) cross-sections. (i.e. our experiment gives a
We now examine the excess of high $p_T$ tracks over VMD expectations (as calculated above). The VMD expectations show a continuous smooth drop of the cross-section with $p_T$ in agreement with data at low $E_\gamma$ but in marked disagreement with the data at higher $E_\gamma$ where above $p_T \gtrsim 2$ GeV there is a substantial excess in the data. We have calculated the Q.C.D. contribution from Q.C.D. compton + Q.C.D. Bethe Heitler and added on to the VMD part from Donaldson et al., in order to compare with data. The curves shown on the figure are the result of performing this calculation. The agreement, except between $p_T = 1$ and 2 GeV is excellent. In fact we have also estimated the contribution from charm photoproduction ($\lesssim 10 \times$ the charm fraction in hadroproduction) and its $p_T$ dependance and it appears to fill in this short-fall between 1 and 2 GeV both in position and magnitude.

Conclusions

We have exciting indications of significant excess of high $p_T$ particles from $\gamma p$ over $\pi p$ in agreement with Q.C.D. expectations. However we recognise that we have a hard job to get the analysis fully convincing because we need very good pattern recognition to prevent spurious high $p_T$ tracks distorting our results. Hence the data cannot yet be taken as more than an indication of a Q.C.D. hard process. The conclusions are in good agreement with data from NA14.
\[ \frac{dN}{dP_T} \]

**$E_{\gamma} = 14.2$ GeV**

PRELIMINARY

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**$E_{\gamma} = 170$ GeV**

PRELIMINARY
DUAL DESCRIPTION OF A CONFINED COLOUR FIELD

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ABSTRACT

The result of perturbative QCD can be formulated in two dual or complementary
ways, either in terms of quarks and gluons or in terms of colour dipoles. The
relation between the two descriptions is similar to that between a lattice and
its dual lattice. If the confined colour field behaves like a vortex line in a
superconductor, then the dipoles form a chain along the vortex line.
In this talk I want to show that the results of perturbative QCD (in the case of e.g. an e⁺e⁻ annihilation event) can be formulated in two dual or complementary ways, either in terms of quarks and gluons or in terms of colour dipoles [1].

Due to the confinement property a colour field does not spread out as an ordinary electromagnetic field, when the charges are separated. It was early suggested that it may be compressed into vortex lines like the magnetic field in a type II superconductor [2]. Such a vortex line consists of a thin core, which is kept together by currents circulating around it, and surrounded by a more extended magnetic field, which at large distances is exponentially damped. The field of such a vortex line is the same as that of a chain of dipoles lined up along the vortex (see fig. 1). (Due to the damping in the superconducting medium this field is not just the sum of the fields from the first and last charges in the dipole chain.)

The dynamics of a vortex line is that of the massless relativistic string as long as it is not very strongly curled up, where the scale is determined by the thickness of the core. In the Lund model gluons emitted in an e⁺e⁻ annihilation event are assumed to act as excitations or kinks on a stringlike field [3]. Thus the field is stretched from the quark via the gluons to the antiquark. This picture implies a set of observable asymmetries and correlations which now have strong experimental support [4]. The gluons are ordered along the string and e.g. a red charge of one gluon and an anti-red charge of the next one form together one dipole in the vortex line chain, as shown in fig. 2.

The stringlike behaviour can also be understood from perturbative QCD [5]. If one gluon has been emitted with a given momentum, then the radiation of further soft gluons is that given by two dipoles, one formed by the quark and the gluon and the other by the gluon and the antiquark. There should also be subtracted, with the relative weight 1/N_c² the emission corresponding to one dipole formed by the quark and the antiquark. Thus soft gluons are emitted in those regions where the string model and the experiments give an excess of hadrons. If this scheme is generalized to more gluons we find that a system with n gluons a quark and an antiquark radiates like n+1 dipoles (with an O(1/N_c²) correction from the quark charges).
When one dipole emits a gluon this dipole is split into two dipoles (see fig. 3). If each dipole could be given a definite colour (e.g. red-antired) it would have been possible to determine which dipole end emitted the soft gluon. For the situation in fig. 3 the red charge in the dipole may be carried e.g. by a $rg$ gluon which emits a $rb$ gluon and turns into $bg$. This is not possible however. Because each dipole is a coherent combination of different colours we are not able to associate the emitted gluon with one of the dipole ends.

The physical state can be described in two alternative ways. Either one specifies the energy-momentum and polarization of all the gluons, quarks and antiquarks (and their colour ordering) or one specifies the energy-momentum and orientation (polarization) of all the dipoles. The dipoles are links which connect the gluons and the gluons are links which join the dipoles (cf fig. 2). Thus the relation between the two ways to describe the state is similar to the relation between a lattice and its dual lattice.

We now want to study the gluon emission from one dipole. If we go to the rest frame of one of the dipoles we have two gluons (or e.g. a gluon and a quark) which move in opposite directions with momenta $p_1$ and $p_2 = -p_1$. The emission of gluons which are soft compared to these momenta, is given by

$$dg = \frac{a_s}{4\pi^2} N_c \frac{dq_L^2}{q_L} dy d\phi$$

Here $q_L$, $y$ and $\phi$ denote the transverse momentum, rapidity, and azimuth angle for the emitted gluon.

We may (somewhat inadequately) regard gluons with $y>0$ ($y<0$) as emitted by the rightmoving (leftmoving) charge. We now make a (large) Lorenz transformation to e.g. the original total cms. For simplicity we assume that this boost is perpendicular to the dipole moment (and in the direction corresponding to $\phi=0$). The momenta after the emission are denoted...
and $q'_i$ in the old and the new Lorentz frame respectively (cf. fig. 4), and we introduce the variables $z$ and $Q^2$ by the relations

$$q'_z = z q'_i \quad (z q'_p)$$
$$Q = (q_1 + q_2)^2 = (q'_1 + q'_2)^2$$

For large enough $z$-values the distribution in eq. (1) now takes the form

$$d\sigma = \frac{\alpha_s(z)}{2\pi} \frac{dQ}{Q^2} \frac{dz}{z}, \quad z \gg Q^2/\pi^2$$

This corresponds to one half of the well-known expression for gluon splitting. The other half corresponds of course to the adjacent dipole. However, the kinematical constraint $y > 0$ implies a cut-off for lower $z$-values. This cut-off corresponds just to the situation where the angle $\theta$ in fig 4 equals the angle $\Psi$.

$$y > 0 \Rightarrow \theta < \Psi$$

We conclude that the result in eqs (3) and (4) corresponds exactly to the splitting of massive gluons into two gluons including the angular cut-offs caused by soft gluon interference, as discussed by Mueller, Marchesini, Webber and others [6].

In the same way gluons with $y < 0$ can be regarded as emitted by the other charge in the dipole. It is very natural that the formula for dipole radiation automatically reproduces the soft gluon cut-off caused by interference. A charge and anticharge do not radiate independently when the wavelength of the radiation is large compared to the relative distance. Of course the choice $y = 0$ as the borderline between the two regions is quite arbitrary. The distribution is smooth in $y$ and the choice of another boundary, e.g. $y = y_0$, only means that some gluons, previously regarded as emitted by one of the gluons, will instead be regarded as emitted by the other.

When a dipole emits a gluon it is split into two dipoles. Thus the production of a final state can be treated as a branching process where dipoles are split into smaller and smaller pieces. In the region where the virtualities are strongly ordered the results of perturbative QCD can therefore be formulated as a branching process in two different equivalent ways, either in terms of gluons which are split into two gluons at each branching, or in terms of dipoles which are split into two dipoles.

The formulation in terms of dipoles has however the great advantage that each branching is completely independent of the rest of the tree. In the gluon formulation the cut-off for low $z$-values depends on the other gluons. Also the
emission is not azimuthally symmetric around the gluon direction.

There will however be differences in the results if the branching processes are terminated at a fixed gluon mass or at a fixed dipole mass. Here we want to argue that a fixed dipole mass is a more relevant termination point. Two opposite charges which move along each other (i.e. with low invariant mass) do not radiate. On the other hand, if they move apart they can emit dipole radiation, even if they are not highly virtual. In Monte Carlo generation programs based on gluon splitting the termination often produces large masses for the colourless clusters (i.e. large dipole masses).

We note however that the termination point is less essential if the final system, after termination of the perturbative development, is treated like a string when it fragments into hadrons. In this case the emission of more and more very soft gluons does not modify the string state.

It is obviously also possible to include the break of the vortex line in two pieces by the process \( g \rightarrow q\bar{q} \). In this case the charges with momenta \( q_1 \) and \( q_2 \) in fig. 4 correspond to a quark and an antiquark, and there is no dipole which connects these charges. This process is suppressed relative to the normal one, first by a factor \( l/N_c \), and second because there is no z-pole in the \( (Q^2, z) \)-distribution.

Our conclusions are that the results of perturbative QCD can be formulated in two dual or complementary ways, either in terms of quarks and gluons or in terms of colour dipoles. The dipole formulation is "natural" if the colour field behaves like a vortex line in a superconductor. It is convenient because in the branching process each branching is independent of the rest of the tree. It can provide a link between perturbative QCD and the string picture.

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PHENOMENOLOGICAL INCORPORATION OF
THE EXACT QCD MATRIX ELEMENT
INTO QCD PARTON CASCADE MODELS

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ABSTRACT
A method for incorporating the exact matrix element into QCD parton cascade
models, in a consistent way, is discussed. The improved model (based on the Webber
model) is tested by comparing its predictions with data taken by JADE at PETRA.
The model predictions agree reasonably well with the data for the energy-energy-
correlation asymmetry and for other jet observables.

It has been shown by the JADE collaboration that the string fragmentation model of the
Lund group[1] combined with the second order QCD matrix element describes better than
independent fragmentation models the particle distribution in three jet events[2], the energy-
energy-correlation (EEC)[3], the charged multiplicity distribution[4] and the $p_t$ broadening
of gluon jets compared with quark jets[5]. Although the Lund model describes the data
quite well at PETRA and PEP energies, at higher energies (at SLC or at LEP) higher order
QCD effects become more important, because they are more visible, since the fragmentation
effects ($<p_t>\approx 0.3$ GeV) become relatively smaller. The following will be missing at higher
energies in any non-cascade models based on $O(\alpha_s^2)$ perturbative QCD (like the Lund model):

1. Events with large jet multiplicities ($N_{jet} > 4$)
2. Continuity of distributions for $qq$ and $qgq$ events$^2$.

Recently, a QCD parton cascade model based on the leading log approximation (LLA)
with soft- and collinear-gluon- interference effects approximated by an angular ordering of
partons has been developed (the Webber model[6]). Parton cascade models are very attrac-
tive, since the above two missing physics points can be admitted simultaneously: events with

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$^2$In the Lund model, $qq$- and $qgq$- distributions are smoothly connected by the fragmentation effects at
PEP/PETRA energies, if a very small value of the infrared cut off parameter ($y_{min}=0.0125$) is chosen[8].
more than 4 partons are naturally produced and the continuity between $q\bar{q}$ and $q\bar{q}g$ events is guaranteed by the existence of many soft gluons near the hard partons. There is, however, a serious problem in the QCD cascade models: the leading log approximation breaks down for hard processes. To solve this problem, the exact perturbative matrix elements have to be incorporated in the LLA parton cascade models.

I present here a method of achieving such an incorporation. The principle of the proposed method is that the gross structure of the event shape is described by the $O(\alpha_s)$ exact matrix element and the detailed structure (near jets) is described by the LLA parton cascade. The gross and the detailed structures are clearly defined by the invariant mass of parton combinations. Hence, the small $y$ region ($y < y_{\text{LIM}}$) is described by the LLA parton cascade and the large $y$ region ($y > y_{\text{LIM}}$) is described by the $O(\alpha_s)$ exact matrix element, where $y$ is the invariant mass squared of partons normalized by $E_{CM}^2$. It is necessary to choose $y_{\text{LIM}}$ sufficiently small that the LLA is good enough in the region $y < y_{\text{LIM}}$. The results should not depend too much on the matching point ($y = y_{\text{LIM}}$) if the value $y_{\text{LIM}}$ is small enough. The following jet algorithm is used to combine the final state partons into jets in order to define the variable $y$.

1. For a given event, the four-momenta of parton pairs are combined if the invariant mass is smaller than the limit ($y_{ij} < y_{\text{LIM}}$).
2. The parton pair with the smallest invariant mass is combined first.
3. Only the following combinations are allowed: $q + g \rightarrow q$, $\bar{q} + g \rightarrow \bar{q}$, $g + g \rightarrow g$, $q + \bar{q} \rightarrow g$. Other unphysical combinations, for example $u + u$ or $u + d$, are not allowed.
4. Keep at least one $q\bar{q}$ pair uncombined, because it couples with the initial virtual $\gamma$ and it is necessary to avoid events with only gluons.

The jet algorithm can be used in a different way, where the combination is terminated if the number of jets is equal to some small number, say two, independent of $y_{\text{LIM}}$. Since the parton branching history depends on the gauge, and the interference effects must be taken into account, the history (which gluon comes from which quark in which order etc.) must not be taken seriously, since it is not physical. Therefore the jet algorithm must be independent of the history.

In the Webber model, there are two event generators: a $q\bar{q}$-generator and a $q\bar{q}g$-generator, where the LLA parton cascade is started from the $q\bar{q}$ or the $q\bar{q}g$ initial state, respectively. The initial parton distribution ($q\bar{q}$ or $q\bar{q}g$) is given by the exact QCD matrix element of $O(\alpha_s)$. In the following I will discuss how to mix the event samples from $q\bar{q}$-generator and $q\bar{q}g$-generator without double counting and with the correct ratio. In order not to have double counting for the two samples, events from the $q\bar{q}$-generator are accepted only if the jet multiplicity is exactly two, and those from the $q\bar{q}g$-generator are only accepted if it is more than two. The jet multiplicity is defined by the jet algorithm with a parameter $y_{\text{LIM}}$ as described above. Double counting is thus avoided. In the next step, the normalization of the two samples has to be fixed. The two event samples should not be mixed with the ratio given by the $O(\alpha_s)$ or $O(\alpha_s^2)$ perturbative QCD calculations, since the higher order effects are effectively already taken into account by the LLA parton cascade in the Webber model. In order to obtain the normalization factor phenomenologically, partons are combined until the event has two jets and then $y$ is defined as the larger jet invariant mass squared normalized by $s$ of the two.

In Fig. 1a and 1b, the $y$ distributions are plotted for the $q\bar{q}$- and for the $q\bar{q}g$-sample, separately, with the parameter $y_{\text{LIM}} = 0.1$. The infrared cut off parameter $y_{\text{min}}$ for cal-
Calculating the $q\bar{q}g$ matrix element is set to 0.03 (this is technically the smallest value) as shown in Fig.1b. Since the LLA is a good approximation at small $y$, the height of the distributions for the two samples has to match in the small $y$. The normalization is thus obtained by requiring the two distributions to match at $y = y_{\text{LIM}}$, as indicated in Fig.1: 
\[
\frac{d\sigma}{dy}(y = y_{\text{LIM}}; q\bar{q}) = \frac{d\sigma}{dy}(y = y_{\text{LIM}}; q\bar{q}g).
\]

The matching point $y_{\text{LIM}}$ must be far from $y_{\text{min}}$, and at the same time, it must be small enough so that the LLA is a good approximation. In summary, the complete set of the events are taken from the $q\bar{q}$-sample with $y \leq y_{\text{LIM}}$ and from the $q\bar{q}g$-sample with $y > y_{\text{LIM}}$ and the normalization of the samples has to be fixed at $y = y_{\text{LIM}}$.

The model is experimentally tested by comparing the predicted distributions of jet observables with data. The first observable tested is the asymmetry in the energy-energy correlation (EEC). Since the EEC asymmetry is an additive observable, in the mixing of the $q\bar{q}$- and the $q\bar{q}g$-event samples, the total asymmetry $A(\theta) \cdot \sin \theta$ is a weighted sum of $A_{q\bar{q}}(\theta) \cdot \sin \theta$ and $A_{q\bar{q}g}(\theta) \cdot \sin \theta$:

\[
A(\theta) \cdot \sin \theta = f_{q\bar{q}} \cdot A_{q\bar{q}}(y \leq y_{\text{LIM}}; \theta) \cdot \sin \theta + (1 - f_{q\bar{q}}) \cdot A_{q\bar{q}g}(y > y_{\text{LIM}}; \theta) \cdot \sin \theta,
\]

where $f_{q\bar{q}}$ is the fraction of events from the $q\bar{q}$-generator obtained by the matching. As shown in Fig.2a, the prediction of the original Webber model in the large $\theta$ region, does not agree with the data (dots), since the LLA is poor in that region. The EEC asymmetry $A(\theta) \cdot \sin \theta$ after the incorporation of the exact matrix element is shown by the full curve, compared with the data (dots) in Fig.2b. The model prediction agrees reasonably well with the data, except for the small angle region.

The QCD scale parameter $\Lambda_{\text{QCD}}$ in the model was optimized to get a good agreement of the EEC asymmetry distribution with the data after incorporating the exact matrix element. The small value of $\Lambda_{\text{QCD}}$ (0.15 GeV) may partially be due to using $p_T^2$ as an argument of $\alpha_s$ in the model. As a test of the matching point invariance of the method, $y_{\text{LIM}}$ was

1 The discrepancy between the prediction and the data at small $\theta$ is essentially due to the overconstraint of the kinematics in the model: all the final state partons are on mass shell.

2 In the Webber model, the argument $Q^2$ of $\alpha_s$ is defined at each branching and is approximately $p_T^2$ of the branching. Therefore the $\alpha_s$ is a running coupling constant even within one event. In order to be consistent with this feature of the running coupling constant in the Webber model, the $O(\alpha_s)$ exact $q\bar{q}g$ cross section was modified as follows:

\[
\alpha_s(Q^2) = \frac{x_1^2 + x_2^2}{(1 - x_1)(1 - x_2)},
\]

where $Q^2$ is not a constant value of $E_{CM}^2$ but the same argument which is used in the Webber model. However, the difference in the EEC asymmetry distribution for the two cases, using the correct $Q^2$ or using $E_{CM}^2$, is small.
varied from 0.10 to 0.08 and to 0.12, and the total EEC asymmetry was plotted using the same procedure. The EEC asymmetry distribution is almost invariant under the change of \( y_{\text{LIM}} \) as shown in the Fig.2c and 2d. Other jet variables are also tested. For example, the sphericity distribution is significantly lower than the data for the original Webber model in the large sphericity region, as shown in Fig.3a. This can be improved by incorporating the exact matrix element (Fig.3b).

The method for incorporating the \( O(\alpha_s^2) \) exact QCD matrix element into the LLA parton cascade model would be an extension of the method just described. The LLA parton cascade is started from \( qq \) or \( qg \) states, as well as from \( q\bar{q}g \) or \( q\bar{q}q' \) states. The parton distributions for the above initial states are based on the \( O(\alpha_s^2) \) exact QCD matrix element. After the LLA parton cascade, the three event samples cannot be simply added because there are overlaps. To avoid double counting, 2-jet events are selected only from the \( qq \) event sample, 3-jet events are chosen only from the \( q\bar{q}g \) event sample, and \( \geq 4 \)-jet events from \( q\bar{q}g + q\bar{q}q' \) sample. The jet multiplicities for the three event samples are defined by the same jet algorithm with the same \( y_{\text{LIM}} \) for the three samples. The normalization of the sizes of the three event samples is determined by the matching condition at \( y = y_{\text{LIM}} \). The normalization of the size of the
sample from the $qg$-event generator to that from the $qg$ generator is obtained in the same way as discussed for $O(\alpha_s)$: the matching condition is $\frac{d\sigma}{dy}(y = y_{\text{LIM}}; qg) = \frac{d\sigma}{dy}(y = y_{\text{LIM}}; qg)$, where $y$ is the larger jet invariant mass squared (normalized by $s$) of the two jets, when the final state partons are combined until the events have two jets. The normalization of the $qg$-event sample to the $qg + qg'g'$-event sample is obtained by the following formula: 

$$\frac{d\sigma}{dy}(y = y_{\text{LIM}}; qg) = \frac{d\sigma}{dy}(y = y_{\text{LIM}}; qg + qg'g')$$

where $y$ is the largest jet invariant mass squared (normalized by $s$) of the three jets, when the partons are combined until the events have three jets.

Since it is not at present possible to calculate the exact $O(\alpha_s^2)$ QCD matrix elements, the incorporation of the $O(\alpha_s^2)$ QCD matrix element into LLA parton cascade model is the practical solution to handle the higher order QCD at higher energies. The method which I have discussed here can be used rather generally also for other dressing schemes, e.g. a la Sterman-Weinberg. The improved model can be a powerful tool not only for studying higher order QCD effects but also for controlling the background from ordinary multihadron events to look for new phenomena. For example, to extract heavy quark (top or 4-th generation) signal at SLC/LEP, the control of the spherical event background from higher order QCD is essential.

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ABSTRACT:

Hadron-hadron collisions are compared to lepton-hadron and $e^+e^-$ collisions under the assumption of partons acting as basic fields in all three. In this comparison, more significance is attached to differences observed in various types of correlation rather than to previously observed similarities in more simple-minded distributions. A number of non-trivial differences exist between $e^+e^-$ and lh (hard) collisions on one side and (soft) hadronic collisions on the other. Within a "dynamical" universality of similar color dipole or chain fragmentation in all types of collision, hadron-hadron collisions necessitate a number of chains, some with large angles relative to the others.
I. INTRODUCTION

The concept of "jet universality" is based on early observations of a similarity of particle production in \( e^+e^- \), deep inelastic lepton-hadron, and hard as well as soft hadron-hadron collisions.

This "naïve" universality could not be maintained at closer inspection. Diquarks were found to lead to lower average multiplicity \( \langle n \rangle \) than quarks, with \( \langle n \rangle \) growing slower with the hadronic mass \( W \) than that of quark jets\(^{1-3}\), but it did not seem to matter whether these quarks and diquarks were excited in lh or hh collisions. Similar differences between quark and diquark jets were found in the dispersion D.

In the following we want to show on a number of topics that also this "learned" universality cannot be maintained. It has to be replaced by a "dynamical" universality of (multiple) chain or colour dipole fragmentation.

II. NEGATIVE BINOMIALS

A parameter particularly sensitive to differences in hard and soft collisions turns out to be the parameter \( k \) of the negative binomial form\(^{4-6}\) recently used\(^{7}\) to describe multiplicity distributions up to collider energies. In fig.1 we reproduce \( 1/k \) as obtained by the UA5 Collaboration for non single-diffractive \( p^+p \) data from \( \sqrt{s}=10-900 \) GeV and compare it to \( 1/k \) obtained from fits to published \( e^+e^- \) multiplicity distributions from \( 7-35 \) GeV\(^{8-13}\). Clearly, \( 1/k \) is lower and rises more slowly with \( \sqrt{s} \) for \( e^+e^- \) collisions than for \( p^+p \) collisions.

To see whether the difference between \( pp^+ \) and \( e^+e^- \) is due to a typical hh effect or simply due to the quark-diquark difference discussed above, meson-proton (M\(^+\)p data\(^{14}\)) are compared to \( pp^+ \) and pp data in fig.2. The solid line (with a slope of 0.058) in both sub-figures corresponds to \( 1/k \) for \( p^+p \) of fig.1. The pp data of fig.2a and the M\(^+\)p data in fig.2b roughly follow the line (for the small difference between pp and M\(^+\)p data see ref.15).

On the other hand, the up data expected to be similar to M\(^+\)p data from "learned" universality, instead follow the trend of the \( e^+e^- \) data. The corresponding slopes of the dashed line fits are 0.023±0.007 for up and 0.016±0.003 for \( e^+e^- \) collisions. We have to conclude, that (soft) hh collisions show a \( 1/k \) behavior different from that of (hard) \( e^+e^- \) and lh collisions.

In fig.3a,b one can see for \( pp^+ \) data and \( e^+e^- \) data, that negative binomials also give perfect fits to the multiplicity distributions for rapidity intervals around the center. In both cases, \( k \) increases with the size of the interval. At \( y=0 \), \( k \) is equal for M\(^+\)p and pp collisions at the same energy\(^{15}\), but different for \( e^+e^- \) collisions. It would be important to see the values for up data.
Fig. 1 Parameters $n^{-1}$ and $k^{-1}$ for negative binomials and $D^2/\langle n^2 \rangle$ for $p^+ p$ data\textsuperscript{7} and $k^{-1}$ for $e^+ e^-$ data\textsuperscript{13}, as a function of $\sqrt{s}$. 

Fig. 2 The parameter $k^{-1}$ from a) fits to $n>6$ pp data and $n>2$ $e^+ e^-$ data, b) to $n>6$ $M^+ p$ and to $u^+ p$ data. The solid line is that from fig. 14. The dashed lines are linear fits to the $e^+ e^-$ and $u^+ p$ points, respectively\textsuperscript{4}. 

Fig. 3 Multiplicity distributions in the indicated (pseudo)-rapidity intervals for $p p$ collisions at 540 GeV (fig. 3a,c\textsuperscript{16}) and $e^+ e^-$ collisions at 29 GeV (fig. 3b\textsuperscript{13}). The solid lines in figs. 3a,b are negative binomial fits, in fig. 3c DTU predictions\textsuperscript{20}. 

What do models tell us about the above observations? From fig. 1 one can see that at large energies $D^2/\langle n^2 \rangle$ is rising with energy due to the increase of $1/k$. KNO scaling \textsuperscript{18} predicts $D^2/\langle n^2 \rangle$ to be independent of energy and is therefore excluded for $h h$ collisions in the full rapidity region.

A beautiful comparison of the two classes of (partially) stimulated emission and cascade models is performed in ref. \textsuperscript{6}. Experimentally, stimulated emission can now probably be excluded from $k$ values being larger for negatives than for all charged particles\textsuperscript{15}. 

The Lund model as it stands\textsuperscript{18)} gives too narrow multiplicity distributions for $\text{hh}$ collisions, even at low energies. However, an interesting new two-chain Lund model (LUND '86) with gluon emission\textsuperscript{19)} remains to be tested. Closer to negative binomials come the predictions from the DTU model\textsuperscript{20)}, even though the simple functional form itself cannot be derived from the model. In Fig.3c, a comparison with the UA5 data is shown. The bare prediction is still slightly too narrow. There may be room for another mechanism.

More fundamentally, Malaza and Webber\textsuperscript{21)} derive QCD predictions for the first five moments of the multiplicity distribution, and find that they are close to those of a negative binomial distribution.

For further discussion of this challenging and quickly developing topic see refs.\textsuperscript{16,22-25).}

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Fig. 4 a) The $q_x$ distribution for like charge pairs divided by unlike pairs for $\text{pp}$ collisions at $\sqrt{s}=63$ GeV\textsuperscript{26)}, b) the radius $r$ as a function of the charge multiplicity at $\sqrt{s}=53$ and 63 GeV\textsuperscript{27)}, c) the $q_x$ distribution for like pairs divided by uncorrelated background and d) $r$ as a function of the multiplicity for $\text{e}^+\text{e}^-$ collisions at 29 GeV\textsuperscript{29}). The solid lines are fits to a Bessel function in $q_x$.
III. BOSE-EINSTEIN CORRELATIONS

An interesting difference between $p^+p$ and $e^+e^-$ collisions can be observed in the multiplicity dependence of the size of the meson emitting region. The radius $r$ and an incoherence $\lambda$ can be estimated from the correlation of two identical bosons at small (transverse) distances $q_t$ in momentum space. Recent measurements of the radius $r$ in $\alpha x$, $pp$ and $p\bar{p}$ collisions come from the AFS and SFM collaborations. In both cases, $r$ is slightly larger than $1\text{fm}$ and the incoherence is $\lambda=0.5$. The AFS collaboration observes a dependence of $r$ on $n$, as shown in fig.4b. In high multiplicity events, the bosons appear to originate from a larger space-time region.

For $e^+e^-$ collisions, Bose-Einstein correlations have been measured by TPC and TASSO. Using the same (spherical) parametrization as in the ISR experiments grants similar $r$ and $\lambda$ as for $pp\bar{p}$ collisions. However, here $r$ does not depend on $n$ (fig.4d).

Introduction of Bose-Einstein correlations into existing string models looks quite natural and good results have already been obtained for $e^+e^-$ collisions.

IV. TRANSVERSE MOMENTUM DEVELOPMENT

A handy distribution to trace hard effects is $\langle p_t \rangle$ or $\langle p_t^2 \rangle$ versus $x$ (the "sea-gull"). Neutrino experiments have shown that already at $W<10\text{GeV}$ the sea-gull is lifting its wings, in particular the current fragmentation wing, as $W$ increases. Fig.5a gives the sea-gull for $\mu\bar{\mu}$ collisions compared to the Lund model with standard three-jet parameters (solid curve), no three-jet events (dashed), no soft gluons (dot-dashed) and no soft gluons but $\langle k_t^2 \rangle=0.88\text{ GeV}^2$ (dotted). According to this parametrization, soft and hard gluons seem responsible for the high $\langle p_t^2 \rangle$ values.

A very similar behavior is observed for $e^+e^-$ annihilation (fig.5b). The narrower jet has little energy dependence, while the wider jet shows rapid increase of $\langle p_t \rangle$ with energy. The curve corresponds to a QCD independent jet fragmentation model, but predictions from the three-jet string model are similar.

However, a rise of the sea-gull wings is also observed in $hh$ collisions! Like for $lh$ collisions, this rise has been observed already at lower energies and is now seen to persist at $\sqrt{s}=22\text{ GeV}$ (fig.5c). Here, the increase is visible in both wings and may be the onset of hard parton scatters and/or gluon emission. In fig.5d, the standard low-$p_t$ Lund model cannot reproduce the effect.

At higher energies, the increasing importance of the intermediate $p_t$ region is observed in the form of "mini-jets". These are defined with the UA1 jet-finding algorithm as jets with transverse energy $E_{T,j}>5\text{ GeV}$ and an axis with
Fig. 5 Values of $<p_T>$ as a function of $x$ for a) $pp$ collisions with $40 < W < 400$ GeV, b) $e^+e^-$ data with respect to the thrust axis (the curves are described in the text). Values of $<p_T>$, weighted by phase space for c) $K^+p + \pi^-X$ from 12.7 to 250 GeV/c and d) $\pi^+p + \pi^-X$ at 250 GeV/c compared to standard low-$p_T$ Lund.

$|n|<1.5$ and azimuth angle $>30^\circ$ from the vertical. The acceptance corrected fraction of events containing at least one mini-jet (semi-hard component) increases roughly logarithmically from 12% at 200 GeV to as much as 35% at 900 GeV.

Events without mini-jets (soft component) is characterized by low $<n>$, large $D$ and low $<p_T>$, the semi-hard component by large $<n>$, small $D$ and high $<p_T>$. As shown in fig. 6 for 200 and 900 GeV, the $n$ dependence of $<p_T>$ is much less pronounced for the semi-hard component than for the soft component.
The development of low-$p_t$ models in the direction of semi-hard effects is well on its way. LUND'86\(^{19}\) has the flexibility of both of a rotation of the two color dipoles with respect to each other and/or of gluon radiation within the dipoles. These options may be enough to explain the increase of $\langle p_t \rangle$ in the sea-gull wings at fixed-target energies.

At the collider, more chains seem to be needed as they appear in DTU\(^{44,45}\). Bopp et al.\(^{45}\) compare two classes of models within DTU, one where gluon emission induces large fluctuations in the parent parton transverse momentum, the other where partons acquire transverse momentum via a hard scattering. They find that both approaches describe the data equally well, and that the transition from soft to semi-hard processes is a smooth one.

V. AZIMUTHAL CORRELATIONS

An important question in hadronization is short range order from $q\bar{q}$ pair production. Studying this from non-strange mesons is hampered by the $q\bar{q}$ combinatorial background. What is needed is a flag identifying the pairs having been created together. Because of the strange sea suppression, there is less combinatorial background for $s\bar{s}$ pairs and this flag does indeed exist, there.

For $e^+e^-$ annihilation good strangeness identification is available in the TPC detector at $\sqrt{s}=29$ GeV. This collaboration observes\(^{66}\) significant short range $K^+K^-$ correlations in $\gamma$. It is well reproduced by the Lund model\(^{18}\) and by the Webber QCD model\(^{47}\).
A more stringent test than the associated strangeness density used above is the azimuthal angular correlation $\Delta \phi$ between the transverse momenta of pairs of strange particles as a function of $\Delta y$. For $\Lambda \Lambda$ pairs, azimuthal correlation has been observed in MARKII at 29 GeV. Similar $K^+K^-$ correlations are seen in the exclusive final state $K^-p\Lambda pK^+K^-\pi^+\pi^-$ at 32 GeV/c.

The results of a systematic study of the $\Delta \phi$ correlation in pp collisions at 360 GeV/c are given in fig. 7. There, the asymmetry parameter $B = \frac{N(\Delta \phi<\pi/2) - N(\Delta \phi<\pi/2)}{N_{all}}$ is given for pairs with (a) opposite strangeness and small $\Delta y$, (b) same strangeness and small $\Delta y$, and (c) same or opposite strangeness, but large $\Delta y$. In general, class (a) has larger asymmetry $B$ than the expectedly uncorrelated classes (b) and (c), and larger $B$ than expected from independent emission (dashed line). However, the asymmetry for class (a) (and also (b)) pairs is smaller than expected from the standard Lund model.

The azimuthal correlation can of course also be studied for $\bar{c}\bar{c}$ in $D\bar{D}$ production. An asymmetry has indeed been observed in $\pi^-p$ collisions at 360 GeV/c. Also there, the Lund model overestimates the effect.

We believe that the overestimation of $B$ in Lund is related to the underestimation of the height of the sea-gull wings.

VI. POLARIZATION

A particularly interesting difference between $hh$, $lh$ and $e^+e^-$ collision is to be expected in hyperon polarization. While the previous comments were on typical hadronization properties, polarization is at least partially determined...
by the production (excitation) mechanism on the parton level.
Because of space limitation I have to refer to my comparison in ref. 54). Here, I just want to say that the available data on hyperon polarization show the differences expected for hh, lh and e⁺e⁻ collisions, but a more differential study in higher statistics data would be welcome for the latter two.

VII. DIFFRACTION DISSOCIATION

So far, we were mainly concerned with a comparison of e⁺e⁻, lh and non-diffractive hh collisions. An important question left is that of diffractive hadron production. Is the decay of the diffractively excited system more or less isotropic or is it elongated like a fragmentation chain?

The R608 Collaboration 55) has studied the exclusive channels pp → (Λ⁺K⁺) p and pp → (ΛΛ⁺p) p at √s=63 GeV, with the bracketed system carrying momentum near that of the beam. A difference from isotropic decay is clearly observed in the Gottfried-Jackson angular distribution of the decay products in fig. 8. The Λ⁺ which can carry a ud diquark of the beam proton, is peaked in the direction of that proton in fig. 8a. In fig. 8c), the K⁺ probably carrying the remaining u quark is peaked in the opposite direction. The Φ does not carry any proton valence quark and is more central (fig. 8b).

Fig. 8 Gottfried-Jackson angular distribution of the particles indicated in the diffractively produced forward system of pp collisions at √s=63 GeV. The solid lines correspond to isotropic phase space events passed through the acceptance of the apparatus.

Fig. 9 Pseudorapidity distributions of charged tracks from the fragmentation of diffractive states of average mass 〈M〉 as indicated, at √s=546 GeV. The arrows show the expected position of the center of the cluster and that of the inner edge.
The above observation is in agreement with the idea\textsuperscript{56} of pointlike pomeron-quark coupling. This leads to the back-scattering of one quark in the proton with a continuing spectator diquark as in deep-inelastic scattering (for earlier thoughts in this direction see refs.\textsuperscript{57,58}). The consequent elongation of the diffractively produced system along the pomeron-proton collision axis is observed in the exclusive final states $pp \rightarrow (p^{+}p^{-}p^{+}p^{-})p$ by the same collaboration\textsuperscript{59} and at 360 GeV/c\textsuperscript{52} and earlier in $\gamma p \rightarrow (p^{+}p^{-}p^{+}p^{-})p$ by the Omega Photon Collaboration\textsuperscript{60}. The hadronization of the diffractive system is well described by a Lund string\textsuperscript{59,61} similar to that of $\ell\ell$ collisions.

Early results on inclusive diffraction dissociation\textsuperscript{62,63} derived from (pseudo-) rapidity distributions were inconclusive. At collider energies, however, a rapidity plateau develops\textsuperscript{64} at the highest diffractive masses (see fig.9). The central rapidity density of the diffractive cluster rises with the excitation mass $M$ in a way very similar to the rise of that in non-diffractive events with $\sqrt{s}$ (not shown). The same holds for $\langle n \rangle$ compared to that for non-diffractive events. The results can be reproduced\textsuperscript{65} within DTU, where chains are stretched between valence constituents of the excited proton and sea constituents of the non-excited one.

But what about $q\bar{q}$ systems? They are simpler than $q(q\bar{q})$ systems and more straight-forward to compare to $e^+e^-$ results. The obvious place is to look in high energy meson diffraction dissociation. There, the disadvantage of the relatively low energy is compensated by the increased rapidity range on one hand, and the availability of very differential data on the other.

The NA22 Collaboration\textsuperscript{66} separates the inclusive single diffractive

![Graph](image_url)
component by a combined rapidity and rapidity gap method. In fig. 10a we compare preliminary results for the thrust distribution of the K+ diffractive system (with average excitation mass of 7.4 GeV) to PLUTO data at 7.7 GeV. One can see that the diffractive system is at least as elongated along the thrust axis as the e⁺e⁻ data. Fig. 10b shows the excitation mass (respectively s) dependence of <n> for the diffractive system and for e⁺e⁻ collisions (both including charged pions from K_s decays). The agreement justifies further investigation.

VIII. SUMMARY AND CONCLUSIONS

Hadron-hadron collisions have been compared to lepton-hadron and e⁺e⁻ collisions under the assumption of partons acting as basic fields in all three. In this comparison, more significance has been attached to differences observed in correlations rather than to previously observed similarities in simpler distributions. The results are summarized in the following table:

<table>
<thead>
<tr>
<th>Effect</th>
<th>e⁺e⁻, 1h</th>
<th>(soft) hh</th>
<th>Models</th>
</tr>
</thead>
<tbody>
<tr>
<td>Neg. Binomials</td>
<td>e⁺e⁻ ≈ wp</td>
<td>M⁺p ≈ pp</td>
<td>stimul. em., casc.</td>
</tr>
<tr>
<td>1/k vs. √s</td>
<td>small</td>
<td>large and steep</td>
<td>- &quot; - DTU</td>
</tr>
<tr>
<td>k at y=0</td>
<td>large</td>
<td>small</td>
<td>DTU?</td>
</tr>
<tr>
<td>k vs. ∆y</td>
<td>rapid increase</td>
<td>slow increase</td>
<td>k⁻ &gt; kᶜᶜ</td>
</tr>
<tr>
<td>negatives</td>
<td>?</td>
<td></td>
<td>stimul. em. out</td>
</tr>
<tr>
<td>Bose-Einstein</td>
<td>r&gt; 1 fm</td>
<td>= &gt; 1 fm</td>
<td>easy to incorporate, first results in string models</td>
</tr>
<tr>
<td>n dependence</td>
<td>no</td>
<td>≈ yes</td>
<td>DTU</td>
</tr>
<tr>
<td>λ</td>
<td>= 0.5</td>
<td>= 0.5</td>
<td>Lund '86, DTU</td>
</tr>
<tr>
<td>Pₜ development</td>
<td>wings rise</td>
<td>wings rise</td>
<td>hope in Lund, DTU</td>
</tr>
<tr>
<td>&quot;sea-gull&quot;</td>
<td>yes</td>
<td>yes</td>
<td>for hh too strong</td>
</tr>
<tr>
<td>&quot;mini-jets&quot;</td>
<td>yes</td>
<td>yes</td>
<td>as expected</td>
</tr>
<tr>
<td>Pₜ correlations</td>
<td>yes</td>
<td>yes</td>
<td>P-Y analogy</td>
</tr>
<tr>
<td>Polarization</td>
<td>e⁺e⁻ : no</td>
<td>transverse</td>
<td>Lund '86, DTU</td>
</tr>
<tr>
<td>A</td>
<td>1h : longit.</td>
<td>(rise with Pₜ)</td>
<td></td>
</tr>
<tr>
<td>Diffraction</td>
<td>meson ≈ e⁺e⁻</td>
<td>proton = 1h</td>
<td></td>
</tr>
</tbody>
</table>

A number of non-trivial differences exist between e⁺e⁻ and 1h (hard) collisions on one side and (soft) hadronic collisions on the other. Within a "dynamical" universality of similar color dipole or chain fragmentation in all types of collision, hadron-hadron collisions necessitate a number of chains, some with large angles relative to the others.
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RAPIDITY DEPENDENCE OF MULTIPlicITIES IN
NON-DIFFRACTIVE n+p AND pp COLLISIONS AT 250 GeV/c.

NA22 COLLABORATION

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ABSTRACT:

Multiplicity distributions of charged and negative particles for non
(single)-diffractive (NSD) events in n+p and pp collisions at 250 GeV/c
(\(\sqrt{s}=22\) GeV) are presented for various rapidity intervals.
- The data are well described by the negative binomial (NB) distribution.
- Comparing the n+p and pp data, the \(k\) and \(n\) parameters of the NB are different
  for the full phase space, but similar for the central region.
- In the central region, the relation \(k_{ch}^{-1/2} k_{\bar{n}}\) holds between the \(k\) parameters
  of the charged and negative multiplicities. This favours the interpretation
  of NB's in terms of cascade models above (partial) stimulated emission models.

Negative binomials\(^{f)}\) appear to describe charged multiplicity distributions
rather well. This holds for full phase space in NSD \((\bar{p})\) collisions over the
whole available energy range (\(\sqrt{s}=10-900\) GeV) and for limited (pseudo) rapidity
intervals for pp at \(\sqrt{s}=540\) GeV and e\(^+\)e\(^-\) annihiliations at \(\sqrt{s}=29\) GeV\(^{1)}\). Knowledge

\(^{f)}\) The negative binomial distribution is given by

\[
P_n(n,k) = \frac{(n+k-1)!}{n!(k-1)!} \left(\frac{k}{k+n}\right)^n \left(1-\frac{n}{k}\right)\]

\(^{1)}\)
on the multiplicity distribution in various restricted regions of phase space is limited for hadronic data at energies below the collider and absent for negative particles. Here, new results on \( \pi^+p \) and \( pp \) collisions are presented.

The experiment (NA22) has been performed at CERN with the European Hybrid Spectrometer (EHS). Charged particle momenta are measured over the entire solid angle. The present NSD sample consists of \( \sim 6800 \pi^+p \) and \( \sim 2500 \, pp \) events. More details can be found elsewhere\(^2,3\).

The multiplicity distributions of charged and negative particles for full phase space and in selected rapidity intervals \(|y|<y_{\text{cut}}\) were fit to the NB distribution. (The rapidity plateau extends to about \( \pm 1.5 \) units; the kinematical limit for a pion is \(|y|=5\).) Fig. 1 and 2 show the \( \pi^+p \) data, expressed in KNO form, together with the best fit. Each successive distribution is shifted down by a factor ten. Except for 2 and 4 prongs in the full distribution, the errors are statistical only. The quality of all fits is good, both for the \( \pi^+p \) and \( pp \) data. The values of the \( \bar{n} \) and \( k \) parameters are shown in fig. 3 and 4 versus the size of the rapidity interval. (Except for full phase space, errors are statistical.) Comparing the \( \pi^+p \) and \( pp \) data, the \( k \) and \( \bar{n} \) parameters of the NB are different for full phase space, but similar for the central region. For the
full charged $\pi^+ p$ ($pp$) multiplicity $n_{ch}=9.21^{\pm}.05$ ($8.76^{\pm}.08$) and $k_{ch}=17.3^{\pm}.1$ ($12.4^{\pm}.8$). (The difference remains, when subtracting 2 from the total multiplicity to account for leading particles$^4$; $k_{ch}=7.3$ (5.1).) Comparing the behaviour of charged and negative particles in the central region, one observes $n_{ch}=2n_-$ (as expected), and $k_{ch}=1/2k_-$. The NB distribution can be generated by at least two mechanisms: 1) by (partial) stimulated emission of identical bosons 2) by cluster production and decay (cascading)$^4,5)$. Since the charged particles mainly comprise two kinds of identical bosons ($\pi^+$ and $\pi^-$), models of the first type predict $k_{ch}=2k_-$. This is in conflict with our data.

I thank P. van Hal and L. Scholten for their help in the analysis and L. Van Hove for interesting discussions.

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FORWARD–BACKWARD MULTIPLICITY CORRELATIONS
IN PP–COLLISIONS AT $\sqrt{s} = 200$ AND 900 GEV

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Presented by B. Holl
Physikalisches Institut, Bonn University, Germany

Abstract

Preliminary results on forward–backward multiplicity correlations at $\sqrt{s} = 200$ and 900 GeV are presented. We observe up to $\sqrt{s} = 900$ GeV an increase of the strength of forward–backward multiplicity correlations proportional to ln $s$. The analysis of these correlations in terms of a simple cluster approach yields values for the effective cluster size for the decay in charged particles of $k_{\text{eff}} \sim 2.6$. This value is in a good agreement with those determined at ISR energies and at $\sqrt{s} = 546$ GeV.
I. Introduction

A possibility to get a deeper understanding of the mechanism of particle production in high energy reactions is the investigation of correlations among particles emitted at various values of rapidity $y = 1/2 \ln(E + p_j)/(E - p_j)$. As reported by different experiments, the study of the two particle rapidity correlations leads to the conclusion that the final state particles are produced in clusters where the decay width of these clusters is $\sim 0.7$ units in rapidity [1]. This observation may be interpreted as an effect of resonance productions.

Besides the above short-range correlations, data from the ISR [2] and from the CERN pp-collider at $\sqrt{s} = 546$ GeV [3] show the existence of correlations extending over a larger range in rapidity. Furthermore the data show an increase of these long-range correlations with energy. The predictions of the Dual-Parton-Model [4,10] and of the UA5 cluster Monte-Carlo, which has been tuned to describe the basic features of multiparticle production [5,6], are compared with the data.

We present here new results on forward-backward multiplicity correlations at $\sqrt{s} = 200$ and 900 GeV. The results are based on $\sim 2000$ ( $\sim 1000$ ) minimum bias events at 900 GeV ( 200 GeV ). The data were taken with the UA5 streamer chamber detector in March/April 1985 with the CERN pp-collider operated in a pulsed mode. The UA5 detector allows the measurement of the center-of-mass scattering angle $\theta$ of the charged particles produced (there is no magnetic field). As an approximation of the rapidity $y$ we use the pseudorapidity $\eta = -\ln \tan(\theta/2)$. Further details on the detector can be found in [5], on data taking during the pulsed collider run and on first results in [7-9].

II. Observation of forward-backward multiplicity correlations

We define two symmetric, non-overlapping intervals in the pseudorapidity space $|\eta| < 4.0$ and call $n_F$ the multiplicity of charged particles in the forward $\eta$-interval $\eta_1 < \eta < \eta_2$ ( $\eta_1$ and $\eta_2$, respectively, denote the upper and lower limit of that interval) and $n_B$ the multiplicity in the backward interval $-\eta_2 < -\eta < -\eta_1$.

The fig. 1a shows a scatter plot of events on the $n_F, n_B$ plane for 900 GeV for the case that the forward and backward intervals are in contact ( $\eta_1 = 0$ ). Experimentally one has found that the average of $n_B$ at fixed $n_F$ as function of $n_F$ is remarkable well described by a linear function of $n_F$ (see also fig. 1b), i.e. $n_F$ and $n_B$ are linear correlated. The straight line in the scatter plot is obtained by computing $< n_B >$ at fixed $n_F$, where the slope and intercept are taken from a straight line fit, $< n_B(n_F) > = a + b \cdot n_F$. The slope parameter $b$ is identical to the correlation coefficient of $n_F$ and $n_B$ (see equation 2). Due to the presence of short-range correlations one would expect such correlations, because clusters produced at $\eta \sim 0$ can emit their decay products into the forward and backward region simultaneously. Introducing a gap of 2 units in pseudorapidity between the forward and backward region one can decouple these regions from the short-range correlation effects. Fig. 1b shows $< n_B(n_F) >$ vs $n_F$ at 900 GeV with a gap of $\delta \eta (\sim 2 \cdot |\eta_1|) = 2.0$ between the two regions. Even if the two regions are decoupled by 2 units in pseudorapidity there is still a remaining linear correlation between $n_F$ and $n_B$, which proves the
Fig. 1. (a) Two dimensional distribution of events as function of $n_B$ and $n_F$ for the $\eta$-range: $0 < |\eta| < 4.0$ at 900 GeV. (b) Average value of $n_B$ taken at fixed $n_F$, as a function of $n_F$ for the $\eta$-range: $1.0 < |\eta| < 4.0$ at 900 GeV.

Fig. 2. The correlation strength $b$ as function of a central pseudorapidity gap $\delta \eta$.

Fig. 3. The correlation strength $b$ as function of the energy for three different $\eta$-intervals.
existence of long-range correlations. In addition fig. 2 shows for $\sqrt{s} = 200, 546$ and 900 GeV, respectively, that the correlation strength $b$ depends on the gap size $\delta \eta$ between the forward/backward region. The error bars indicate the statistical error only. The curves are shown in this figure are Monte-Carlo predictions based on the Dual-Parton-Model [10]. From fig. 2 one observes that the dependence of the correlation strength on the gap size is reproduced by the Dual-Parton-Model.

The energy dependence of the correlation strength $b$ for three different forward/backward regions are shown in fig. 3. In the central region the observed correlations are dominated by the short-range correlations. Apparently $b$ still rises linearly with $\ln s$ in the energy range considered (the straight lines are drawn to guide the eye).

III. Analysis of forward-backward multiplicity correlations

As is seen in fig. 3, the correlation strength is reduced but not zero if one introduces a gap of 2 units in pseudorapidity between the forward and backward regions. An explanation in terms of a simple cluster approach – generation of hadrons in groups with a decay distribution reflecting effects of short lived low mass resonance decay and Short Range Order – will be described below. From the correlation between the quantities $n_F$ and $n_B$ it follows that deviations (fluctuations) from the average number of particles in one region are accompanied by similar fluctuations in the other region. If one assumes that the observed charged particles at fixed number of $n_S = n_B + n_F$ are independently emitted with equal probability to fall into the forward or backward region and uses the fact that even in smaller $\eta$ ranges than the full phase space the charged multiplicity distribution is well described by a negative binomial distribution [11], one may express the correlation strength in terms of the two parameters $(N, K)$ of the negative binomial distribution:

$$b = \frac{\bar{N}}{2K + \bar{N}}. \quad (1)$$

Inserting the values for $\bar{N}$ and $K$ for the range $0 < |\eta| < 4$ at 546 GeV one obtains $b=0.81 \pm 0.01$. The comparison with the corresponding one from fig. 3 clearly shows that one obtains too high fluctuations. Therefore the assumption that the charged particles independently branch (binomially distributed with $p=1/2$) into the forward/backward region is wrong.

It can be shown that the following identity holds [3]:

$$b = \frac{\text{cov}(n_F, n_B)}{\sqrt{\text{var}(n_F)\text{var}(n_B)}} = \frac{1/4D_S^2 - <d_S^2(n_F)>}{1/4D_S^2 + <d_S^2(n_F)>>} \quad (2)$$

where $D_S^2$ is the variance of the $n_S$-distribution and $d_S^2$ the variance of the $n_F$-distribution at fixed $n_S$ and $<>$ denotes an average value over the $n_S$-distribution. Under the assumption of all clusters being independently generated and binomially branched into the forward/backward region and all charged particles emitted from a cluster remaining in the same region, one obtains:

$$<d_S^2(n_F)> = \frac{1}{4} k_{\text{eff}} <n_S> \quad ,$$
where $k_{\text{eff}}$ is the so-called effective cluster size and is a function of the first two moments of the unknown cluster decay distribution ($k_{\text{eff}} = \bar{k} + d_{k}^2 / \bar{k}$) \cite{12}. This allows the determination of the effective cluster size independently from the investigations of the two-particle correlation function.

In fig. 4a (4b) the quantity $k_{\text{eff}}(n_S) = 4d_{S}^2(n_F)/n_S$ is plotted vs $n_S$ at 900 GeV (200 GeV) for different forward/backward regions $\Delta \eta = |\eta_2| - |\eta_1|$ and fixed gap size of 2 units in pseudorapidity. If the size of the forward/backward region is reduced from $\Delta \eta = 3$ to $\Delta \eta = 2$ and $\Delta \eta = 1$, respectively, the probability that a cluster emits a particle outside this region increases. Therefore, due to these leakages of cluster products a cluster will appear smaller than generated. This is seen in the data and is also reproduced by the UA5 cluster Monte-Carlo, as indicated by the curves. As an input for the cluster formation the Monte-Carlo uses for the mean charged particles per cluster a value of $<k> \sim 1.8$ and an effective cluster size of $k_{\text{eff}} \sim 2.6$. The effective cluster size of $k_{\text{eff}} \sim 2.6$ is in a good agreement with the values obtained at 546 GeV from our analysis of the forward–backward multiplicity correlations and from the semi inclusive two–particle pseudorapidity correlations \cite{13} and with earlier ISR results \cite{14}. Therefore it appears that $k_{\text{eff}}$ is independent of energy and the rise in $b$ results from the increase in $D_S^2 / <n_S>$.

![Fig. 4.](image)

Fig. 4. The ratio $4d_{S}^2(n_F)/n_S$ as a function of the total multiplicity $n_S$ for three different forward/backward intervals at 900 GeV (fig. 4a) and 200 GeV (fig. 4b).
IV. Conclusions

We observe up to $\sqrt{s} = 900$ GeV an increase of the strength of forward–backward multiplicity correlations proportional to ln $s$. The analysis of these correlations yields values for the effective cluster size of $k_{\text{eff}}(= k + d^2 k/\bar{k}) \sim 2.6$ in a good agreement with earlier ISR results and results at 546 GeV. Therefore the continuous growth with energy of the correlations strength is associated with the growth of $D_2^f < n_S>$.

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We review properties of the negative binomial distribution, along with its many possible statistical or dynamical origins. Considering the relation of the multiplicity distribution to the density matrix for Boson systems, we re-introduce the partially coherent laser distribution, which allows for coherent as well as incoherent hadronic emission from the k fundamental cells, and provides equally good phenomenological fits to existing data. The broadening of non-single diffractive hadron-hadron distributions can be equally well due to the decrease of coherence with increasing energy as to the large (and rapidly decreasing) values of k deduced from negative binomial fits. Similarly the narrowness of e⁺e⁻ multiplicity distribution is due to nearly coherent (therefore nearly Poissonian) emission from a small number of jets, in contrast to the negative binomial with enormous values of k.
I. NEGATIVE BINOMIAL DISTRIBUTION

The negative binomial distribution

\[ p_n^k = \frac{(n+k-1)!}{n!(k-1)!} \left( \frac{\bar{n}/k}{1+\bar{n}/k} \right)^n \left( 1+\bar{n}/k \right)^{n+k} \]  

(1)

with \( \bar{n} \) (the average multiplicity) and \( k \) as the parameters, has been found to give an excellent account of hadronic multiplicity distributions\(^1\). In particular the recent fit to non-single diffractive data from all energies by the UA5 group is especially interesting in that the parameter \( k \) is required to decrease rather rapidly with energy. Put differently, the scaling form of (1) (a special case of the gamma distribution)

\[ \bar{n} p_n^k \sim \psi_k(z) = \frac{k^k}{(k-1)!} z^{k-1} e^{-kz} \]  

(2)

\[ z = n/n^- \]

does not exhibit KNO scaling, as would be the case were \( k \) constant. These facts, in particular the energy dependence of \( k \) have inspired many theoretical conjectures about the meaning of (1) and its physical meaning.

The negative binomial distributions occurs in many physical and mathematical contexts. Here we mention of few examples, referring to standard mathematical texts\(^3\) and our forthcoming review article\(^4\) for more information on this and related distributions.

(1) \( p_n^k \) is a generalized Bose-Einstein distribution composed of \( k \) (integral) cells of equal average occupancy \( \bar{n}/k \) (see Knox\(^1\) and ref. 3).

(2) \( p_n^k \) is the superposition of Poisson distributions for the particular case of the Poisson transform\(^4,5\) of the weight \( f(x) \)

\[ p_n = \int_0^\infty dx f(x) \frac{(\bar{n})^n e^{-\bar{n}}}{n!} \]  

(3)
with the weight \( f(x) = \psi_k(x) \) given by Eq. (2). This formula suggests the possibility \( 6-8 \) that the observed broad hadronic distributions are the consequence of an average (in the event sample) over varying inelasticities or equivalently impact parameters, for continuous \( k \).

(3) \( P_n^k \) corresponds to the counting distribution characteristic of a Gaussian field ensemble, whether in semiclassical photocount theory \( 2,9 \) or in the representation of the oscillator by the diagonal coherent state representation \( 9,10 \). For applications to hadronization, \( k \) would be the average number of effective cells (or emitters). Although there is no fundamental basis for \( k \), most people imagine that the number of emitters should increase with energy.

(4) \( P_n^k \) can be derived as a composite Poisson-logarithmic distribution \( 11-13 \) in which clusters produced with a Poisson distribution decay into the final hadrons via a logarithmic distribution. This picture (as do the others mentioned here) requires further elaboration to become compelling.

(5) \( P_n^k \) is the solution of various \( 3,4,14-17 \) probability evolution equations in the parameter \( \bar{n} \). In these cases the mathematics is more clear than the physical processes allowing the reduction of the many-body problem to a few degrees of freedom obeying the appropriate equations.

(6) \( P_n^k \) is the long time distribution whose time dependent Poisson kernel \( f(x,t) \) obeys a suitable stochastic differential equation. At least two cases are known \( 6,18 \).

(7) \( P_n^k \) can result from a Cantor set type of cascade structure (including parallel or composite cascades) in which \( P_n^k \) is the fraction of a line occupied at the \( n \)th stage \( 19 \). This interpretation is closely connected with the interpretation of the \( k=1 \) Bose-Einstein distribution as a "geometric" distribution \( 3 \).

The foregoing list in no way exhausts the rich variety of contexts in which the negative binomial distribution occurs in nature \( 20 \). We have emphasized those which may have relevance to the particle physics multihadron production problem.

II. PROTOTYPE HADRONIZATION DENSITY MATRICES

The counting distribution \( P_n \) can in principle be obtained from the projection of the wave function \( \psi(t) \) on the \( n \) particle states at \( t=0 \), i.e. when the produced hadrons have separated. Writing \( \psi(\omega) \) as \( \psi_{\text{out}} \) we have
for a so called "pure" state. The "in" density matrix $\rho_{in} = |\psi_{in}><\psi_{in}|$ is related to $\rho_{out}$ via the $S$ "matrix" $S = |\psi_{in}><\psi_{out}|$ by $\rho_{out} = S\rho_{in}S^*$, indicating the relation of (4) to the usual formulation in terms of phase space integrations over the squared $S$-matrix. Having said this, we admit that a dynamical evaluation of $\rho_{out}$ is not easier than that of $S$. Nevertheless, one can make educated guesses on the structure of $\rho_{out}$ on the experimental results and accumulated from statistical physics, particularly the sophisticated results from quantum optics$^9$-$^{10}$. Such results can then provide well-formulated goals for more ambitious dynamical schemes such as jet calculus$^{21}$, dual topological models$^{22}$, etc. This framework suggests the merit of deriving the "stochastic essence" from the full exclusive event by searching for suitable probabilistic equations for inclusive variables. Although traditional in other branches of science, particle physics has heretofore made little use of these techniques.

The final hadrons are to good approximation described by a set of free Bose fields, whose creation and destruction operators ($a,a^+$) are nothing but free harmonic oscillators. $\rho_{out}$ is therefore some function of the outgoing $a$'s and $a^+$'s. Although the actual $a$ and $a^+$ variables are equipped with momenta and other degrees of freedom, it turns out to be fruitful to consider a prototype model with one (or a few) effective oscillators. The most popular oscillator states, the number states $|n> = (a^+)^n|0>/(n!)^{1/2}$ have ill-defined phase and hence are only indirectly related to classical-like field motions. This does not matter for the incoherent thermal ensemble, whose (mixed) density matrix

$$
\rho = \frac{\exp(-\beta H)}{\text{Tr} \exp(-\beta H)} = \sum_{n=0}^{\infty} \frac{N^n}{(1+N)^{n+1}} |n><n|
$$

is diagonal in the number basis. Recall that in this case the occupation number $N$ of the Bose-Einstein distribution is given by $N^{-1} = \exp(\beta\omega)-1$. 

\[\begin{align*}
\rho_{out} &= |\psi_{out}><\psi_{out}| \\
\rho_{in} &= |\psi_{in}><\psi_{in}| \\
S &= |\psi_{in}><\psi_{out}| \\
\rho_{out} &= S\rho_{in}S^* \\
\end{align*}\]
Suppose, however, we have the opposite case of an oscillator undergoing sinusoidal motion. In this case we expect a Poisson distribution for probabilities of the system being found in the nth excited state. As is well-known\(^4,9,10\), the most suitable states in this case are the coherent states \(|\alpha\rangle\) which can be defined (for any complex \(\alpha\)) by

\[
|\alpha\rangle = e^{-\frac{1}{2}||\alpha||^2} \sum_{n} \frac{\alpha^n}{n(n!)^{1/2}} |n\rangle
\]

The Poisson distribution follows immediately on identifying the mean multiplicity \(S\) with \(\langle |\alpha||^2\)

\[
\rho = |\langle \alpha|\alpha \rangle|
\]

\[
P_n = \frac{S^n e^{-S}}{n!}, \quad S = |\alpha|^2
\]

The motion \(\langle \alpha|x(t)|\alpha \rangle \sim |\alpha|^2 \cos(\phi - \omega t)\) where \(\phi = \arg \alpha\). Although \(\langle \alpha|\alpha \rangle\) has many off-diagonal elements in the number basis, the counting process is not sensitive to them. Hence observation of (7) in no way implies that the actual physical system has the full classical-like phase structure.

Next suppose that instead of the pure state \(|\alpha|\alpha \rangle\) we have a mixed ensemble with real weight function \(\Phi(\alpha)\)

\[
\rho = \int d^2 \alpha \Phi(\alpha) |\alpha|\alpha \rangle\langle \alpha|
\]

This representation has considerable great generality than might be surmised from the foregoing. Moreover, as one can easily see the diagonal element \(\langle n|\rho|n \rangle\) leads directly to the Poisson transform, Eq. (3), which thereby inherits this greater generality.

As our first example we note that a Gaussian weight leads to the Bose-Einstein distribution:

\[
\Phi(\alpha) = \exp(-|\alpha|^2/N)/\pi N
\]

\[
P_n = \frac{N^n}{(1+N)^{n+1}}
\]
(For k modes the direct product \(\exp(-\Sigma|\alpha_i|^2/(N/k))/(\pi N/k)^k\) leads directly to the negative binomial, Eq. (1). Note that these results are compatible with but do not require thermal equilibrium.

A very interesting generalization, which actually arises in a model of a single-mode laser\(^9\), is the displaced Gaussian weight whose \(P_n\) interpolates between (7) and (9)

\[
\Phi = \exp(-|\alpha-\beta|^2/N) / \pi N
\]

\[
P_n = \frac{N^n}{(1+N)^{n+1}} \exp\left(-\frac{S}{1+N}\right) L_n\left(-\frac{S}{1+N}\right)
\]

Here the average multiplicity is \(\bar{n} = S+N\) with \(S = |\beta|^2\); \(L_n\) is the usual Laguerre polynomial (positive for negative argument.)

The notation is chosen so that \(S\) measures the strength of the coherent signal and \(N\) the strength of the (Gaussian) noise. We shall refer to (10) and its generalization to \(k\) (equal strength) cells:

\[
P_{nk} = \frac{(N/k)^n}{(1+N/k)^{n+k}} \exp\left(-\frac{S}{1+N/k}\right) L_{n-k}^{-1}\left(-kS/N\right)
\]

as the partially coherent laser distribution (PCLD). These formulas were originally derived for the photocount distribution for \(k\)-mode lasers whose emitting modes have a (common) signal and noise ratio as defined above. This suggests application to the description of hadron counting for emissions from a set of cells having to first approximation a common \((S,N)\).

III. DESCRIPTION OF MULTIPLICITY DISTRIBUTIONS

The PCLD Eq. (11) was first used\(^5,24\) to describe hadronic multiplicities in 1983, although equivalent physics was assumed for moments (for finite rapidity differences) as early as 1978 by the Marburg\(^25\) group. We note that (11) depends on three parameters \((S,N,k)\) as opposed to just two for the negative binomial, Eq. (1). We shall use the equivalent set \((\bar{n} = N+S, m = (N/S)^{\frac{k}{2}}, k)\). Note that as \(N/S \to \infty\) (11) goes over to the negative binomial, while \(N/S \to 0\) leads to the Poisson. We have used the noise to signal amplitude \(m = (N/S)^{\frac{k}{2}}\) rather than the
usual S/N ratio for the following reason. Near the Poisson limit the
shape of the wings of the distribution is very sensitive to a small
amount of noise. Hence m can be a few percent, in some sense very close
to Poissonian, yet to the eye the curve looks quite different from
Poissonian (see Fig. 1 of ref. 5 for illustration of this fact).

Since (11) has an extra parameter, it is not surprising that it
leads to multiplicity fits as good or better than the negative binomial.
What is not visible at first sight, however, is that one can trade an
increasing k for an increasing S/N: either will narrow the distribution
in a way which accommodates the data equally well from a $\chi^2$
criterion. Examples were given\textsuperscript{26} by us at the Lund conference; more will be given
elsewhere. Due to space limitations we here only assert again the
result: fits to multiplicity distributions alone cannot distinguish
between the negative binomial from the partially coherent distribution
with smaller k and non-zero S/N. However, analysis of the pp
forward-backward correlation\textsuperscript{28} does constrain the parameters, indicating
that N/S > 1 but not that (1) is really in force. What are the
consequences of this ambiguity in the parameter space of (11)? The most
important ones are:

(1) The large values of k obtained by UA5 by fitting Eq. (1) to
non-single-diffractive data can be eliminated by allowing coherent
emission to be substantial at low energy, disappearing completely at
higher energies (so that the energy dependence of KNO plot could
stabilize at higher energies.) This point of view, stressed recently by
the Marburg group\textsuperscript{29}, shows how the puzzling rapid decrease of k can be
replaced by a more plausible increase of randomness of the emitting
fields with increasing energy.

(2) Recent measurements of $e^+e^-$ annihilation to hadrons at 29 GeV
have given precise charged multiplicity distributions. These data were
very well described by (1) with very large values of k (ranging up to
100). In 1984 we claimed that existing data were almost Poissonian, the
deviations being due to a small amount of noise superposed on almost
coherent-state behavior for one or two quark jets (i.e. k = 1 or 2 is
literally the number of sources). Since for $k \to \infty$ the negative binomial
approaches the Poisson, these views are not very different
mathematically even though we have no idea how to interpret k = 50, or
even 20, in a physical way. Recently\textsuperscript{31} we have analyzed the data of ref. 30 to try to discriminate phenomenologically between these alternatives. It is very hard to distinguish, even with the aid of F/B correlation data and restricted rapidity intervals, although the nearly Poissonian limit looks somewhat better.

To summarize, the replacement of the negative binomial distribution (1) by the partially coherent distribution for both hadron-hadron and e\textsuperscript{+}-e\textsuperscript{-} multiplicity distributions. In each case the physical picture is intuitively simple. For h-h we have emissions from a small average number of effective cells (whose number could even be constant). The energy dependence of the KNO plot is then to be interpreted as the decrease of S/N; current collider results are nearly at the negative binomial limit. It is therefore tempting to speculate that the C\textsubscript{n} moments will saturate beginning by Fermilab collider energies. For e\textsuperscript{+}-e\textsuperscript{-} hadronizations we can, as in ref. 5 continue to identify k as the number of jets (except at the lower energies), which is small. The narrowness of the distribution is due to the largeness of S/N. What is missing in this parametrization by (11) is any understanding of why KNO scaling should hold (as it seems to) in e\textsuperscript{+}-e\textsuperscript{-} annihilations. The pure Poisson does not scale, and we have neither a dynamical or statistical explanation of how N/S should be tuned to conform to the apparent experimental validity of KNO scaling in e\textsuperscript{+}-e\textsuperscript{-} annihilations. We hope that the next generation of experiments will shed light on this question.

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27. C. C. Shih (unpublished, 1984); to appear in ref. 4.


ABSTRACT

The shape and the energy-variation of multiplicity distributions after the jet-component is removed agree with the predictions of the branching model whose dynamical contents and other consequences are discussed. Other phenomenological aspects of multiplicity distributions are also discussed, including the dependence on the rapidity window and the necessity for two components.
I. Phenomenology

It is a remarkable fact that the probability \( P_n \) for producing \( n \) charged particles at a given c.m. energy \( \sqrt{s} \) with an average multiplicity \( \bar{n}(s) \) is given by the negative binomial distribution\(^1\)

\[
P_n(k, \bar{n}) = \frac{(n+k-1)!}{n!(k-1)!} (\frac{\bar{n}}{k})^{-k} \frac{1}{(1 + \frac{k}{\bar{n}})^{n+k}}
\]

predicted by some theories\(^2\). If \( k \) were a constant, it would have predicted also the energy variation, but unfortunately this is not the case. Instead, it can be fitted by the following formula for all energies\(^1\):

\[
k^{-1} = -0.104 + 0.058 \ln \sqrt{s}
\]

The width of (1) is given by

\[
\gamma_2 = D^2/\bar{n} = k^{-1} + n^{-1},
\]

where \( D = (n - \bar{n})^{1/2} \) is the dispersion.

Now there is another theory which predicts a different negative binomial distribution

\[
P'_n(k', \bar{n}) = \frac{(n-1)!}{(n-k')! (k'-1)!} (\frac{\bar{n}}{k'})^{-k'} (1 - \frac{k'}{\bar{n}})^{-n-k'}
\]

whose width is given by

\[
\gamma_2 = D'/\bar{n} = k'^{-1} - \bar{n}^{-1}
\]

This is the branching model\(^7-9\) and the purpose of this talk is to discuss this model. Before doing so let us first examine the phenomenology of (4). Detailed fits using (4) are not available, but its resemblance with (1) makes it quite likely that it can also fit the shape of the distribution. Even without fits, we can estimate \( k' \) by equating (3) and (5). Then we get

\[
k'^{-1} = k^{-1} + 2/\bar{n}
\]

Using (2), (6) and the measured values of \( \bar{n} \), we have computed \( k' \) and find it to be energy independent! For example, for \( \sqrt{s}=22(\text{NA22}) \), 62.2(\text{ISR}), and 200(\text{UA5}) GeV, \( \bar{n}=8.76, 13.8, \) and 21.6; \( k=12.4, 7.4, \) and 4.6, leading to a value of \( k' \) equal to 3.2, 3.6, and 3.2 respectively. A more thorough analysis of this will appear in a forthcoming paper\(^3\). It turns out that \( k' \) is consistent with being a constant at all energies up to about \( \sqrt{s}=200 \) GeV, beyond which it may have a slow decrease. This is consistent with the \text{UA1}\(^4\) observation that KNO scaling can be restored after the removal of the jet component from the data. Without its removal, the mismatch between
the jet and non-jet components would widen the distribution and lead to a smaller \( k' \). Since there is very little jet component below \( \sqrt{s} \)pp energies, jet removal is important only at the super-collider energies.

There is then a good case for refitting the data after jet-removal with (4) and we hope that this will be done soon. If the constancy of \( k' \) is upheld by such a detailed analysis, then it is truly remarkable because now (4) can explain not only the shape but also the energy variation of the multiplicity distribution.

The fact that we have to separate the soft (non-jet) and the hard (jet) components in order to have a simple distribution (4) argues for two different dynamical mechanisms for the two components. What are they? We suggest that the soft component involves little color separation and interaction occurs directly between color-neutral clusters in such a way that it is governed mostly by phase space considerations (see next two sections). In contrast, the hard component is driven by color interaction, or QCD, which allows many soft gluons to be produced. This accounts for the observation (4) that the hard component has a large \( \bar{n}(s) \), and its multiplicity distribution is far narrower than that for the soft component. This latter is because most of the (soft) gluons do not have sufficient energy to branch out again. If none of them can branch out again, then the factor \( (n-L) \) on the right-hand side of (6) would be replaced by unity and the solution is the (narrow) Poisson distribution. Since we do have hard gluons which can branch out, the distribution is likely to be wider than Poisson but is still understandably narrow.

Another phenomenological aspect that we want to remark on concerns the observed dependence
of $\gamma_2$ on the pseudo-rapidity interval $\eta_0 (|\eta| < \eta_0)$ in which the data is taken: $\gamma_2$ increases with decreasing $\eta_0$. Qualitatively this can be understood as a result of energy-momentum conservation as follows. Most of the particles in an event with large $n$ must have small $\eta$ in order to conserve energy. Thus small-$\eta_0$ distributions are biased towards large $n$ and are thus wider (larger $\gamma_2$). In Fig. 1, a numerical result based on this idea is shown.

The physical basis of (4) will be discussed in the next two sections.

II. The Branching Equation

It can be shown that if $\phi(z)=\bar{n}_n$ obeys KNO scaling at very high energies, then $P_n$ would satisfy the following evolution equation

$$\frac{dP_n}{dt} = \sum_{k} a_k \{ (n-k)P_{n-k} - nP_n \}$$

(7)

at $n \gg 1$ and $\bar{n} \gg 1$. The evolution parameter $t$ is an increasing function of $s$ or $\bar{n}$, $a_k$ are some parameters and $k$ can be any real number. In particular, if $k \geq 1$ are positive integers and if (7) is valid for all $n$, then we will refer to it as the branching equation. To justify this name and explain its physical content, consider the simplest case $k=1$, and normalize $t$ so that $a_1 = 1$. Then it follows from (7) that $t=\bar{n}_n \bar{n}(s)$, and the equation describes production of particles by branching, or division, much like the biological cells do. Branching can happen to any of the particles present, and the probability for that to occur is proportional to $dt$. This is the content of (7) for $k=1$ and similar interpretations can be given to the general case. The solution of (7) for $k=1$ with the initial condition $P_n = \bar{n}_n k'$ at $t=\bar{n}_k \bar{n}k'$ is given by (4) and (5), where $k'$ is the initial number of clusters produced.

III. Dynamical Branching Model

Division of the original $k'$ clusters will continue presumably until energy is exhausted at the analog time $t=\bar{n}_n \bar{n}(s)$, though it is not easy to justify this statement directly from (7). The difficulty arises because (7) is a probability equation that knows nothing about energy-momentum conservation. To prove the statement and to gain a dynamical understanding of the content of (7), a dynamical model fully incorporating energy-momentum conservation which leads to (7) has to be constructed. This has been done and (7) appears as a result of the model in the infinite energy limit. The model is controlled mostly by phase space though sufficient flexibility is built in to accommodate any observed $\bar{n}(s)$. The reason why (7) results at infinite energy is because the model is based on the branching tree.
diagrams, and that energy is so plentiful for large $s$ that energy conservation can largely be ignored. But this is no longer the case at lower energies and (7) has to be modified. The leading correction can be computed and the modified evolution equation becomes

$$\frac{dP_n}{dt} = (n-1) \left[1-f(n-1,t)\right]P_{n-1} - n[1-f(n,t)]P_n,$$

$$f(n,t) = 0.06 c^2 / [\xi^2(n)[t-\xi n b]^2],$$

if for simplicity we take $\bar{n}(s)=b(s')^{c/4}$. The function $\xi(n)$, shown in Fig. 2, is a slowly decreasing function of $n$. If this decrease is ignored, then $f(n,t) \approx f(t)$, and a change of variable from $dt$ to $d\xi=dt[1-f(t)]$ will restore (8) and (7) with $\xi=1$, so that (4) is still the solution. We have also solved (8) numerically with full $\xi(n)$ and shown that there is practically no difference between this solution and (4), at least in the $\bar{p}p$ energy range.

It is also possible to calculate the rapidity distribution of the produced particles from this dynamical model. Since the branching tree diagrams are s-channel like, one might expect the long chains in the diagram to correspond to jets. Calculation shows that this is not the case; these chains actually correspond to combs in t-channel diagrams, as far as the rapidity distribution is concerned. The reason is because the dynamics in the model is phase-space-like rather than QCD-like, so that the result resembles more t-channel (smooth, potential) scattering than s-channel resonances or jets.

IV. Conclusion

In conclusion, we have shown that the solution (4) of the branching equation can probably explain both the shape and the energy variation of the multiplicity distribution. We have also discussed the dependence of the distribution on the rapidity interval $\eta_0$, and have suggested a possible physical mechanism distinguishing the soft and the hard components of the
distribution. Dynamical models for the branching equations can be constructed, from which one obtains the branching equation in the infinite energy limit, and the modified branching equation (8) at finite energy. Rapidity distributions of produced particles can also be computed and the result is t-channel-like.

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MULTIPLICITY DISTRIBUTIONS AT HIGH ENERGIES AND THEIR INTERPRETATIONS

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Abstract: Ambiguities in presenting and interpreting the data on multiplicity distributions for non-diffractive high energy p(p)p interactions are indicated and discussed. It is shown that drawing any detailed conclusions from the successful fit and fitted parameter values may be very misleading. The distributions for limited (pseudo) rapidity intervals are also discussed.

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In this talk few remarks will be presented on the multiplicity distributions and their interpretations.

Let us start by discussing the normalized moments of multiplicity distributions for full phase-space. The compilation of non-diffractive pp and $\bar{p}p$ data for $\sqrt{s}$ above 10 GeV presented by the UA5 group\(^1\) and supplemented by 200 and 900 GeV data\(^2\) shows apparently that the normalized moments $c_q = \frac{n}{n^q}$ are constant for $\sqrt{s} < 60$ GeV, but increase significantly at higher energies. On the other hand, $l/k$ parameter of the successful fits to the negative binomial distributions

$$P(n) = \frac{\Gamma(n+k)}{n\Gamma(k)} \left( \frac{\bar{n}}{\bar{n}+k} \right)^n \left( \frac{k}{\bar{n}+k} \right)^k$$

seems to increase linearly with ln $s$ in all the energy range considered, suggesting that the KNO scaling\(^3\) at lower energies was accidental. We shall note that the $l/k$ parameter determines (independently of $\bar{n}$) the normalized factorial moments $f_q = \frac{n(n-1)\ldots(n-q+1)}{n^q}$, which are just the moments expected originally to be constant (up to the $1/\ln s$ terms) in the KNO analysis\(^3\). There is really no preference between choosing $c_q$ or $f_q$ as the parameters showing the energy (in)dependence of multiplicity distributions. Moreover, in the compilation\(^1\) one counted all charged hadrons of the final state. One can argue that a better choice would be to subtract two charged hadrons already present in the initial state, and/or to count only every second charged hadron to account for charge conservation. All these choices lead to the same asymptotic values of normalized moments, but in the accessible energy range the results vary significantly. The energy dependence of normalized moments shown in fig. 1 is quite different for various choices. The only common feature is an increase above the top ISR energy and it is a matter of personal choice if we say that this trend is new, the same as in lower energy data, or reversed. Thus it is clear that the multiplicity distribution alone does not allow for any firm conclusions concerning the energy dependence of multiple production.

Fig. 1. Normalized moments $c_q$ for $n=n_{ch}$ (black squares) and $n=n_{ch}-2$ (black dots with error-bars) and normalized factorial moments $f_q$ for $n=n_{ch}$ (open squares), $n=n_{ch}-2$ (open circles) and $n=n_{ch}/2-1$ (crosses) as functions of CM energy $\sqrt{s}$. 
The second remark concerns the interpretation of the successful fit of data to the negative binomial distribution (denoted further as NBD). Let us note first that the standard arguments against the original (Planck's) interpretation of NBD (as a superposition of k sources, each emitting particles according to the geometrical distribution) are that fitted k values are non-integer and decrease with energy \(^1\)). However, it is easy to check that for each \(k \gg 2\) the distribution fitted as NBD with non-integer \(k\) may be fitted as well (with first six normalized moments differing by less than 1 \%) as a superposition of two NBD-s with integer \(k_1, k_2\) values bracketing \(k\). Moreover, the broadening of distribution (decreasing \(k\)) may be to a large extent explained by growing phase-space \(^4\) and/or by wide fluctuations of energy attributed to different sources. Thus the original interpretation may well be the best one.

However, attaching too much significance to the fitted values of \(k\) and other resulting parameters seems to be unjustified. We may see it clearly in the example of a model \(^5\) in which NBD results from the Poissonian distribution (with average \(\bar{c}\)) in the number of clusters, each decaying into hadrons according to the logarithmic distribution

\[
p(h) = \frac{k}{\bar{c}h} \left(\frac{\bar{c}h}{\bar{c}+k}\right)^h
\]

Assuming at least one charged hadron in each cluster one finds

\[
\bar{c} = k \ln(1+n/k)
\]

from which one obtains the values of \(\bar{c}\) surprisingly constant at a level of 8 for 504\(s<1000\) GeV. However, it has been shown \(^6\) that assuming instead at least one hadron of any charge in each cluster and attributing randomly zero or non-zero charge to each hadron (with probabilities 1/3 and 2/3, respectively) one finds instead

\[
\bar{c} = k \ln(1+3n/2k)
\]

which yields obviously quite different values of \(\bar{c}\), as shown in fig.2. This assumption may seem more plausible, as it leads also to a quite succesful prediction for correlations between the numbers of charged and neutral hadrons \(^6\)). Nevertheless, the values of \(\bar{c}\) are no more uniquely determined than before. Assuming in turn NBD for pairs of charged hadrons (which looks reasonable if one wants to take into account charge conservation, and which gives equally good fits) one finds a new value \(k'\) given approximately by \(1/k'=(1/k-1/n\), and if each cluster contains at least one pair, one gets
\[ c = k' \ln(1 + n/2k') \]  

resulting again in different values, also shown in fig. 2. Further variations can be introduced by subtracting two initial hadrons etc.

Moreover, let us note that five of the lower energy points correspond now to negative \( k' \). This makes formula (2) obviously absurd (negative probabilities), but \( c \) behaves there as smoothly as for positive \( k \). Thus it seems unlikely that the values of \( c \) extracted by one of possible ways from data may have any physical meaning. One may add more generally that well known abundance of resonances in the final state makes any straightforward interpretation of stable hadrons' multiplicity distribution at finite energies rather unreliable.

It has been suggested that more information can be extracted from an analysis of multiplicity distributions in restricted (pseudo) rapidity intervals. At 540 GeV such an analysis resulted in good fit to NBD with \( k \) values increasingly (linearly?) with interval length\(^7\). One can interpret this result in terms of cluster size effects\(^5\), varying number of contributing sources etc., but the effect is already quite well described\(^8\) in a minimal model - an uncorrelated cluster emission model, in which the multiplicity distributions for the full phase-space are constrained to fit the data. In this model "clusters" are defined as objects needed to describe short-range rapidity correlation seen in the data and may be regarded as "effective resonances", possibly including local charge conservation effects. Thus, unlike the clusters of previously discussed model, they have well known decay parameters and little freedom is left for adjustments. We refer to\(^8\) for more detailed discussion of the model and its comparison with data. Let us mention here only that a small but non-negligible discrepancy between the experiment and model predictions disappears if one modifies the original assumption of a common shape of all semi-inclusive rapidity distributions to account for narrower distributions at higher multiplicities, as indicated by data. This effect may seem to result trivially from the energy-momentum conservation, but the data suggest that it it stronger at collider\(^9\) than ISR\(^10\). (Note that the SPS fixed target data presented at
this meeting also seem to agree quite well with the minimal model without modification.) Thus the analysisresults in two conclusions: 1) a new high multiplicity component with a narrow rapidity distribution seems to appear at collider energies (as suggested already by a multitude of models and data analyses), and ii) apart from this effect, the multiplicity distributions in rapidity intervals seem to follow just from the shape of the full multiplicity distribution and the uncorrelated emission of cluster for fixed global multiplicity. Therefore, no extra dynamical assumptions are needed to describe these data and they cannot be regarded as strong evidence in favour or against any model. It is only relevant if the model takes into account the existence of clusters, describes properly full multiplicity distributions and has structure approximately compatible with i) and ii).

To conclude, we have shown that the presentation and interpretation of data on multiplicity distributions is highly ambiguous. It seems very unlikely that one can build a good model for multiple production starting just from the regularities discovered in these data. It is certainly much more important to test directly if any model fits the data well, than to inquire if it produces "naturally" negative binomial distributions, and if it provides a "physical" reason for the observed energy- and interval length dependence of k or other fit parameters. Obviously, this does not diminish the value of NBD fits as an excellent way of summarizing and memorizing the main features of data. However, it remains to be seen if these fits are as deeply significant as they are useful.

References:

ABSTRACT

The UA5 experiment has measured the ratio of the inelastic $\bar{p}p$ cross-sections at 900 and 200 GeV to be $1.20 \pm 0.01 \pm 0.02$. This result is compared to fits to data on total cross-sections and on the ratio of the real to imaginary part of the forward scattering amplitude. These fits, using analytic functions for the forward scattering amplitude, give different predictions for the measured ratio for the two cases $\sigma_L$ - constant or $\ln^2 s$, and the UA5 result favours the case with a constant cross-section at very high energy. Comparisons are also made with the highest energy cosmic ray results.
1. INTRODUCTION

The pp and $\bar{p}p$ total cross-sections $\sigma_t$ were found to increase over the c.m. energy range $\sqrt{s} = 15-62$ GeV (ISR energies) with a few mb [1]. Recently $\sigma_t$ was measured at $\sqrt{s} = 546$ GeV by the UA1 and UA4 experiments [2,3], showing an increase over the values obtained at the ISR by some 20 mb. The increase was found to be consistent with a $ln^2 s$ dependence of $\sigma_t$. Recently much attention has been given to the asymptotic form of the total cross-section [4-6]. Does $\sigma_t$ continue to increase at the maximum rate allowed by the Froissart bound [7], $\sigma_t \propto ln^2 s$, or does $\sigma_t$ increase more slowly or attain a constant value?

Block and Cahn [4,6] and Bourrely and Martin [5] have used an even scattering amplitude of the form

$$F(s) = \frac{A + B[ln s/s_0 - i\pi/2]^2}{1 + C[ln s/s_0 - i\pi/2]^2}$$

(1)

to describe the total cross-section. If $C = 0$ then $\sigma_t \propto ln^2 s$ but if $C \neq 0$ then $\sigma_t$ is constant. In this paper we compare the recent measurement by the UA5 experiment [8] of the inelastic cross-section ratio

$$R \equiv \frac{\sigma_{inel}(900 \text{ GeV})}{\sigma_{inel}(200 \text{ GeV})}$$

(2)

with predictions from fits with either $C = 0$ or $C \neq 0$ in (1) to different sets of data on $\sigma_t(pp)$, $\sigma_t(pp)$ and $\varphi$, the ratio of the real to imaginary part of the forward scattering amplitude. We also compare these fits to accelerator data with higher energy cosmic ray data [9-15].

2. THE UA5 CROSS-SECTION EXPERIMENT

The experimental method is described in ref. [8] and the reader is referred to that paper for details. The CERN $pp$ Collider was operated in March-April 1985 in a pulsed mode, with the proton and antiproton beams circulating at a flat top of 4 s length at 450 GeV/c and at a flat bottom of 8 s length at 100 GeV/c. The total cycle length was 21 s. The luminosity of the machine was much lower than during normal dc operation but gave adequate trigger rates for the measurement of the inelastic rates. The basic principle of the UA5 cross-section experiment was that although the absolute luminosity of the Collider was known only to 10-15 %, the ratio of the luminosities at $\sqrt{s} = 900$ and 200 GeV was known to about 1 % [8]. Thus a measurement of the ratio of the corresponding interaction rates allows, in principle, an accurate determination of R in (2).
The UA5 detector and standard analysis procedures have been described elsewhere [16]. Two large streamer chambers, 6m×1.25m×0.5m, were placed above and below the SPS beryllium beam pipe. This gave a geometrical acceptance of about 95% for |η|<3, falling to zero at |η|=5 (η is the pseudorapidity, η=−ln tan θ/2 where θ is the c.m. emission angle). The trigger for the chambers was provided by scintillation counter hodoscopes at each end of the chambers covering the pseudorapidity range 2<|η|<5.6. The streamer chambers were triggered at the flat bottom (\(\sqrt{s} = 200\) GeV), at the flat top (\(\sqrt{s} = 900\) GeV) and also during the ramps. Two different triggers were used: the 2-arm trigger and the 1-arm trigger. During the 1 s deadtime of the streamer chambers, rates for the 1-arm and 2-arm triggers were collected. The data presented here are based on about 60k triggers, of which 25k had streamer chamber photographs that were examined.

The rates \(N_1\) and \(N_2\) for the 1-arm and 2-arm triggers are related to the luminosity \(L\) through \(N = L\sigma\). The relations between the physical processes (elastic scattering, single diffraction SD and non-single diffraction NSD) with the different triggers are illustrated in fig. 1. The trigger cross-sections \(\sigma_1\) and \(\sigma_2\) are linear functions of the physical cross-sections \(\sigma_{SD}\) and \(\sigma_{NSD}\) with constants, trigger efficiencies \(\varepsilon\), that are determined by simulations [8,17].

![Fig. 1](image-url) Relations between different physical processes and the different triggers. The figures refer to 900 GeV.

The relations are:

\[
\begin{pmatrix}
\sigma_1 \\
\sigma_2
\end{pmatrix} =
\begin{pmatrix}
\varepsilon_1^{SD} & \varepsilon_1^{NSD} \\
\varepsilon_2^{SD} & \varepsilon_2^{NSD}
\end{pmatrix}
\begin{pmatrix}
\sigma_{SD} \\
\sigma_{NSD}
\end{pmatrix}
\]
Here \( \epsilon_{1\text{ SD}} \) is the 1-arm triggering efficiency for single diffractive events etc. Given measured rates \( N_1' \), the relevant cross-sections can be determined through eq. (3).

A very important part of the analysis work was to correct the measured rates for background caused by, e.g. beam-gas interactions. These corrections were made using the streamer chamber photographs. For the 2-arm triggers, the fraction of beam-beam interactions was found to very high, > 90%. For the 1-arm triggers, however, this fraction was much lower being about 8% and 27% for 200 and 900 GeV, respectively. The ratios of the trigger rates at 900 to those at 200 GeV was found to be constant during each run and also from run to run.

The final result for the ratio (2) of the inelastic cross-sections at 900 and 200 GeV is

\[
R_{\text{inel}} = 1.20 \pm 0.01 \pm 0.02
\]

The first error is statistical, the second systematic. The latter takes into account estimated errors in background subtraction, luminosity ratio and trigger efficiencies (\( \epsilon \) in eq. (3)) as determined by simulations. One should note that a large part of the systematic errors cancel in the ratio. For details the reader is referred to ref. [8].

3. VALUES OF \( \sigma_t(900) \) AND \( \sigma_{\text{inel}}(900) \)

In order to relate \( \sigma_t \) and \( \sigma_{\text{inel}} \) we need the ratio \( \sigma_{\text{el}}/\sigma_{\text{inel}} = (\sigma_t - \sigma_{\text{inel}})/\sigma_t \). This ratio has been measured at \( \sqrt{s} = 546 \) GeV to be \( 0.215 \pm 0.005 \) [3], whereas at ISR energies it is lower, around 0.18 [1,18]. We estimate the values at 200 and 900 GeV to be \( 0.19 \pm 0.01 \) and \( 0.23 \pm 0.01 \), respectively (See also fits to this ratio in e.g. ref. [4].) The relation between \( R_{\text{inel}} \) (eq. (3)) and \( R_t \), the corresponding ratio of total cross-sections, then becomes

\[
R_{\text{inel}} = (0.95 \pm 0.02)R_t
\]

Using the recent fit to \( \sigma_t \) by Amos et al. [1], which gives \( \sigma_t(200 \text{ GeV}) = 51.6 \pm 0.4 \) mb, and the above values of \( \sigma_{\text{el}}/\sigma_t \), we then get

\[
\begin{align*}
\sigma_t(900 \text{ GeV}) &= 65.3 \pm 0.7 \pm 1.5 \text{ mb} \\
\sigma_{\text{inel}}(900 \text{ GeV}) &= 50.3 \pm 0.4 \pm 1.0 \text{ mb}
\end{align*}
\]

The first error is statistical, the second systematic. The values obtained for \( \sigma_t(900) \) agrees with the dispersion relation fit of Amos et al. [1] that gives \( \sigma_t(900) = 65.8 \) mb. Fig. 2 shows the energy dependence of the total cross-section. The line is a fit using dispersion relations taken from Amos et al. [1]. Our value at 900 GeV falls nicely on the curve.
Energy dependence of the total cross-sections for pp and $\bar{p}p$. The figure is from Amos et al. [1] and the UA5 point at 900 GeV has been added. The line is a dispersion relation fit [1].

![Energy dependence graph]

4. COMPARISON WITH FITS TO DATA USING ANALYTIC FUNCTIONS

Fits to $\sigma_t$ and $\phi$ for pp and $\bar{p}p$ using an even scattering amplitude of the type (1) has been done by Block and Cahn [4] using data in the c. m. energy range 5 - 63 GeV and by Bourrely and Martin [5] using ISR energy data and a UA4 $\sigma_t(1 + \phi^2)$ value at $\sqrt{s} = 546$ GeV [3]. Using their results and the relation $R_{\text{inel}} = (0.95 \pm 0.02)R_t$ we get the predictions for $R_{\text{inel}}$ shown in table 1.

Table 1. Predicted values of $R_{\text{inel}}$ compared to the UA5 measurement.

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<th>C = 0</th>
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<tr>
<td></td>
<td>$\sigma$ \to\ constant</td>
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<tr>
<td>Measured UA5</td>
<td>1.20$\pm$0.01$\pm$0.02</td>
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Fig. 2.

Block and Cahn [4] and Bourrely and Martin [5] using ISR energy data and a UA4 $\sigma_t(1 + \phi^2)$ value at $\sqrt{s} = 546$ GeV [3]. Using their results and the relation $R_{\text{inel}} = (0.95 \pm 0.02)R_t$ we get the predictions for $R_{\text{inel}}$ shown in table 1.

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From table 1 we conclude that using (1) as a form for the scattering amplitude the UA5 measurement favours the case $C \neq 0$, i.e. $\sigma_t \propto \ln^2 s$. Block and Cahn have updated their fit, and at this meeting results were presented [3] in agreement with the above conclusion.

5. COMPARISON WITH COSMIC RAY DATA

The highest energy cosmic ray data extend the available energy range up to about 100 TeV, i.e. two orders of magnitude higher than the present maximum at accelerators. Although cosmic ray data have generally very large errors, a comparison can still shed light on the very high energy behavior of $\sigma_t$. The proton-air inelastic cross-section is determined from a study of the depth distribution of the maximum for air showers. The value for $\sigma_{t}^{(pp)}$ is then derived using Glauber theory [19]. Fig. 3 shows a recent compilation of data on $\sigma_{t}^{(pp)}$ taken from Lindsley [9]. The two lines are fits to data on $\sigma_t$ and $\phi$ for pp and pp. Included were data in the energy range 19 - 26 GeV from Carrol et al. [1], 23 - 62 GeV from Amos et al. [1] and the UA4 value at 546 GeV [3]. The fits were performed using the form (1) for the scattering amplitude, either letting $C$ be a free parameter or fixing $C = 0$.

![Fig. 3](image-url)

A compilation of cosmic ray data on $\sigma_t$ for pp interactions, taken from Lindsley [9]. The curves are fits to accelerator data discussed in the text. The lower one corresponds to $C \neq 0$, i.e. $\sigma \propto \ln^2 s$ and the upper one to $C = 0$, i.e. $\sigma \propto \ln^2 s$. 
We conclude that cosmic ray data fits either solution equally well, although a given experiment, e.g. 'Akeno', may well favour one solution.

6. CONCLUSION

The UA5 cross-section experiment has given the following results:

\[
\begin{align*}
\sigma_{\text{inel}}(900 \text{ GeV}) & = 50.3 \pm 0.4 \pm 1.0 \text{ mb} \\
\sigma_t(900 \text{ GeV}) & = 65.3 \pm 0.7 \pm 1.5 \text{ mb} \\
R & = \frac{\sigma_{\text{inel}}(900 \text{ GeV})}{\sigma_{\text{inel}}(200 \text{ GeV})} = 1.20 \pm 0.01 \pm 0.02
\end{align*}
\]

Compared to predictions from fits of data on \(\sigma_t\) and \(\phi\), using analytic functions of the form (1), the result on \(R_{\text{inel}}\) favours the solution with \(C' = 0\), i.e. a total cross-section that attains a constant value at very high energy. The highest energy cosmic ray results cannot distinguish between the two solutions.

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LARGE- \( t \) ELASTIC SCATTERING AND DIFFRACTION DISSOCIATION AT THE CERN SPS COLLIDER

UA4 Collaboration

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ABSTRACT

Proton-antiproton elastic scattering was measured at the centre-of-mass energies of 546 and 630 GeV in the four-momentum transfer range \( 0.45 \leq -t \leq 2.2 \text{ GeV}^2 \). The \( t \)-distribution shows an exponential decrease with a pronounced break at \( -t \approx 0.9 \text{ GeV}^2 \) followed by a more gentle decrease. Diffraction dissociation \( \bar{p}p \rightarrow \bar{p}X \) was also studied. The pseudorapidity distribution of charged particles, produced in the fragmentation of the system \( X \), is shown to be in striking agreement with the prediction of the Dual Parton Model.
We report recent measurements of proton-antiproton elastic scattering and diffraction
dissociation at a centre-of-mass energies of 546 and 630 GeV, performed at the CERN
SppS Collider by the UA4 Collaboration. Comparisons with theoretical expectations will
be made for both reactions.

1. ELASTIC SCATTERING

Elastic events were detected by high resolution wire chambers and scintillation
counters hodoscopes, placed in special movable sections of the vacuum pipe ("Roman
pots")2. Data were collected during several Collider runs at $\sqrt{s} = 546$ and 630 GeV. At the
lower energy the betatron function at the crossing point in the horizontal and vertical
plane were $\beta_h = 2$ m, $\beta_v = 1$ m for the low-$\beta$ mode and $\beta_h = 1.3$ m, $\beta_v = 0.65$ m for the
squeezed-$\beta$ mode. At $\sqrt{s} = 630$ GeV the corresponding figures were $\beta_h = 1$ m and $\beta_v = 0.5$ m.

With these optics our detectors covered the angular range $2.5 \leq \theta \leq 4.6$ mrad in the
low-$\beta$ mode and $2.8 \leq \theta \leq 4.6$ mrad in the squeezed-$\beta$ mode (the upper limits being due to
the shape of the machine quadrupoles placed just in front of the inner pots). These figures
correspond to a $t$-range of $0.45 \leq -t \leq 1.55$ GeV$^2$ at $\sqrt{s} = 546$ GeV that becomes
$0.7 \leq -t \leq 2.2$ GeV$^2$ at $\sqrt{s} = 630$ GeV.

The trigger was provided by the left-right coincidence of the trigger counters of two
opposite arms. It was a 6-fold coincidence for the low-$\beta$ and a 8-fold coincidence for the
squeezed-$\beta$ mode. The deflection through the machine quadrupoles ensured that particles
coming from the crossing region and accepted by the trigger had a momentum within about
10% of the beam momentum. Calibration and detection efficiency of the wire chambers
were studied using the data collected with the elastic trigger itself. In the track
reconstruction procedure a hit was required in at least five of the eight drift planes and in
at least one of two proportional planes in a telescope. Due to the redundancy in the number
of planes a fully efficient detection could be ensured over the telescope acceptance.

The analysis was carried out according to the following steps:

(a) A single track was demanded in each telescope of two opposite arms without tracks in
the other telescopes.

(b) Events were rejected if a hit was present in a system of scintillation counter
hodoscopes covering the pseudorapidity range $2.5 \leq \eta \leq 5.6$ on both sides of the
crossing region.
(c) A three-standard-deviation cut was applied to the distributions of the vertical coordinate $y_0$ of the proton and antiproton trajectories at the centre of the crossing region. A cut on the quantity $y_0$ is equivalent to a momentum analysis on both scattered particles: a three-standard-deviation cut corresponds to a selection of events where the momentum of both $p$ and $\bar{p}$ is within 2% of the beam momentum. The natural momentum spread of the circulating beam is only a few times $10^{-4}$.

(d) A three-standard-deviation cut was then applied to the collinearity distributions of the proton and antiproton trajectories in the vertical plane.

The uncertainty on the scattering angle is determined by the intrinsic angular spread of the beam. The observed width of the collinearity distributions provides a direct measurement of this quantity which agrees well with the value calculated from the known machine parameters. The r.m.s. value of the $t$-resolution was $\Delta t = 0.06 \sqrt{t}$ at $\sqrt{s} = 546$ GeV and $\Delta t = 0.08 \sqrt{t}$ at $\sqrt{s} = 630$ GeV.

The position and the angle of the beam at the crossing point were continuously monitored during the data taking and checked comparing the trend of the observed angular distributions for the two arms (up-left x down-right) and (down-left x up-right).

The geometrical acceptance of the detectors as a function of $t$ was carefully studied by varying the accepted interval of $\Theta_h$ (the horizontal component of the scattering angle) in order to ensure clearance inside the quadrupole vacuum chamber. The acceptance was calculated with a Monte-Carlo that takes into account the size, the angular spread and the angle of the incoming beams. It was checked that the $t$-dependence of the data remained the same when further reducing the accepted interval of $|\Theta_h|$.

The actual value of the scattering angle was calculated by taking the average of the observed production angles of the proton and antiproton, corrected for the beam angle.

The absolute normalization of the elastic differential cross sections at $\sqrt{s} = 546$ GeV was obtained by smoothly joining the data points to previous measurements at lower momentum transfer, which, in turn, were normalized to the optical point. An alternative method is based on the use of a luminosity monitor which consists of the left-right coincidence of scintillation counters covering the angular range from $0.4^\circ$ up to $1.5^\circ$, symmetrically on both sides of the crossing region. At $\sqrt{s} = 546$ GeV the monitor detected 86% of the inelastic non diffractive cross section. No elastic scattering data at low-$t$ being available at $\sqrt{s} = 630$ GeV, we have normalized the elastic rates measured at this energy to the elastic cross section at $\sqrt{s} = 546$ GeV using the ratio of the counting rates of the luminosity monitor as taken at the two energies. A 2% correction was applied to account for the energy dependence of the effective monitor cross section.
Results from the present experiment\(^5,6\) are shown in fig. 1 compared with pp data measured at ISR at \(\sqrt{s} = 53\) GeV\(^7\).

The data show an exponential decrease with momentum transfer up to \(-t \sim 0.9\) GeV\(^2\) followed by a shoulder with no sign of the dip, present in the ISR pp data. The cross section at the second maximum is more than one order of magnitude higher than at the ISR. In the momentum transfer region \(1.1 \leq -t \leq 2.2\) GeV\(^2\) covered at \(\sqrt{s} = 630\) GeV the data can be fitted using a simple exponential, \(d\sigma/dt = \exp(bt)\), with a slope \(b = 2.7 \pm 0.1\) GeV\(^{-2}\). As can be seen in fig. 1, the differential cross section at \(-t = 2\) GeV\(^2\) still remains larger than at the ISR. However, the trend of the Collider data is such that, as the momentum transfer increases, they seem to approach the ISR data.

Proton-antiproton elastic scattering in the energy region from the ISR to the SPS Collider, and beyond, has been recently discussed in several theoretical papers (ref. 10 for a recent review) where also definite predictions are made concerning the kinematic range covered in the present experiment.

In fig. 2 our data are compared with predictions from same current models of hadron elastic scattering. Two of these models from Donnachie and Landshoff DL\(^{11}\) and from Gauron, Leader and Nicolescu GLN\(^{12}\) account for the difference on shape between Collider and ISR data to the presence of an odd contribution to the elastic amplitude dominant at large-\(t\). These models predict also a different shape in \(\bar{p}p\) and pp scattering as observed at \(\sqrt{s} = 53\) GeV at the ISR\(^8,9\). In contrast with the two above quoted approaches, most of the existing models, based mainly on a geometrical picture of the hadron collision, as the one from Bourrely, Soffer and Wu BSW\(^{13}\) also presented in fig. 2, predict that \(\bar{p}p\) and pp elastic cross sections approach each other as the energy increase in conflict with the \(\sqrt{s} = 53\) GeV data.
2. DIFFRACTION DISSOCIATION

The reaction $\bar{p}p + \bar{p}X$ was studied at the centre-of-mass energy $\sqrt{s} = 546$ GeV. The antiproton scattered at very small angles was observed in the elastic detectors discussed above. Charged particles produced by the proton fragmentation were observed in the vertex detector consisting of the UA4 counters and drift chamber telescopes, and of the central detector of the UA2 Collaboration. The trigger required a particle through the "Roman pot" telescopes on the outgoing antiproton side, in coincidence with at least one charged particle in the pseudorapidity range $2.5 \leq \eta \leq 5.6$ in the proton hemisphere. Events were analyzed by requiring the presence of an interaction vertex, reconstructed by the vertex detector, and of a single track in the "Roman pot" telescopes. The antiproton momentum was determined from the deflection of its trajectory by the magnetic field of the machine quadrupoles. Defining the variable $x = p/p_0$, where $p$ is the measured antiproton momentum and $p_0$ the beam momentum, the mass $M$ of the system $X$ is given by $M^2 = (1 - x)s$.

Data were collected with the Collider running in the low-$\beta$ mode. The acceptance of the forward spectrometer covered the ranges $x \geq 0.9$ and $0.5 \leq -t \leq 1.5$ GeV$^2$. A momentum resolution of 0.6% (r.m.s. value), was determined from magnetic analysis of elastically scattered particle.

Inelastic diffractive processes can take place if the coherence condition is satisfied, i.e. if the longitudinal momentum transfer is smaller than the inverse of the strong interaction radius. As the centre-of-mass energy increases, this condition becomes satisfied for larger and larger masses. The validity of such classical argument at the energy of the Collider, with production of heavy diffractive states, was shown by the data reported earlier$^{14}$, where the typical quasi-elastic peak for $M^2/s \lesssim 0.04$ was clearly visible in the invariant differential cross section. At fixed $t$ and for $0.01 \leq M^2/s \leq 0.04$, the $1/M^2$ dependence of the mass distribution already seen in the FNAL and ISR energy range was shown to persist for diffractive states of masses up to more than 100 GeV.

Further information on the features of diffraction dissociation events can be obtained from the study of the produced particles$^{15}$. Pseudo-rapidity distributions of charged tracks from the fragmentation of diffractive states are shown in fig. 3 for several values of the diffractive mass $M$. An average mass value $\langle M \rangle$ was assigned to each mass interval taking into account the observed $1/M^2$ behaviour and the effect of the mass resolution. Clusters are centered at the expected value $\eta_X = \log(\sqrt{s}/M)$ and their width increases with $M$ in good agreement with the prediction of longitudinal phase space (multiperipheral-type) models and clearly disfavouring isotropic decay (fire-ball) models of the diffractively produced states.
The similarity of the dissociation of a diffractively produced system of mass $M$ to the hadronization resulting from collisions at c.m.s. energy $\sqrt{s} = M$, is further illustrated in fig. 4. Integration of the pseudorapidity distributions of fig. 3 yields the average multiplicity $\langle n_{ch} \rangle$ of charged secondaries produced in the fragmentation of the X-system. Data from the present experiment are plotted in fig. 4 at $\sqrt{s} = M$ together with a compilation\(^{16}\) of $\langle n_{ch} \rangle$ as a function of $\sqrt{s}$ for inelastic, non-diffractive, collisions. The growth of $\langle n_{ch} \rangle$ with $M$ closely follows the observed rise of the charged multiplicity with $\sqrt{s}$. However, a closer inspection of fig. 3 shows that, contrary to the case of inelastic pp collisions, the pseudorapidity distribution are not symmetric with respect to the center of mass of the cluster: the average multiplicity in the backward hemisphere (i.e. towards the non-excited $\bar{p}$) is bigger than in the forward one. This asymmetry is similar the one observed in the inelastic np collisions.

We shall see in next section as all these features are quantitatively reproduced in the Dual Parton Model (DPM).
2.1 Diffraction dissociation in the Dual Parton Model \(^{17}\)

In the DPM, high mass single diffractive dissociation is described by the two-chain diagram shown in fig. 5 \(^{18}\). These two chains or strings are stretched between valence constituents of the excited proton and sea constituents of the non-excited antiproton (which have to form a color singlet). Due to the \(X^{-1}\) behaviour of the sea quarks and gluons structure functions, the antiproton inclusive spectrum will have the characteristic \(1/(1 - x_p)\) behaviour, typical of a triple-Pomeron Regge approach, well verified by the data as described above. Moreover, at fixed \(X = 1 - x_p = M^2/s\) the two chains in the diagram of fig. 5 are identical to those appearing in the dominant (two-chain) component of inelastic \(\pi^0 p\) scattering at center-of-mass energy \(\sqrt{s} = M\) shown in fig. 6. In fact the two chains \(q_s - (qq)^p\) and \(q_s - (gq)^p\) appearing in fig. 5 are the same as those in fig. 6 except that the upper quark and antiquark are from the antiproton-sea. Therefore the momentum distribution functions are in general different in the two cases. However, it turns out that when fixing \(X\), i.e. when fixing the mass \(M\) of the diffractive final state, the two momentum distribution functions became identical. In fact in the DPM the joint momentum distribution of a sea quark-antiquark pair is obtained from soft physics arguments (dual Regge behaviour). It was shown \(^{19}\) that when the \(x\)-values of the two members of the pair are equal \((x_1 = x_2\) in fig. 5) one obtains the characteristic \(x_1^{-1}\) (for \(x_1 \rightarrow 0\)) behaviour for each of them (the same as in deep-inelastic structure functions). However, when \(x_1 \gg x_2\) one obtains a \(x_2^{-3/2}\) \((x_2^{-1/2})\) behaviour for the fastest (slowest) member of the sea pair. A momentum distribution function of the sea pair, consistent with the above results can be obtained as follows:

\[
\rho(x_1, x_2) = C \int dx X^{-1} \int dx \rho_q^\pi(x) \delta(xX - x_1) \delta((1 - x)X - x_2) \frac{C}{x_1 + x_2} \frac{1}{(x_1 x_2)^{1/2}}
\]

where \(\rho_q^\pi = x^{-1/2}(1 - x)^{-1/2}\) is the momentum distribution functions of a valence quark in a \(\pi^0\). The factor \(X^{-1} = (x_1 + x_2)^{-1}\) is typical of sea constituents (both in deep-inelastic structure function and in dual models \(^{21}\)) and, as mentioned above, gives rise to the standard \((1 - x_p)^{-1}\) form of the diffractive peak. The second term is just the quark momentum distribution functions in a \(\pi\). It is now obvious that, at fixed \(x_1 + x_2 = X\) (i.e. at fixed \(M\)), the momentum distribution function of the sea quark-antiquark pair is the same as the one of the valence quark-antiquark pair in a \(p\) (i.e. \(x_1 + x_2 = 1\)). Thus, at fixed \(M\), the two chains of the graph in fig. 5 are identical to the corresponding chains in a \(\pi^0 p\) collision (fig. 6) at center-of-mass energy \(\sqrt{s} = M\).
The computation of the rapidity distribution, $dN/dy(M^2)$, of the diffractively produced state of mass $M$ can therefore be performed using the standard convolution$^{19,20}$ between momentum distribution functions and fragmentation functions$^{21}$ for each of the two chains of fig. 5, and adding them together. The results are compared with the experimental data in fig. 3 where the rapidity has been converted into pseudorapidity using the approach in ref.$^{22}$. In order to reproduce the data in the first mass interval ($M < 50 \text{ GeV}$), one has to integrate $dN/dy(M^2)$, weighted by $1/M^2$, with the experimental gaussian resolution folded in. This has a negligible effect in the other mass intervals where one can take a fixed value of $M$ equal to the $\langle M \rangle$ of the considered interval.

The full line in fig. 4 shows the integrated multiplicity computed up to $M = 600 \text{ GeV}$ using the 2–chain graph of fig. 5 only. As in non-diffractive inelastic collisions, high order contributions containing more chains will become increasingly important with increasing $M$. The weights of these higher order contributions are not well known. For instance, the ratio of the 4-chain to the 2-chain components is proportional to the ratio of the 4-Pomeron over the 3-Pomeron coupling. If the Pomeron-Pomeron total cross-section is as small as recently estimated$^{23}$, we expect that the weights of the multi-chain components will be smaller than for inelastic collisions; multiplicity in diffractive events at $M = 546 \text{ GeV}$ will then be lower than for non-diffractive events at $\sqrt{s} = M$ (fig. 4).

3. CONCLUSIONS

Data on large-$t$ elastic scattering measured by the UA4 Collaboration at the CERN Collider were reported. The $t$-ranges covered were $0.45 \leq -t \leq 1.55$ at $\sqrt{s} = 546 \text{ GeV}$ and $0.7 \leq -t \leq 2.2$ at $\sqrt{s} = 630 \text{ GeV}$. Main feature of these data is the presence of a break in the
\( \frac{d\sigma}{dt} \) at \( -t \sim 0.9 \text{GeV}^2 \) followed by a shoulder with cross section as large as more than one order of magnitude than at the ISR. A comparison of the experimental data with theoretical expectation was done. High energy (ISR and Collider) large-\( t \) data were found compatible either with models predicting a non-vanishing odd contribution to the elastic scattering amplitude or with models based on a geometrical picture of the hadron collisions. The geometrical models, however, cannot account for the observed difference between pp and \( \bar{p}p \) cross-sections in the dip region as observed at \( \sqrt{s} = 53 \text{ GeV} \). A conclusive statement on the presence of an asymptotically non-vanishing odd term in the hadron elastic scattering amplitude cannot be made before further comparisons between high energy pp and \( \bar{p}p \) cross-sections.

A study of diffractive dissociation at the Collider was also presented. Data on pseudorapidity distribution of charged particles produced in the fragmentation of the diffractive states definitively disfavour an isotropic decay (fire-ball like) of these states. It was shown how the Dual Parton Model quantitatively describes multihadron production also in high-mass diffraction dissociation. A prediction was made on the possibility that at higher masses the shown similarity between diffraction dissociation and \( p\bar{p} \) inelastic interaction at \( \sqrt{s} = M \) can be broken due to the smaller contribution of higher order multi-chain graphs in diffractive collisions than in non-diffractive inelastic ones.

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Abstract

Soft hadronization data for pp collisions, obtained with EHS–RCBC exposed to 360 GeV/c protons, are analyzed with a quark–diquark fragmentation mechanism. Satisfactory agreements between data and model calculations support for an existence of rather a tight spectator diquark in the proton.

Using the particle ratios in the central region, the LUND model parameters are stringently tested. Results for local compensation of transverse momentum are compared with LUND model. The valon radius, in the case of three independent valons, is estimated.
Introduction

Topics covered in this talk are more or less related to models for soft hadronization. We assume a reaction mechanism of soft hadronic processes to be two step processes; collisions described in quark-parton level and subsequent hadronization mechanisms. A quark or a diquark could be regarded in soft collisions as extended objects, so-called dressed valence quarks (valons) in which whole effects of gluons and sea quarks are involved. Through 1.1 to 1.3, we stress the importance of the quark-diquark fragmentation mechanism in the soft proton-proton collisions. Here the diquark means an object which acts as a single cluster not splitting into two quarks in the hadronization process. In 1.1, the predictions of various models are compared with experimental Feynman $X_F$ distributions. For the quark structure of a proton, we adopt two viewpoints; a proton consists of a quark-diquark (in the SU(6) limit) and of three independent quarks. Collisions are described in a viewpoint of the DTU model limited in two sheets. For hadronization mechanisms of two colour neutral systems after exchange of "held back quarks", we adopt the independent cascade picture of the Field and Feynman or a string based (Lund) model. In 1.4 and 1.5, the LUND (single colour flux tube) model, full model describing soft hadron-hadron collisions, are examined. The model parameters are determined neatly and the string tear-up mechanism is tested in terms of the local transverse momentum compensation.

The data used in this analysis were obtained from the experiment at CERN with the European Hybrid Spectrometer (EHS) and the 80 cm Rapid Cycling Bubble Chamber (RCBC) exposed to 360 GeV/c protons. With its capabilities for 4W detection, particle identification with ISIS, SAD, $\gamma$ detection with IGD, FGD and precise momentum resolution in the whole region, the EHS is suitable place to study soft hadronic collisions. A detailed description of the set-up and data reduction of $V^0$ and $\pi^0$ were reported in refs. 1-4. The primary aim of this experiment was to investigate strange particle production, so that we took data in association with $V^0$ in a bubble chamber. It should be noted that the positive ($h^+$) and negative ($h^-$) particle distributions are $V^0$ triggered one.

1.1 Quark-diquark Fragmentation Mechanism for non-DD

We compared the data with model calculation with or without a diquark [5]. Fig. 1 shows the single particle $X_F$ distribution for $\Lambda^0$. We also show the predictions in cases of three independent quarks (3q) and quark-diquark (q-dq) fragmentation with the F.F model. The full LUND model prediction is also shown. The $X_F$ distribution with independent quark model (3q F.F.) is striking disagreement for the $\Lambda^0$. Making a comparison between the full LUND and the q-dq
F.F. models, we found both describe the data fairly well as shown in Figure 1.

In our model, $\Lambda^0$ is mainly created from the spin 0 valence (ud) diquark combined with a s-quark in the sea. This result is in accordance with an assumption of a semi-classical model, which explains the $\Lambda^0$ polarization observed in the quark fragmentation picture [4,6].

1.2 Primordial Motion of Quarks in a Proton

Transverse momenta, $p_T$, of hadrons produced in pp collisions can arise not only from the hadronization of quarks but also from the primordial motion of quarks in a proton. We thus introduce the primordial transverse momentum $k_T$ of quarks according to a function $\exp(-k_T^2/2\langle k_T^2 \rangle)$, where $\langle k_T \rangle$ is the average transverse momentum [7].

Figures 2 shows the $p_T^2$ distributions of $h^+$, where we can see at least two components of slopes if the spectra are decomposed into the exponential form. We examine model predictions derived from the F.F. model with the quark-diquark configuration for the proton.

One clearly sees in Fig. 2 that the model with $\langle k_T \rangle = 0.4$ GeV/c gives the most reasonable agreement with data especially for $p_T^2$ higher than 0.4 (GeV/c)$^2$.

For the $X_F\cdot p_T$ correlations (seagull-effect), we show the predictions from 3q-F.F., q-dq-F.F. and q-dq-Lund models. We find a significant role of the diquark especially for $h^+$ and $\Lambda^0$ as shown in fig.3(a)-(b).

1.3 Application to High Mass Diffraction Dissociation

Using another sub-sample, we have 2862 four prong events without strange particle decays. Among these we have selected 195 uniquely 4C fitted beam diffraction dissociation (DD) events.

There has been a naive idea that the quark-diquark jets could be observed in the proton DD. In the following analysis, we take a Gottfried-Jackson frame, the rest system of invariant mass $M_X$. The direction of the incident momentum of beam proton is taken to be $\cos \phi = 1$. Fig.4 shows the average of transverse and longitudinal momenta as functions of the invariant mass of the excited system, $M_X$. Obviously, we find the "$p_T$-suppression" above $M_X = 3$ GeV (high mass DD).

We show in fig.5 the Gottfried-Jackson angular distributions for high mass DD. There is remarkable forward and/or backward peaking. Especially protons succeed large fraction of incident momenta. We assumed for excited mass distribution in a conventional form of $1/M_X^2$ for high mass DD. Then one string quark-diquark fragmentation is calculated using the Lund model in the rest system of exited proton, where the diquark moves in the direction of incident
proton. The smooth curves in the figs.5 are the results of simulation. Striking agreements are observed [8].

1.4 Determination of Model Parameters

In the above sections, we used the model parameters which we fixed using various distributions in e^+e^- annihilations based on the jet universality. In order to determine model parameters neatly, the particle ratios are measured carefully in the central region, where produced particles are free from initial quark contents; i.e. in the very small interval of |X_F| < 0.04 for pions, for example, the particle yields for positive and negative pions are almost the same, which ensures the independence from initial quark contents [9]. Resulted particle ratios are shown in Fig.6. Particle ratios for K^-/π^- and p/π^- are experimentally 0.05 ± 0.02 and 0.010 ± 0.002 and with full LUND model 0.050 and 0.021, respectively.

Initial model parameters used are 0.3 and 0.09 for the sea-strange-quark and the sea-diquark picking-up rates, respectively. Our first-approximation result shows that the sea-diquark picking-up rate could be lowered to 0.05 instead of 0.09. If this trend of raising this parameter in e^+e^- annihilations on one hand and the lower value in low p_t hadronic collisions on the other hand comes out stable, this difference may signal something special beyond the "jet universality".

1.5 Local Transverse Momentum Compensation

Through the string tear-up mechanism, the short range correlation for the strangeness compensation is expected for particle pairs which carries (anti-)strangeness. For those particle pairs, the local compensation of transverse momentum is also expected, even if this effect is smeared through hadronization process. The azimuthal angle between transverse momenta of particle pair is measured. Using the numbers of pairs NO and NS, where the transverse momenta vectors are in the opposite-side and in the same-side, respectively, the asymmetry parameter B is defined as the ratio ; (NO - NS)/ (NO + NS). Particle pairs are classified as follows:

1) The rapidity gap is smaller than 2, where the local compensation is likely to be expected;
   1.a) includes s+s- pair, therefore the prominent effect is expected,
   1.b) which do not contain a common s+s- pair, then the effect is not likely expected.
2) The rapidity gap is greater than 2, therefore the effect is almost
unexpected, even for the particle pairs which carry (anti-)strangeness.

Calculated results, using 1337 $\Lambda K$ and KK pairs and 2894 $\Lambda h^+$ pairs, are shown in Fig. 7. In the case of independent emission of secondary particles, expected value of the asymmetry parameter $B$ is 0.1.

As shown in Fig. 7, the LUND model reproduces the observed tendency as expected, however it predicts slightly stronger correlations [10].

1.6 Inelastic Overlap Function and the Valon Radius

Using the differential elastic cross section, we obtained the inelastic overlap function $G_{\text{in}}(b)$ as a function of impact parameter $b$ as shown in Fig. 8.

Using this inelastic overlap function, we estimate the radius of the valon with a simple geometrical model. According to the idea that the transparency of hadrons are reflection of matter distribution inside the hadrons; the quark-gluon matter inside the hadron is distributed not uniformly but concentrates within discrete regions, we assume the proton consists of three independent valons. The condition of inelastic interaction is $d_{ij} < 2r_v$ for at least one pair of valons $i$ and $j$ from various protons, where $d$ is the impact parameter of pair of valons. Thus the model is controlled by the two parameters only - effective radii of hadron and valon $r_h$ and $r_v$, respectively. Estimated $r_h$ and $r_v$ are $0.98 \pm 0.03$ fm and $0.203 \pm 0.003$ fm, respectively [11]. The fraction of events for which only a single hit of valon pair takes a place is large (68%) as expected.

The analysis with the quark-diquark picture for the proton is now under way.

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FIG. 1

\[ \frac{d \sigma}{d Q} (\text{GeV/c}^2) \]

FIG. 2

\[ \frac{d \sigma}{d p_T^2} (\text{GeV/c}^2) \]

FIG. 3 (a)

\[ \langle p_T \rangle (\text{GeV/c}) \]

FIG. 3 (b)

\[ \langle p_T \rangle (\text{GeV/c}) \]

FIG. 4

\[ \langle p_T, M_X \rangle, \langle M_X \rangle \]
$a) \pi^+$

$bl \pi^-$

$cl \pi^-$

$FIG. 5$

$FIG. 6$

$FIG. 7$

$FIG. 8$
The origin of transverse momenta in minimum bias events is investigated within the framework of the Dual-Parton model. It is concluded that the scattering process at SPS-COLLIDER energies has to contain a substantial "semihard" component, which does not drastically affect the general event structure.
Yesterday you heard a talk about the dual parton model. The model explains essentially all features of highly inelastic processes, at least in an approximate way. To be specific, the model understands the limited charge transfer, the rise of the rapidity spectrum, various forms of long range correlations, the KNO scaling violation and the basic structure of diffractive events. In my talk I will address one particular aspect.

As you properly remember, the dual parton model assumed a factorization in the scattering between an initial process which is responsible for the production of strings and a final process which is responsible for their decay. The decay is parametrized as in \( e^+e^- \) annihilation or lepto-production in a scaling region; the initial string production is described by one or more Pomeron-exchanges.

The initial exchange of one or more Pomerons involves only small transferred momenta and the dual parton model is therefore essentially a longitudinal model. It is well known that hard or semihard processes will play with increasing energy a more and more important role and it is interesting to investigate the actual situation with minimum bias events at collider energies in this regard.

How can one observe a small transverse momentum of the initial partons? Their effect on most quantities turns out to be minute, as long as one assumes that the transition happens continuously, i.e. if a soft Pomeron-exchange slowly obtains a semi-hard gluon-exchange like component without otherwise changing the topological structure (i.e. the number of strings). Even the expected rise in the global average transverse momentum is quite diluted as the momentum of the string end is usually shared among several final particles.
A quite sensitive observable for the transverse structure is the correlation between the average transverse momentum and the associated multiplicity. If one selects events with an enhanced multiplicity in a certain rapidity interval, one typically chooses multi-Pomeron events i.e. events which contain more of the shorter additional strings. The partons entering the string on each side carry primordial $p_t$. This initial $p_t$ will be less diluted in these shorter strings than in the longer initial strings. An enhancement of the multi-string contribution therefore means an increase in the average $p_t$.

To be quantitative about this effect we first need to estimate the actual value of the transverse momentum obtained in the soft multi-Pomeron process. The eikonal width of our Reggeon calculus parametrisation of the total cross sections translates approximately in a mean square value of a Gaussian distribution of 0.7 GeV. Can this value explain the data? For collider energies the situation is definite, the eikonal transverse momenta are not enough to produce a sufficiently strong increase of the semi-inclusive average transverse momenta. To get a satisfactory representation of the data one needs an extra contribution to the primordial transverse momentum of partons.

To estimate the size of such an extra semi-hard contribution we consider the following parametrization of the hard process:

$$\text{Prob}(p_t^2) = \text{const.} \times (p_t^2 + (3.\text{GeV})^2)^{-2}$$

and adjust the probability to have such a hard scattering to occur. Satisfactory fits are obtained, if 40% hard scattering was used for 64 GeV and 55% for 540 GeV. Especially for in multi-string events the actually obtainable transverse momentum is severely restricted by kinematic constraints and the final transverse momentum on the partons averages only at 0.66 and 1.22 GeV. These values should not be taken too quantitatively as they depend on the details of the implementation. However, we tried many different implementations, no drastic effects could be obtained and the basic conclusion of a significant non soft component at SPS-COLLIDER energies is unavoidable.
To see if the semihard transverse momenta are reasonable in the framework of non-soft perturbative QCD, we borrowed the transverse momentum of semi-hard gluons calculated by GRIBOV, LEVIN and RYSKIN and obtain the fit given in the figure. Actually the calculation was before the considerably changed 84 data were available and the curve is therefore to a certain degree a prediction. It is obviously somewhat daring to combine both models as the relation between sea quarks and gluons is not clarified, but the point that the required $p_t$ actually follows quite reasonable expectations, is clearly demonstrated.

To put things together:

a) A careful analysis shows that at SPS energies the minimum bias events are no longer soft by a significant amount. However we found no drastic change in general structure of the events and except for very sensitive quantities like the considered correlation the required modification stays almost invisible.

b) Somehow the Pomerons start, with increasing energies, to look like semihard gluons and it would be nice to obtain a better theoretical understanding about this transition.

c) Obviously string ends will look like small jets and the model is therefore consistent with the observation of mini-jets. However mini-jet events do not play any special role and no clear cut separation between hard and soft events is possible. If one would have a unbiased jet-trigger, triggering jetty events one would obtain multi-string events which contain typically large multiplicities. One could observe the inverse of the above correlation. As the UA1 mini-jet trigger actually also favors high multiplicities, it is not easy to isolate this dynamical correlation.
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GLUON-GLUON INTERACTION IN HADRON INDUCED
MULTIPARTICLE PRODUCTION PROCESSES

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ABSTRACT

The observed dependence on energy and that on (pseudo)rapidity of multiplicity
distributions in high-energy nondiffractive pp collisions can be understood
in terms of a statistical model proposed some time ago by the Berliner group.
The result is interpreted in the quark-gluon picture. The possible role of
 gluon-gluon interactions in such processes is discussed.

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Max Planck Gesellschaft.
One of the recent experimental results from the CERN $\bar{p}p$-collider that has attracted much attention is the observed dependence on (the total c.m.s.) energy $\sqrt{s}$ and (pseudo)rapidity of multiplicity distributions in non single-diffractive collisions. I shall, in the first part of my talk, address myself to this and other related phenomena and show that these experimental facts can be understood in terms of a statistical model proposed by the Berliner group. In the second part, I shall interpret these results in terms of the quark–gluon picture and discuss the possible role of gluon-gluon interactions in hadron induced low-$p_T$ multiparticle production processes.

We recall: The underlying physical picture of the proposed model can be summarized as follows: In a typical high-energy nondiffractive hadron-hadron collision event, the colliding objects $P$ (projectile) and $T$ (target) "go through" each other and appear as "leading particles". During this process, $P$ and $T$ lose a considerable part of their energies and momenta. A part of this energy materializes— in general into a number of clusters which subsequently decay into hadrons. The materialization energy is distributed in three distant systems $C^*$, $P^*$ and $T^*$ which are characterized by their locations in rapidity space. ($C^*$ stands for central rapidity region, $P^*$ and $T^*$ stand for projectile and target fragmentation regions respectively).

Under the assumption that each of these systems randomly obtain its materialization energy from two independent sources and that each system forgets its history after the formation (Note that both randomness and memory loss are characteristic properties of statistical processes.), we obtain

$$\langle E_i^* \rangle P(E_i^*) = \frac{4}{\langle E_i^* \rangle} \exp(-2 \frac{E_i^*}{\langle E_i^* \rangle})$$

(1)

Here, $\langle E_i^* \rangle$ is the materialization energy of (that is, the total internal energy of the created clusters in) the system $i$ ($i = C^*, P^*, T^*$), $\langle E_i^* \rangle$ is its average value. $P(E_i^*)$ is the probability of finding the system $i$ in the state characterized by $E_i^*$.

The total internal energy of the clusters is obviously directly proportional to the total transverse energy. Taken together with the empirical fact that the measured transverse energy per charged hadron (at given $\sqrt{s}$) is approximately the same, this leads us to the ansatz $E_i^*$ is directly proportional to the number of charged hadrons produced by the system $i$ ($i = C^*, P^*, T^*$). Hence, we have for the central rapidity region,
This result\(^3\) is in very good agreement with the data\(^6\).

Multiplicity distribution for the full rapidity range can be, and has also been calculated by taking the contributions of the P* and the T* into account\(^3\). The agreement between data\(^6\) and the calculated result has led us to the conclusion that such production processes are statistical processes.

In order to see whether the experimentally observed (pseudo) rapidity dependence of multiplicity distributions is also a statistical effect, let us first consider a limited rapidity window \(W\) in the central rapidity region. Since in this region the products of the central system C* dominates, the distribution \(P_W(n_w)\) for the multiplicity \(n_w\) of the charged hadrons inside the rapidity window \(W\) can be written as

\[
P_W(n_w) = \sum_{n_C=2,4,6,...} \frac{n_C}{n_W} \left( \frac{n_C}{2} \right)^{n_C/2} \frac{n_W}{(1-q_W)}^{n_W/2} (1-q_W)^{n_C/2-n_W/2}
\]

where \(q_W = \langle n_W \rangle / \langle n_C \rangle\); \(\langle n_W \rangle\) is the average multiplicity of charged hadrons inside the rapidity window; \(\langle n_C \rangle\) is the average multiplicity of charged hadrons produced by the system C*. Here, \(q_W\) can be calculated in this model\(^3\) where we have, for the sake of simplicity, assumed that every produced cluster consists of two charged hadrons; the rapidity distribution of the clusters is flat and every cluster decays isotropically in its rest frame. We note that, since in this approximation we only include the contribution of the central system C*, some difference between data and the calculated result for large rapidity windows is expected. This difference should however vanish when the contribution of the P* and the T* systems are taken into account. More detailed study\(^3\) shows that this is indeed the case.

We note that Eq.(3) implies \(P_W(n_w) = P(n_C/2)\) for \(q_W = 1\). That is, if the rapidity window \(W\) in the central rapidity region is on the one hand large enough to include all the hadrons produced by the system C*, but on the other hand not too large so that the contributions from the P* and T* systems are negligible, the multiplicity distribution \(P_W(n_w)\) is expected to satisfy

\[
\langle n_w \rangle P_W(n_w) = 4 \frac{n_w}{\langle n_w \rangle} \exp(-2 \frac{n_w}{\langle n_w \rangle})
\]
where \( <n_w> \) is the average value of \( n_w \) in the rapidity window \( W \). It means in particular that for every given (total c.m.s.) energy \( \sqrt{s} \) (sufficiently high so that the overlapping regions of \( C^* \) with \( P^* \) and \( T^* \) systems in rapidity space are relatively small) there is a rapidity interval \( W \) in which the probability \( P_w(n_w) \) to find \( n_w \) charged hadrons in the interval \( W \) satisfies Eq. (4).

In terms of the quark-gluon picture in which hadrons are made out of valence quarks and gluons, and on the average about 35-50% of the momentum of a high-energy hadron is carried by its valence quarks and the rest by the gluons (and/or sea quark pairs), the hadronic matter created in system \( C^* \) is due to the interaction between the gluons (and/or sea-quark pairs) in the projectile and those in the target. That is, the two sources from which the \( C^* \) system obtains the materialization energies (that is the masses) and the momenta for its clusters are the gluon-part of the projectile and that of the target. Hence, the growth in width as well as in height of the rapidity distribution for increasing total c.m.s. energy \( \sqrt{s} \) means in this picture that the gluon-gluon interaction becomes more and more important for higher and higher bombarding energies.

However, as we have seen in the preceding discussions, dynamical details do not play an important role, as long as we are only interested in the gross features of such processes such as multiplicity distributions, rapidity distributions, long-range multiplicity correlations and even short-range rapidity correlations provided that we do not ask questions such as the following. "What are the hadronic clusters?" "Why do they exist in high-energy hadron-hadron collisions?" "Why are they so small?" "The properties of the hadronic clusters in and/or near the central rapidity region are found to be independent of the quantum numbers of the incident particles, and independent of the bombarding energy. Why?"

In a recent paper, Chao, Gao and myself made an attempt to link these questions with the double role of the gluons (They carry color-quanta and they are responsible for the confining color forces.). We speculate that the hadronic clusters which have been introduced to explain the observed short-range correlations in hadron-induced multiparticle production processes are nothing else but color singlet gluon-clusters. Such gluon-clusters are in general extremely short-lived. That is, they are metastable objects, the decay widths of which are in general much broader than those of ordinary resonances. (That is, the "glueball" candidates are special cases of gluon-clusters)
Because of the limited time (for the talk) and space (in the written version) I can not discuss the details about the gluon clusters. Let me just briefly mention two of the conclusions: First, a new lower limit for the meson-baryon ratio, \( M/B = 12.3 \), is obtained which is more than twice as large as the well-known value given by Anisovich and Collaborators\(^9\) and closer to the experimental data\(^10\). Second, there should be in particular a number of \( I=0, J^{PC}=0^{++} \) glueball candidates with different masses. The recently found resonance by Tanimori et al.\(^11\), and those by Morgen et al.\(^12\), seem to belong to this group.

References and Footnotes

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INCLUSIVE CHARGED PARTICLE DISTRIBUTION IN NEARLY 3-FOLD SYMMETRIC 3-JET EVENTS AT $E_{cm} = 29$ GEV

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ABSTRACT

Results of inclusive charged particle distribution for gluon jets using nearly 3-fold symmetric 3-jet events taken at center of mass energies of 29 GeV in $e^+e^-$ annihilation are presented. The charged particle spectrum for these jets is observed to be softer than that of quark jets with the same jet energy.

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Jets initiated by a quark or antiquark have been measured in great detail by various experiments. However, experimental studies about jets which originate from a high energy gluon are just at the beginning.\textsuperscript{[1],[2]}

We studied inclusive charged particle production in 3-jet events which have nearly 3-fold symmetry under the assumption that these events originate from two quark jets and one gluon jet. This requirement has the following advantage: all three jets have nearly the same energy $\sim 1/3 \ E_{cm}$, the jets have the best possible separation, and the gluon jet has a relatively high energy. On the other hand the production of such events is very suppressed, and it has been shown that quark and gluon jets do not fragment totally independently in an event,\textsuperscript{[3]} but these results show that mainly the soft particles are affected, whereas the high momentum particles follow more the direction of the partons. We compared the particle distribution of these 3-jet events with distributions of events originating from initial quark jets with the same jet energies.\textsuperscript{[4]}

The data sample used in this measurement was collected by the Mark II detector at $E_{cm} = 29$ GeV. The total integrated luminosity of 215 pb\textsuperscript{-1} corresponds to 90000 hadronic events. A subset of these data was used for a measurement of the inclusive charged particle cross section.\textsuperscript{[5]}

The Mark II detector has been described in detail elsewhere.\textsuperscript{[6]} The track selection criteria are the following: a well reconstructed charged track has to pass within 40 mm in radius and 60 mm in $z$ from the event vertex and have at least 100 MeV/c of transverse momentum. Neutral particles assumed to be photons are detected by the central region lead-liquid argon calorimeter modules.\textsuperscript{[7]} A neutral cluster with $E > 150$ MeV and a distance (at the radius of the shower counter) of $d > 300$ mm from the closest charged track is defined as a photon.

Hadronic events are selected by requiring at least 5 well reconstructed charged tracks and a total charged and neutral energy $> 55\%$ of $E_{cm}$. A cluster algorithm\textsuperscript{[8]} which uses the vector momenta of charged and neutral particles partitions the data into n-jet events. Only 3-jet events are retained of which each jet has to have greater than 2 GeV of observed energy and to contain at least 3 (charged or neutral) particles. The jet axes are defined by the vector sum of the particle momenta within each jet, and the jet energies ($E_j$) are calculated from the angles between the jet axes assuming three massless partons. To require almost 3-fold symmetry for the 3-jet events, all three angles between the jet axes have to lie between $100^\circ$ and $140^\circ$. All of the above criteria are met by 560 events which corresponds to about 0.5\% of all hadronic events. Monte Carlo calculations using models having QCD plus fragmentation estimate a background of $(0.4 \pm 0.2)\%$. The model calculations show further
that the fraction of heavy quarks in this sample is the same fraction as for all hadronic events.

The inclusive charged particle distribution is analyzed in terms of the fractional momentum $x_i = p_i / E_j$, where $p_i$ is the momentum of particle $i$, and $E_j$ the energy of the jet to which it is assigned. The fact that all three jets have nearly the same energy implies that a wrong assignment of a particle to a jet is not a severe problem. The data are corrected by a bin by bin correction factor, calculated by different Monte Carlo generators\cite{9,10,11} to correct for the detector inefficiencies and initial state QED radiation. The correction factor increases slightly from 0.85 at low $x$ to 0.92 at high $x$ values.

Figure 1 shows the corrected $x$ distribution of the 3-fold 3-jet events (full symbols), where the errors include both statistics and the uncertainty in the efficiency. To compare this $x$ distribution originating from events containing two quark jets and one gluon jet, each with approximately 10 GeV energy, with $x$ distributions initiated by events containing two quark jets each with the same energy, we use the published data at $E_{cm} = 5.2, 6.5, 14, 22, 29$ and $34$ GeV,\cite{5,12,13,14} In Fig. 2 all these results are shown in terms of their c.m. energy dependence. $1/N_{jet} \, dN/dx$ is plotted for the twelve fixed $x$ intervals shown in Fig. 1 as a function of the beam energy. Within a fixed $x$ interval, all the data points with $N_{jet} = 2$ are fitted to the form suggested by QCD:\cite{15}

$$\frac{1}{\sigma_{tot}} \frac{d\sigma}{dx} = c_1(x)(1 + c_2(x) \ln(s))$$

where $c_1(x)$ and $c_2(x)$ are free parameters. The resulting fits are represented by the curves in Fig 2. For $x > 0.15$ all the different data points agree quite well with the fitted curves, whereas for low $x$ values some larger deviations are visible, probably due to higher background problems for low momentum tracks.

![Fig. 1. The detector corrected inclusive charged particle distribution for 3-jet events at $E_{cm} = 29$ GeV (full symbols) in comparison with the inclusive charged particle cross section of hadronic events at $E_{cm} = 19.3$ GeV, extrapolated from the fitted curves in Fig. 2 (dotted curve) and the inclusive charged particle distribution of a gluon jet of $E_j = 9$ GeV (open symbols).](image)

The $x$ distribution of the 3-fold 3-jet events at $E_{cm} = 29$ GeV is also shown in Fig. 2 for the twelve $x$ intervals. The points are drawn at $E_j = 9.66$ GeV. The deviation between the points and the fitted curves suggests a difference in the $x$ distribution.
The results of the fits within the twelve intervals from Fig. 1 are used to interpolate the cross section at $E_{cm} = 19.3$ GeV. This is shown in Fig. 1 as a dotted curve. The slope of the distribution containing the 1/3 admixture of gluon jets is observed to fall off faster than that of initial quark jets. To extract to first approximation an inclusive charged particle distribution for a gluon jet of $\approx 10$ GeV energy, we adopted the following ansatz:

$$\frac{1}{\sigma_{tot}} \frac{d\sigma}{dx} \text{(gluon jets)} = \frac{1}{\sigma_{tot}} \frac{d\sigma}{dx} \text{(3-jet events, } E_{cm} = 29 \text{ GeV)} - \frac{1}{\sigma_{tot}} \frac{d\sigma}{dx} \text{(all events, } E_{cm} = 19.3 \text{ GeV)}$$

(2)

where for the events at $E_{cm} = 19.3$ GeV the fit results are again used. This cross section is also shown in Fig. 1 (open symbols).

Another way of displaying the data is to calculate the ratio

$$r(x) = \frac{\frac{1}{\sigma_{tot}} \frac{d\sigma}{dx} \text{(3-jet events, } E_{cm} = 29 \text{ GeV)}}{\frac{1}{\sigma_{tot}} \frac{d\sigma}{dx} \text{(all events, } E_{cm} = 19.3 \text{ GeV)}}$$

(3)

as a function of $x$ which is shown in Fig. 3. For comparison, the calculation for the models are also indicated in Fig. 3.

Fig. 2. The inclusive charged particle distribution for a jet versus the jet energy measured by various experiments. The curves represent fits to the different data points for each of the twelve $x$ intervals at till. The inclusive charged particle distribution for the 3-jet events is also shown.

Fig. 3. The ratio of the inclusive charged particle distribution for 3-jet events at $E_{cm} = 29$ GeV to the inclusive charged particle cross section of hadronic events at $E_{cm} = 19.3$ GeV, together with several model predictions.
For the simulation of $q\bar{q}$ and $q\bar{q}g$ events the following steps have to be taken into account in the models, although there is no distinct separation between the different processes: Hard parton emission from quark jets, hard parton emission from gluon jets, fragmentation of light quark jets, fragmentation of heavy quark jets (It has been shown that $x$ distributions of particles from heavy quark jets are softer than that of light quark jets), and fragmentation of gluon jets. Assuming that parts of the hadronisation are similar for quark and gluon jets (e.g. probability of higher spin resonances, limited phase space at low energies, decays of heavy resonances), these things should cancel out in the ratio, and $r$ should be less sensitive to fine tuning of the different models.

The Lund model, in which not the partons themselves but a string, stretched from the quark via gluon(s) to the antiquark, fragments and which contains a $y_{\text{min}}$ cut to distinguish between the parton classes ($y_{\text{min}}$ is the smallest invariant mass of any two partons over the c.m. energy), has $r$ values close to unity over the whole $x$ region. Although the fragmentation of a gluon jet differs quite drastically from a quark jet in this model, it does not show up in the ratio $r$. One reason is that the fixed $y_{\text{min}}$ cut, implement more gluon emission at $E_{cm} = 19$ GeV. A further point in the model is the insufficient simulation of further parton radiation in the 3-jet events according to next higher orders, which are not accounted for, using only calculations up to second order in $\alpha_s$.

On the other hand, the Webber model, which uses leading log evolution for the parton radiation, accounting for multiple parton emission in a better way, exhibits good agreement with the data. In the Webber model a $q\bar{q}$ event at $E_{cm} = 19$ GeV has on the average 4.1 partons, whereas a 3-fold $q\bar{q}g$ event at $E_{cm} = 29$ GeV produces 6.7 partons, resulting in a ratio of parton multiplicity of 1.3. This is still far away from 9/4 which is expected at infinite energies, indicating that the limited phase space plays still an important role at these energies. The limited phase space effect and the higher multiplicity of heavy quark jets cause in addition that the ratio in the charged multiplicity only amounts to 1.15.

A delicate issue in the Webber model is the decay mechanism of heavy clusters, which is normally used for cluster masses $> 3.5$ GeV. These decay according to a string like picture into two lighter clusters, which may pass the same procedure again, if their clusters are still to heavy. But using only this mechanism for the whole quark cascade leads to a particle $x$ spectrum which is even softer than that for a gluon jet. So the following condition arises: If in the model the gluon branching $g \to gg$ is arbitrarily reduced by a factor $4/9$, the parton multiplicity of the quark and gluon jets are nearly equal, but the gluon jet has more heavy clusters such that the final charged multiplicity of the gluon jet and the $r$ distribution are the same as without the factor $4/9$. This means that in the Webber model the changes on the parton level do not result in changes on the hadron level.
Another possibility is to use the parton shower of the Webber model and to stretch at the end of the parton evolution a string from the quark via the gluons to the antiquark and let this string fragment according to the Lund model. The $r$ distribution observed is similar to that of the original Webber model (dotted line in Fig. 3), and now this distribution is sensitive to changes in the gluon branching as one should expect.

In conclusion, we have used nearly 3-fold symmetric 3-jet events produced in $e^+e^-$ annihilation at 29 GeV to investigate the inclusive charged particle distribution of gluon jets. Although the color configuration in these events differs from that in $q\bar{q}$ events, we have as a first approximation compared these distributions with charged particle distribution of all hadronic events at $E_{cm} = 19$ GeV. The comparison suggests a steeper slope for $x > 0.4$ for gluon jets.

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ABSTRACT

We have determined the parameters of $\rho'(1250)$ and $\rho''(1600)$ vector mesons from data on $e^+e^-\rightarrow \pi^+\pi^-$. The analysis was carried out by means of a realistic three-resonances pion form factor model. The latter is based on minimally simplified analytic structure of the pion form factor, reflects its other well-established properties and depends just on physical parameters. The validity of the model is further confirmed in prediction of the behaviour of the P-wave isovector $\pi\pi$ scattering amplitude at the experimentally measurable region.
Due to the electromagnetic nature of the process \( e^+e^- \rightarrow \pi^+\pi^- \) one is allowed to treat it in one-photon exchange approximation. The process provides therefore a model independent experimental information on the pion form factor (ff) in the time-like region. In other words, in order to explain the \( e^+e^- \rightarrow \pi^+\pi^- \) cross section theoretically, it is sufficient to understand the pion ff behaviour, but extended also to the space-like region.

Almost 15 years ago Renard\(^1\) drew an attention to the structures in \( e^+e^- \) annihilation channels. Though he argued that one of the possible explanations is the presence of higher vector resonances, the poor data at that time did not allow to determine their parameters.

In this contribution we show that the new more precise and dense experimental data on \( e^+e^- \rightarrow \pi^+\pi^- \) from Novosibirsk\(^2\) and ORSAY\(^3\) have changed the situation; one is able to determine the parameters of \( \rho(1250) \) and \( \rho^*(1600) \).

To achieve this goal we have constructed\(^4\) a realistic pion ff model, which respects the analyticity and other well-established pion ff properties as precisely as possible and depends just on parameters with clear physical meaning.

The analytic properties of the pion ff \( F_\pi(t) \) used in our model consist of a square-root branch points at \( t = 4m_\pi^2 \) and at \( t_{inel} > 4m_\pi^2 \), of the pole at \( t_\rho = -15.36m_\pi^2 \) and the zero at \( t_z = -8.96m_\pi^2 \) both on the second Riemann sheet, and of the poles \( t_\rho' \), \( t_\rho'' \), \( t_\rho^* \) on unphysical sheets, corresponding to \( \rho(770) \), \( \rho(1250) \) and \( \rho^*(1600) \) resonances. The first branch point at \( t = 4m_\pi^2 \) corresponds to the elastic threshold, the second one at \( t = t_{inel} \) to the beginning of an effective inelastic cut, to be left as a free parameter. The pole \( t_\rho \) and the zero \( t_z \) approximate the so called pion ff left-hand cut.

The cut structure used generates the four-sheeted Riemann surface, on which our model of \( F_\pi(t) \) is defined. The \( \rho(770) \) is placed on the second sheet at \( t_\rho = (m_\rho + i\Gamma_\rho/2)^2 \). Since we expect\(^4\)

\[ t_{inel} < m_\rho^2, m_{\rho'}^2, \rho(1250) \) and \( \rho^*(1600) \) are situated on the third sheet at \( t_\rho' = (m_\rho + i\Gamma_\rho'/2)^2 \) and \( t_\rho^* = (m_\rho + i\Gamma_\rho^*/2)^2 \).

By using the inverse Zhukovsky transformation

\[ W(t) = i [(q_1^+ q)^{1/2} - (q_1^- q)^{1/2}] / [(q_1^+ q)^{1/2} + (q_1^- q)^{1/2}] \quad (1) \]

where

\[ i = (-1)^{1/2}, q_1 = 1/2(t_{inel} - 4)^{1/2}, q = 1/2(t - 4)^{1/2}, m_\pi = 1 \]
we map all four sheets of Riemann surface in t-variable onto the W-plane and all cuts disappear. The variable $W(t)$ possesses the assumed pion $f\bar{f}$ dominant cut structure in t-variable and the same holds for the expansion

$$F_{\pi}[W(t)] = \frac{(W^2 - 1)^m (W - W_0)^{\sum_{n=0}^{L} A_n W^n}}{(W - W_p)^{\nu} (W - W_v) (W - W_v^*)}$$

with arbitrary $L$ and real coefficients $A_n$ (the latter is the consequence of the reality condition $F_{\pi}(t) = F_{\pi}(t^*)$). The factor $(W^2 - 1)^m$ is present to ensure the asymptotic behaviour of the form $t^{-m/2}$ with $m$ being a free positive integer.

It is well known that the resonances are besides $m_V$, $\Gamma_V$ characterized also by the residue at the corresponding pole. The residue contains the ratio of coupling constants $f_{\pi V \pi V} / f_V$, where the universal vector-meson coupling $f_V$ determines the lepton width $\Gamma(V \rightarrow e^+e^-)$ and $f_{\pi V \pi V}$ the partial width $\Gamma(V \rightarrow \pi^+\pi^-)$. In order to find explicit dependence of our model on the ratio $f_{\pi V \pi V} / f_V$, we employ the vector meson dominance (VMD) pole-diagram contributions

$$F_{\pi}(\text{VMD})(t) = \sum_{\nu} \frac{m_V^2 (f_{\pi V \pi V} / f_V)}{m_V^2 - t} = \sum_{\nu} F_{\pi}(\nu)(t)$$

transformed into the $W$-variable:

$$F_{\pi}(\nu)(t) = \frac{m_V^2 (f_{\pi V \pi V} / f_V) [(1 - W^2)(1 - W_{\nu V}^2)]^2}{16 q_1^2 (1 - W_{\nu V}^2 W^2)(W^2 - W_{\nu V}^2)}$$

$$\nu = \rho, \omega, \rho^*$$

Here

$$W_{\rho O} = -W_{\rho O}^* = \lim_{\Gamma_{\rho O} \rightarrow 0} W_{\rho O}, \quad W_{\rho O} = 1/W_{\rho O}^* = \lim_{\Gamma_{\rho O} \rightarrow 0} W_{\rho O}, \quad W_{\rho O} = 1/W_{\rho O}^* = \lim_{\Gamma_{\rho O} \rightarrow 0} W_{\rho O}^*$$

In expression (2), for every resonance only the nearest to the physical region complex conjugate pair of poles is considered. Therefore, we extract only the term with denominator of the form $(W - W_{\nu V})(W - W_{\nu V}^*)$ from (4) and demand its equality to the equivalent term from (2) in the limit $\Gamma_{\nu} \rightarrow 0$. 
This results in the following equations:

\[
\sum_{n=0}^{L} \left(C_{\rho\rho} W_n^n - C_{\rho\rho} W_n^s^n\right) A_n = 2r_{\rho\rho} \quad (5)
\]

\[
\sum_{n=0}^{L} \left(C_{\rho\rho} W_n^n + C_{\rho\rho} W_n^s^n\right) A_n = r_{\rho\rho} + r_{\rho\rho}^s \quad (6)
\]

\[
\sum_{n=0}^{L} \left(C_{\rho\rho} W_n^n + C_{\rho\rho} W_n^s^n\right) A_n = r_{\rho\rho} + r_{\rho\rho}^s \quad (7)
\]

where

\[
C_{\rho\rho} = \frac{(W_{\rho\rho}^2 - 1)^m(W_{\rho\rho}^p - W_{\rho\rho}^z)}{(W_{\rho\rho}^p - W_{\rho\rho}^z)(W_{\rho\rho}^p - W_{\rho\rho}^s)(W_{\rho\rho}^p - W_{\rho\rho}^s)(W_{\rho\rho}^s - W_{\rho\rho}^s)(W_{\rho\rho}^s - W_{\rho\rho}^s)}
\]

\[
C_{\rho\rho}^s = \frac{(W_{\rho\rho}^s - 1)^m(W_{\rho\rho}^s - W_{\rho\rho}^z)}{(W_{\rho\rho}^s - W_{\rho\rho}^z)(W_{\rho\rho}^s - W_{\rho\rho}^s)(W_{\rho\rho}^s - W_{\rho\rho}^s)(W_{\rho\rho}^s - W_{\rho\rho}^s)(W_{\rho\rho}^s - W_{\rho\rho}^s)}
\]

and \( r_{\rho\rho}, r_{\rho\rho}^s, r_{\rho\rho}^s \) are the residua of \( F^{(v)}_n(t) \) at the poles \( \rho_{\rho\rho}^o \), \( W_{\rho\rho}^o \) and \( W_{\rho\rho}^o^s \), respectively.

The threshold behaviour conditions

\[
\text{Im } F_n(t) \big|_{q=0} = \frac{d}{dq} \text{Im } F_n(t) \big|_{q=0} = \frac{d^2}{dq^2} \text{Im } F_n(t) \big|_{q=0} = 0 \quad (8)
\]

lead to the relation

\[
A_0 + R(\rho, \rho, \rho') A_1 = 0 \quad (9)
\]

with

\[
R(\rho, \rho', \rho'') = \left[\frac{(W_{\rho\rho}^p - W_{\rho\rho}^z) - 1}{|W_{\rho\rho}^p|^2} + \frac{2\text{Re } W_{\rho\rho}^p}{|W_{\rho\rho}^p|^2} + \frac{2\text{Re } W_{\rho\rho}^s}{|W_{\rho\rho}^s|^2} + \frac{2\text{Re } W_{\rho\rho}^s - 1}{|W_{\rho\rho}^s|^2}\right]
\]

The normalization condition gives

\[
\sum_{n=0}^{L} A_n W_n^n = C_N(\rho, \rho', \rho) \quad (10)
\]

where

\[
C_N(\rho, \rho', \rho') = \frac{(W_N - W_{\rho\rho}^p)(W_N - W_{\rho\rho}^s)(W_N - W_{\rho\rho}^s)}{(W_N^p)^2 - 1)^m(W_N - W_{\rho\rho}^z)}
\]

and \( W_N = W(t) \big|_{t=0} \).
The present experimental situation does not require more than 5 terms in the expansion of (2). Consequently we restrict ourselves to $L=4$. Solving the closed system of the algebraic equations (5), (6), (7), (9) and (10) for $A_n$, one has the analytic pion ff model with all fundamental properties, which depends only on the following physical parameters: $m$, $t_{\text{inel}}$ and $m_\nu$, $\Gamma_\nu$, $f_{\nu\pi\pi}/f_\nu$ for $\nu = \rho$, $\rho'$, $\rho''$.

The analysis of experimental data compiled in ref. 7) joined with new data from Novosibirsk2) and ORSAY3) by means of (2) gives the results as follows (for data and our fit see fig. 1):

$$\chi^2/\text{ndf} = 1.85 \quad m = 3, \text{i.e. } F_\pi(W(t)) \sim t^{-3/2}$$

$$t_{\text{inel}} = 1.46 \pm 0.01 \text{ GeV}$$

$m_\rho = 760.1 \pm 0.5 \text{ MeV}, \Gamma_\rho = 152.4 \pm 0.6 \text{ MeV}, f_{\rho\pi\pi}/f_\rho = 1.24 \pm 0.01$

$m_\rho' = 1217.4 \pm 0.3 \text{ MeV}, \Gamma_\rho' = 323.8 \pm 4.3 \text{ MeV}, f_{\rho'\pi\pi}/f_\rho' = -1.10 \pm 0.01$

$m_\rho'' = 1531.6 \pm 15.5 \text{ MeV}, \Gamma_\rho'' = 454.0 \pm 18.3 \text{ MeV}, f_{\rho''\pi\pi}/f_\rho'' = -0.07 \pm 0.01$

The parameters of $\rho(770)$ and $\rho''(1600)$ are in satisfactory agreement with their known values8). In addition, a clear evidence for $\rho'(1250)$ in $e^+e^-\rightarrow\pi^+\pi^-$ is obtained, with the mass and the width in accordance with earlier, nowadays essentially abandoned hints.

The last remark concerns the further test of validity of our pion ff model; it predicts9) the behaviour of the P-wave isovector $\pi\pi$ phase shift and inelasticity, which reproduce the experimental data quite well (see fig. 2).

REFERENCES

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Fig. 2

\[ \delta_1' \]

\[ \eta_1' \]
The electromagnetic form factors of $\pi$ and $K$ mesons have been measured by scattering these particles from the electrons in a liquid hydrogen target. The variation of the form factors with four momentum transfer gives the root mean square of the pion charge radius, $<r^2>_{\pi} = 0.663 \pm 0.006$ fm and that of the $K$ meson $<r^2>_K = 0.63 \pm 0.05$. The $K$ meson result is provisional.

A measurement of the reaction $\pi^- + e^- + \pi^- + \pi^0 + e^-$ gives a result in good agreement with PCAC with 3 colours.
The elastic scattering of pions on electrons is described by the following diagram:

![Diagram](image)

where $q$ is the four momentum transfer in the collision due to the exchange of a virtual photon. The pion photon vertex involves a form factor $F$ which is a function of $q^2$ only. The differential scattering cross-section is given by the expression:

$$
\frac{d\sigma}{dq^2} = \frac{4\pi\alpha^2}{q^4} \left( 1 - \frac{q^2}{q_{\text{max}}^2} \right) F^2(q^2)
$$

where $F$ would be 1 for a point charge. The mean square radius of the pion charge distribution is given by: $\langle r^2 \rangle = \lim_{q^2 \to 0} -6 \frac{d^2F}{dq^2}$.

The pion form factor was measured using an incident energy of 300 GeV. With this incident energy the maximum value of $q^2$ is $0.288 \text{ (Gev)}^2$. The portion of this experiment covering the range $0.014 < q^2 < 0.122 \text{ (Gev)}^2$ has already been published. This gave the root mean square charge radius $\langle r^2 \rangle^{\frac{1}{2}} = 0.657 \pm 0.012 \text{fm}$.

We have now completed the analysis up to $q^2 = 0.26 \text{ (Gev)}^2$. Beyond this the number of events is negligible. The analysis for $q^2 > 0.12 \text{ (Gev)}^2$ presents special problems. At lower $q^2$ an accurate knowledge of the scattering angles of the two particles is sufficient to identify which trajectory corresponds to the electron and which to the pion because the electron angle is larger than the largest scattering angle experienced by the pion. In the neighbourhood of $q^2 = 0.15 \text{ (Gev)}^2$ the pion and the electron have the same momentum and angle. In the region of $0.12 < q^2 < 0.18$ approximately, a knowledge of the angles alone allows two possibilities for the value of $q^2$ depending on which trajectory is assumed.
to be which particle. This is called the ambiguity region. This can only be resolved by using additional information to identify the electron and the pion. This additional information is provided by electro-magnetic shower detectors which form part of the forward spectrometer. Each of these shower detectors is divided into a front portion of a thickness of 4 radiation lengths which is followed by another 22 radiation lengths. The identification of electrons and pions is improved considerably by the information of the longitudinal shower development provided by this subdivision.

Fig. 1 shows the overall layout of the experiment and Fig. 2 shows the arrangement of the shower detectors. Fig. 3 illustrates the way in which the information from the shower detectors is used. The vertical axis gives the ratio of the shower energy absorbed by the front portion to the total energy absorbed by the whole shower detector. The horizontal axis gives the ratio of the total shower energy to the energy of the incident charged particle as measured by the magnets of the forward spectrometer. Fig. 3a contains only particles which have already been unambiguously identified as electrons by the scattering kinematics and Fig. 3b contains only those particles which have similarly been identified as pions. The dotted lines define very safe cuts for particle identification. Events between these cuts are identified by a detailed examination of the particle trajectories. They correspond mainly to electrons which strike the sides of the central holes of the shower detectors. It should be borne in mind that only those events are wrongly attributed in which both particles have been misidentified.

Fig. 4 shows the complete result of this experiment and an inset giving the results of the best previous measurement by an American Russian collaboration\(^2\) who were the pioneers in this field.

The result fits the parametrisation proposed by Dubnicka and Martinovic\(^3\) which is linked to the known \(\pi \pi\) phase shift in the region of the \(\rho\) meson. This fit has a \(\chi^2\) probability of 55%. It gives a pion charge radius of \(\langle r^2_{\pi}\rangle^{\frac{1}{2}} = 0.663 \pm 0.006\text{fm}\).
We also measured the kaon charge radius using 250 Gev kaons identified by a differential gas Cerenkov counter in the incident beam line. In every other respect this experiment was very similar to the 300 Gev pion experiment described above. The result, which is provisional, is shown in Fig. 5 giving a charge radius of $\langle r_K^2 \rangle^{1/2} = 0.63 \pm 0.05$ fm.

At the same time as data were taken for $\pi^- - e$ elastic scattering at 300 Gev, we also looked for events of $\pi^- + e^- \rightarrow \pi^- + \pi^0 + e^-$. This work has been published\(^4\). Such events were identified by looking at the missing mass from charged particle kinematics when the shower detectors detected $\gamma$ rays consistent with the decay of a $\pi^0$, Fig. 6. The result is in good agreement with the predictions of PCAC with three colours (Fig. 7).

References

MEASUREMENT OF SOFT PHOTONS AT 63 GeV/c

The Axial Field Spectrometer Collaboration

*Presented by J.A. Thompson*

ABSTRACT

The production of soft ($p_\perp < 1$ GeV) photons has been studied in the Axial Field Spectrometer apparatus at the CERN Intersecting Storage Rings at $\sqrt{s} = 63$ GeV. Limits on a direct component are presented and compared with results at $\sqrt{s} = 12$ GeV.
We report on a measurement of low transverse momentum γ's at a centre-of-mass energy = 63 GeV in the Axial Field Spectrometer (AFS)\textsuperscript{1,2} at the CERN Intersecting Storage Rings (ISR). The photons are observed as electromagnetic showers in two added high-granularity NaI γ detector walls\textsuperscript{3}. A previous experiment at 12 GeV has reported an excess below p\textsubscript{T} of 50 MeV in the c.m. We examine our data for such an excess.

The important parts of the AFS apparatus for this investigation were the central drift chamber subtending rapidity ±1.0, and the NaI γ detector. The drift chamber operated in an axially symmetric magnetic field of 0.1 T in the low-field configuration and at 0.5 T for the high-field configuration. The momentum resolution in the low-field configuration is given by \( \delta p/p = 0.1p \), where p is in GeV/c. The drift-chamber information was used in track studies of minimum-ionizing showers in the NaI detector.

The shower detector consisted of two walls of high-granularity NaI crystals read out with vacuum photodiodes. Each wall covered a solid angle of 0.6 sr and consisted of 600 crystals in a 30 (vertical) by 20 (horizontal) matrix. The 13.6 cm (5.3 radiation lengths) long NaI crystals are 3.5 cm × 3.5 cm at the front, tapering to 4.0 cm × 4.0 cm at the rear, arranged so that a particle from the intersection region traversed a roughly constant amount of NaI, only weakly dependent on the angle of incidence.

The NaI detector was calibrated: i) through the study of e\textsuperscript{0}′s and γ′s from the two-γ decay mode; ii) through comparison of electron energy deposition in the NaI with electron momenta measured in the drift chamber; and iii) through the energy deposition of minimum-ionizing particles in individual crystals. The response of the NaI alone, which varied significantly with energy for the low-energy γ′s of interest to us, was determined to within 10% by the consistency of the various methods.

The data for this analysis were collected in 'minimum-bias' runs, with a loose trigger corresponding to an inelastic collision with two charged particles in the final state. A sample of 220,000 events was taken, 140,000 with the AFS field at 0.1 T (low field) and 80,000 with the standard (high-field) AFS configuration. The low-field sample corresponds to an integrated luminosity of 5.9 \times 10\textsuperscript{32} cm\textsuperscript{-2}; the high-field sample to 3.1 \times 10\textsuperscript{32} cm\textsuperscript{-2}. A comparable sample of 'apparatus empty' events was accumulated, with the equipment randomly strobed. These events were used to control for apparent γ′s in the NaI detector, faked by the electronics or other non-event-associated sources.

To reduce cosmic-ray and beam-gas interactions, a vertex with two charged tracks was required in the intersection diamond. No second interaction was allowed within ±15 ns of the nominal event time.

Since we normalized the spectra of photons to the number of charged tracks observed in our detector, it is important that we measure the track spectrum carefully. We calculated the efficiency and acceptance curves for tracks and pairs by using a detailed simulation of our detector. Minimum-bias generated events were generated according to the measured spectra of Alper et al.\textsuperscript{4} and Guettler et al.\textsuperscript{5}, and followed inside the detector taking into account all the usual processes, such as dE/dx, multiple scattering, nuclear interactions, bremsstrahlung, and pair production. Digitizings corresponding to the different parts of the detector, identical in format to real data, were analysed by the same programs that analysed the real data.

Observed tracks were required to have at least 30 degrees of freedom, a \( \chi^2 \) per degree of freedom of less than 10, rapidity within ±0.8, \( \cos(\theta) < 0.7 \), and an absolute value of the azimuth outside the range 1.1–2.0 to avoid blind spots in the drift chamber. Electrons were removed from the track sample by a dE/dx cut in the drift chamber (single tracks) or invariant mass less than 20 MeV (pairs). Two tracks satisfying the above cuts were required, consistent with the trigger.

The threshold to trigger a shower definition was 10 MeV. The cut defined a lower limit for the energy of recognized showers; a cut of 15 MeV was used in the final analysis. Photon candidates were required to be within a fiducial volume of ±53 cm vertical, by ±34 cm horizontal at 113.4 cm from the centre of the interaction diamond. Showers which matched a track (within an ellipse with axes 8 cm vertical by 16 cm horizontal) were excluded from consideration.

Measurement of the photon spectrum includes removal of spurious showers

\[
\text{True observed photons} = N_{\text{obs}} - N_{\text{noise}} - N_{\text{unseen tracks}},
\]

where

\[
N_{\text{tot}} = \text{total observed photons},
\]
\[
N_{\text{noise}} = \text{apparent photons from electronic noise method},
\]
\[
N_{\text{unseen tracks}} = \text{showers associated with charged tracks where the charged tracks are not seen or the photons are not properly associated with the charged tracks}.
\]
Subtraction (or exclusion) of false photons from electronic sources was important for photons with an energy in the laboratory system of less than 100 MeV. This important background was removed in two different ways, and the consistency between the two methods, of approximately 10% per energy bin, is an indication of our systematic errors. In the first method, suspect 'noisy' blocks are removed from the data, and full event processing is carried out both for the data and for corresponding data taking with a fake 'empty' apparatus from the data with an out-of-time trigger. These 'empty' events then yield an apparent photon/interaction spectrum characteristic of the electronic noise sources. This noise spectrum, normalized to the total number of data events, is subtracted from the photon spectrum in the data. We have checked, via a detailed shower shape study of the data and empty event photons, that the empty subtraction method seems to be reliable in subtracting the apparent noise events in the data.

The second method of removing the noise events is to make a cut on the radius of accepted showers. Two variations of this method were tried. In the first variation, events with very large radii photons were removed from the sample. In the second, the large-radii photons were removed and the photon spectra were corrected for losses from the radius cuts, using the expected shower shape from the EGS to determine the losses. For the different methods, the corrections are quite different fractions of the spectra in the low-energy bins with which we are concerned, but the subtracted results were consistent to within 10% of the final spectrum.

The remaining background is from showers associated with charged tracks in the NaI, for which the tracks have not been recognized in the drift chamber. The consistency between the two walls is particularly useful for this correction, since the energy associated with the unseen tracks is constant in the laboratory frame, but comes at relatively lower energies in the c.m. in wall 2, which is toward the direction of the c.m. motion.

A final source of systematic error is shower splitting by the reconstruction program. For the results here we add nearby showers, assuming that they arise from reconstruction-program imperfections, and correct geometrically for chance overlaps. The opposite extreme—assuming no splitting in the residual spectra—has also been tried as a limiting case. This extreme method does not change the spectrum appreciably and changes the overall number of photons by less than 5%. Neutrons are a small contamination, estimated at less than 2% of our observed photon spectrum.

Photons are expected to come primarily from decays of \( \pi^0, \eta, \omega, \) etc. We assume isospin invariance and calculate our \( \pi^0 \) spectrum from the yield of \( \pi^+ \) and \( \pi^- \): \( \pi_0 = (\pi^+ + \pi^-)/2 \). As for the charged-track spectra in the NaI method, which extend up to 2 GeV/c, the spectra were joined at approximately 600 MeV/c, with an upward shift of the Alper parametrization by 10%, required by our charged-track spectra and consistent with quoted normalization uncertainties. This shift has a negligible effect on our final results, less than 10% of overall systematic uncertainties. To the photons coming from \( \pi^0 \)'s one has to add the contribution of sources that do not obey the isospin invariance or do not proceed through an intermediate \( \pi^0 \) state: for example, \( \eta \rightarrow \gamma \gamma \) and \( \eta \rightarrow 3\pi^0 \rightarrow 6\gamma \). Shower-finding efficiency and NaI response determined for our apparatus, using the EGS (Fig. 1), are applied to this true \( \gamma \) spectrum. The resulting \( \gamma \) spectrum is then renormalized, using the ratio of our observed charged tracks to the calculated charged-track spectrum from the Monte Carlo calculation.

![Fig. 1](image)

Fig. 1 a) Efficiency of shower reconstruction. b) Fraction of the total energy observed in the NaI.
Fig. 2 Data points compared with predictions based on the extremes of the band in Fig. 1b.

Limits on the NaI response come from the consistency of the NaI electron calibration, with the NaI response inferred from the $E_{\text{NaI}}/E_{\text{tot}}$ ratio observed for the high-energy $\gamma$'s from $\pi$ and $\eta$ decay. The NaI electron calibration, for electrons in our energy range but with a slightly different detector configuration than that used for data-taking, is in rough agreement with the EGS prediction, while the $E_{\text{NaI}}/E_{\text{tot}}$ ratio from the higher-energy triggered data indicates a higher response in the data than that found from the EGS. We take the midpoint of the corrections suggested by the two methods as the nominal NaI response, with limits given by the band in Fig. 1.

Figure 2 is a comparison of the observed and predicted photon spectra, normalized to charged tracks. The two curves shown arise from the limits in our uncertainty of the NaI response, to which the predicted observed photon spectrum is quite sensitive. An overall 11% normalization uncertainty has not been included in the curves shown. The error bars on the data points reflect systematic uncertainty estimates.

The total observed photons show an excess compared to prediction of $13 \pm 15\%$. The excess in the lowest bin is $(15 \pm 15)\%$. In a bin kinematically comparable to this, Chliapnikov et al. found an excess of $(30 \pm 6)\%$. Thus we can exclude, at the $2\sigma$ level, substantial growth of this excess between $\sqrt{s} = 12$ GeV and $\sqrt{s} = 63$ GeV.

REFERENCES

REGGE TRAJECTORIES IN THE QUARK MODEL

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ABSTRACT

We prove that the large angular momentum behaviour of the leading Regge trajectory of a meson (q̅q) or a baryon (qqq) can be obtained by minimizing the classical energy of the system for given angular momentum. A two-body quark-antiquark linear potential plus relativistic kinematics produces asymptotically linear Regge trajectories for mesons. For baryons we take either a sum of two-body potentials with half strength or a string of minimum length connecting the quarks, and find in both cases that the favoured configuration is a quark-diquark system and that the baryon and meson trajectories have the same slope. Short-distance singularities of the potential are shown to be unimportant.
1. - INTRODUCTION

In the 1960's, everybody knew about Regge poles. Actually they were thought to explain everything. Now they are out of fashion but they still exist! In the crossed channel they manifest themselves to describe two-body reactions. For instance, the difference between the $p\bar{p}$ and $pp$ total cross-sections is well described from $E_{lab} = 10$ GeV to $E_{lab} = 2000$ GeV by

$$\sigma_{p\bar{p}} - \sigma_{pp} = \text{const.} \cdot E^{-0.55}$$

corresponding to the exchange of the $p$ trajectory$^1)$. In the direct channel it is striking to see trajectories of mesons going up to $J = 6$ and of baryons up to $J = 17/2$. These trajectories are remarkably linear and approximately parallel, i.e., the angular momentum is a linear function of the square of the particle mass.

In the 60's this fact was taken for granted by those who applied Regge's ideas to particle physics even though there was not the faintest justification for this! Regge's original work was done with non-confining potentials, and, at very large energies, trajectories turned around in the complex $J$-plane to end up at some negative integer.

Now we know that hadrons have a composite structure, and that mesons are quark-antiquark pairs and baryons three-quark systems. The use of a potential interaction has met with considerable success in the description of hadrons, especially those made of heavy quarks, but also to a certain extent those made of light quarks. There is more and more support for the belief that the quark-antiquark potential is linear at a large distance, and, to the extent that the quark-antiquark system can be regarded as a relativistic string for large angular momentum, the slope of the Regge trajectory has been connected with the string tension. Here we shall do something slightly heretical and regard the potential as producing an instantaneous interaction between relativistic point-like quarks.

In the case of baryons, the most natural prescription is to take two-body forces with $V_{QQ} = 1/2 V_{Q\bar{Q}}$. However, for large separations there are good reasons to believe that the potential energy between three quarks is proportional to the length of the string of minimum length connecting the three quarks. We shall in fact consider both cases.
Several authors have proposed models in which the baryons are made of a quark-diquark system. Then in particular the parallelism of the meson and baryon trajectories becomes very natural. However, we still have to understand why it might be so.

Although the ground states and low-lying excited states of baryons have been well studied, the excited states with large angular momentum have only been touched upon and do not lend themselves so easily to numerical study. The main remark of the present paper is that, as for mesons — where, as we shall see, it is completely clear — the leading Regge trajectory of a baryon (i.e., the sequence of ground-state wave functions and energies with increasing angular momentum \( J \)) can be obtained in the large \( J \) limit, by minimizing the classical energy of the system for given \( J \). Quantum effects only play a role in preventing the collapse of a subsystem caused by short-range singularities of the potential. For a linear two-body potential and for a string, one finds that the configuration minimizing the energy is a quark-diquark system. With relativistic kinematics one proves that the trajectories tend to become linear, and, unavoidably, since the colour of a diquark system in a baryon has to be a \( \bar{3} \), the potential energy is the same as in a meson and the trajectories become parallel.

2. — THE TWO-BODY CASE: RELATIVISTIC KINEMATICS

For simplicity we take only the extreme relativistic case, i.e., the Hamiltonian is, taking \( c = 1 \),

\[
H = \frac{p_1^2}{2m} + \frac{p_2^2}{2m} + V(r_{12})
\]

(1)

or, in the c.m. system,

\[
H = 2p + V(r)
\]

(2)

Provisionally we restrict ourselves to a purely linear potential,

\[
V(r) = \lambda r
\]

(3)

Here the minimum of the classical energy for a given angular momentum will give a lower bound for the energy of the ground state. In fact, we can give an explicit proof by generalizing an inequality of Herbst, who proves the operator inequality
to be valid for any state. If we restrict ourselves to states of angular momentum $J$, we have the following inequality:

$$\langle J|\hat{p}|J \rangle \geq 2 \left( \frac{J(J+1)}{F(J+\frac{1}{2})} \right)^{\frac{1}{2}} \langle J|\frac{1}{r}|J \rangle \geq (J+\frac{1}{2}) \langle J|\frac{1}{r}|J \rangle \quad (4)$$

Therefore, the quantum energy of a system of two particles of momenta $\hat{p}$ and $-\hat{p}$ will satisfy the inequality

$$E_Q(J) \geq \text{Im} \int \left\{ \frac{2J+1}{r^2} + V_{\text{lin}}(r) \right\} \quad (5)$$

while the minimum of the classical energy is

$$E_c(J) = \int_{\text{lin}} \left\{ \frac{2J}{r^2} + V_{\text{lin}}(r) \right\} . \quad (6)$$

So if we take the linear potential (16) we get

$$E_c(J) = 2\sqrt{2} \sqrt{J} \sqrt{J} \quad (7)$$

and hence, if we believe that this gives the leading behaviour of the quantum ground-state energy for large $J$, then

$$J(t) \sim \frac{1}{8\lambda} t + \cdots \quad (8)$$

We can prove that Eq. (7) indeed gives the leading behaviour by bounding above the Hamiltonian (2), by using the operator inequalities

$$\hat{r} \leq \frac{1}{2} \left( \frac{\hat{p}^2}{\hat{x}} + x \right) \quad (9)$$

$$\hat{x} \leq \frac{1}{2} \left( \frac{\hat{r}^2}{\hat{r}} + y \right) \quad (10)$$

and get

$$2\sqrt{2} \sqrt{J} (J+\frac{1}{2}) \leq E(\chi, J) \leq 2\sqrt{2} \sqrt{J} \left( 2n + J + \frac{3}{2} \right)^{\frac{1}{2}}. \quad (11)$$
A more precise but less rigorous estimate is

\[ t(J) \approx 4(J + \frac{1}{2}) + 4\sqrt{2} \alpha + 2\sqrt{2} + \ldots \]  

(12)

This means that daughter trajectories are asymptotically parallel to the leading trajectory.

3. - THE THREE-BODY CASE

For the interaction between the three quarks constituting a baryon, we have taken two extreme cases:

i) A sum of two-body interactions adjusted in such a way that if two quarks are close to one another, the potential between the quark-diquark system is identical to the quark-antiquark potential,

\[ V = \frac{a}{2} \left[ V(r_{12}) + V(r_{23}) + V(r_{13}) \right] \]  

(13)

where \( V \) is the \( qq \) two-body potential, and in the special case (3)

\[ V = \frac{a}{2} \left( r_{12} + r_{23} + r_{13} \right). \]  

(14)

There is no rigorous justification for this choice. The rule (28) holds for one-gluon exchange contributions. The other remark is that a diquark system has colour \( \bar{3} \) and therefore looks like an antiquark. Finally, let us indicate that experience has shown that the application of this rule to the calculation of ground-state energies of baryons has met with remarkable success, for instance, in the calculation of the \( Q^- \) mass by J.-M. Richard.

ii) The potential energy could be proportional to the minimum length of a Y-shaped string connecting the three quarks. Again the strength is adjusted in such a way that when two quarks coincide it agrees with the quark-antiquark potential,

\[ V = \lambda \inf_{P} (r_{1P} + r_{2P} + r_{3P}). \]  

(15)
There are good reasons for believing that Eq. (15) holds for large separations between the quarks. For instance, this is what one gets in the static Born-Oppenheimer bag model.

We shall assume that, like in the two-body case, the classical energy gives a lower bound of the ground state energy. Therefore, we wish to minimize the classical energy.

If \( J \) is the total angular momentum

\[
\vec{J} = \vec{p}_1 \times \vec{p}_1 + \vec{p}_2 \times \vec{p}_2 + \vec{p}_3 \times \vec{p}_3, \quad (\vec{p}_1 + \vec{p}_2 + \vec{p}_3 = 0),
\]

we notice that if we project the points 1, 2 and 3 and momenta \( \vec{p}_1, \vec{p}_2 \) and \( \vec{p}_3 \) on a plane perpendicular to \( \vec{J} \), then \( \vec{J} \) is unchanged, the kinetic energy is reduced, and for monotonous two-body potentials as well as for the string interaction (15) the potential energy is reduced. Hence we can restrict ourselves to a motion of the three particles in a plane.

We shall not give the details of the minimization in the plane, which is a very amusing geometrical exercise. One ends up with two possibilities:

1) A quark-diquark configuration (which is the only possibility for the string) with an energy

\[
E = 2p + \frac{3}{2} \rho
\]  (16)

and \( J = p \rho \) where \( p \) is the impulsion of the isolated quark and \( \rho \) the quark-diquark distance. Hence, minimizing,

\[
E = 2\sqrt{2} \sqrt{\rho} \sqrt{J}
\]  (17)

2) An equilateral triangle, formed by the three quarks, with an energy

\[
E = 3p + \frac{3}{2} 3\sqrt{3} R
\]  (18)

and \( J = 3pR \), when \( R \) is the distance of the quarks to the centre.

Hence
now since $3^{3/4} > 2$ (because $27 > 16$), we see that the quark-diquark configuration is always favoured.

Using the same type of majorizations, as in the two-body case:

$$2 J c^2 < A + \frac{A c^2}{A}, \quad 2 r_{ij}^2 < B + \frac{r_{ij}^2}{B},$$

it is easy to prove that the classical minimum gives the leading behaviour of the ground state energy for large $J$.

The conclusion is that both baryon and meson trajectories have the same slope asymptotically:

$$J(t) \sim \frac{t}{8a}$$

The fact that the coefficient does not agree with what one gets from the Nambu-Goto string should not be worrying. What matters is the comparison between the baryon and meson cases.

For references we draw the reader's attention to the CERN preprint TH.4259/85 (1985) to appear in Zeitschrift für Physik C. However, since then I have had communications of explicit calculations in the meson case, by S. Godfrey [Phys. Rev. D31 (1985) 2375] who uses the same Hamiltonian as us, and by H. Baacke who uses a Dirac or a Klein-Gordon equation (private communication at this conference). Both find "precocious" linearity of the trajectories, as do the calculations of Schnitzer, Preparata and collaborators, Basdevant and Boukraa.
CHARM HADROPRODUCTION

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ABSTRACT

A brief review of charm hadroproduction results is presented. In order to avoid the confusion that has long irked this subject the discussion is devoted solely to single proton target data.
It is becoming quite difficult to write an original article on charm hadroproduction. At least two reviews are written per annum [1] and yet new results are accumulated at a rate which would justify a review article with something new to say perhaps every 5 years. For reasons noted below, the next review article written on this subject ought to make a serious and useful contribution. This article, alas, can do little more than revisit some old data, emphasise a practical point of view and hold out the promise of some new and useful data to be published later this year.

After a brief discussion of the A-dependence problem in which we underline our reasons for using only single-proton target data in this article, we give the present situation on the charm total cross-section versus energy, charm differential cross-section properties and $D\bar{D}$ correlation data.

1. The A-Dependence Problem

One of the most annoying problems that has plagued the interpretation and synthesis of charm hadroproduction data is the so-called A-dependence problem. For several practical reasons, most charm hadroproduction data have been accumulated using a variety of nuclear targets (Be, C, Fe, W, Freon, etc.). The only simple way to compare these data with each other and with hydrogen and collider data is to extrapolate all data to a selected target material; usually we choose a single proton as the selected target. Unfortunately it is not known how to perform the extrapolation. It is usually (incorrectly [2]!) assumed that $\sigma(A) = \sigma(p).A^\alpha$, where $\sigma(A)$ and $\sigma(p)$ are the nuclear and proton cross-sections, $A$ is the nuclear mass number and $\alpha$ is a parameter to be determined experimentally. Several attempts have been made to measure $\alpha$, under different conditions, with conflicting results. For example, for light flavour production, Barton et al [2] have shown that $\alpha$ is a strong function of $x_F$ and Cronin et al [3] have demonstrated that $\alpha$ depends on $p_\perp$. $J/\psi$ production supports $\alpha = 1$ [4] and, traditionally, it has been assumed that $\alpha = 1$ (independent of $x_F$ and $p_\perp$) for open charm hadroproduction as well [1,5]. Fermilab experiment E613 [6] using a
beam dump technique and assuming that the neutrinos leaving the dump come from charm particle semi-leptonic decays obtains $\alpha = 0.75 \pm 0.05$ independent of neutrino energy. In fact, $\alpha = 1$ is strongly vetoed by E613; comparing the neutrino production rate with Be and W targets, they obtain $0.91 \pm 0.13$ while $\alpha = 1$ would predict 2.3.

The practical point of view noted above, therefore, is that whereas heavy target data can give valuable information on charm particle properties (such as masses, lifetimes, branching ratios, etc.), there is no reliable way to use these data to investigate charm hadroproduction characteristics. Under these circumstances we prefer to restrict ourselves to data collected with a single proton target. This is especially important for the cross-section results discussed in Sections 2 and 3(b).

At this point it is worth emphasising that this situation should improve within a short time. Fermilab experiment E769 [7] plans to measure the A-dependence of charm hadroproduction. Since E769 will use the same equipment and the same technique as the very successful E691 [8] we are optimistic that charm A-dependence data will soon exist.

2. Charm Total Cross-sections

With the restriction noted in the previous section, Table 1 gives all charm production cross-section measurements. Even though these data are the most reliable available, there are still problems. The BNL-MPS experiment [9] did not see any charm signal and can only quote a 95% confidence level upper limit. The LEBC-EHS experiments NA16 [10], NA27 [11] and E743 [12] all suffer from very low statistics charm samples. (Note that the E743 samples are labelled (charm)$^{\pm}$ and (charm)$^{0}$. This is because the E743 results are based on a topological analysis and it is conceivable that both samples are contaminated. Comparison with the NA27 data indicates that both are D samples, however.) The ISR experiments SFM [13] and SFM-CBF [14] had to apply quite tight selection criteria to see any evidence of charm and therefore their results are affected by enormous correction.
factors. For the pion induced data (first 9 entries in Table 1) all samples have been fully analysed and there is not much likelihood of the Table 1 situation improving; in the near future at least. The proton situation is much healthier. The NA27 results are based on an analysis of only \( \approx 25-30\% \) of the available data; the E743 results are based on \( \approx 5\% \). In both these experiments new results will be published later this year. Both experiments should achieve small total cross-section errors (\( \approx 1-2\% \)). Both SFM experiments took large quantities of new data just before the demise of the ISR. At least one of these \[15\] plans new results later this year. Thus the proton situation of Table 1 can only get better. This is important, since with the present data all we can say is that \( \sigma \) increases with \( \sqrt{s} \) and that has been known for almost a decade. Details of the excitation curve await the new data.

One final point worth a mention concerns the F-meson. None of the experiments contained in Table 1 have reported any evidence for the F. The NA27 collaboration performed a sophisticated analysis based on charged particle identification \[11\] and obtained the 90\% confidence level upper limit given in Table 1 (the errors are systematic).

3. Charm Differential Cross-sections

a) \( P_{\perp} \) Behaviour

The only charm hadroproduction characteristic that is apparently unaffected by nuclear effects is the \( P_{\perp} \) behaviour. All charm particles appear to be produced according to \( e^{-aP_{\perp}^2} \) where \( a \approx 1 \text{ (GeV/c)}^{-2} \). In other words, the average \( P_{\perp} \) is \( \approx 1 \text{ GeV/c} \) \[1\]. The \( P_{\perp} \) effect seen in light flavour production is quite subtle \[3\] and presumably the present statistical significance of hadroproduced charm is inadequate to detect such effects.

b) \( X_F \) Behaviour

Two problems have caused confusion when comparing \( X_F \) distributions from different experiments. If the A-dependence is as complicated for charm production as it is for light flavour production \[2\] then heavy nucleus data samples have
$x_F$ distributions which are distorted by nuclear effects. For this reason, again we only consider single proton target data in this section.

The second problem originates in the form of the function that is used in different experiments to fit the $x_F$ distribution. In principle one ought to fit the invariant distribution $E d\sigma/dx_F \sim (1-x_F)^n$ to obtain the exponent $n$. However, in low statistics experiments this procedure is not always practical and experimenters tend to fit $dN/dx_F \sim (1-x_F)^n$ to obtain $n$. Thus one cannot easily compare the values of $n$ obtained and published by different experimenters. It turns out that this problem is avoided when one uses only the data of the experiments listed in Table 1 since all these experiments use the second form. However, it is a general problem to bear in mind.

Back in 1981 [16], the SFM-CBF group published evidence of a leading hadron effect in charm production. This was manifested as a rather flat $x_F$ distribution ($n=0.4$) for $\Lambda_c$ production in pp collisions. The NA16 and NA27 collaborations [17] have reported a similar, if more complicated, effect in D-meson production by incident $\pi^-$. Both these LEBC experiments observe an $x_F$ distribution which does not fit well to $(1-x_F)^n$. Dividing the data into the so-called leading states ($D^-, D^0, D^{*-}, D^{*0}$) and non-leading states ($D^+, \bar{D}^0, D^{*+}, \bar{D}^{*0}$) both experiments report a striking difference in the respective $x_F$ distributions. Fitting them to $(1-x_F)^n$, NA16 obtains $n=1\pm1$ and $6\pm3$ for leading and non-leading samples and NA27 obtains $n=1.8^{+0.6}_{-0.5}$ and $7.9^{+1.6}_{-1.4}$. The NA27 collaboration has analysed the $D^*$ sample separately and find $n=4.3^{+1.8}_{-1.5}$ independent of quark content. The suggestion from all this is that there are several charm hadroproduction mechanisms. Charm particles containing a quark in common with the beam particle tend to have flat $x_F$ distributions. All other charm particles (including all $D^*$ states ??) are produced $\approx$ centrally ($n \geq 4$).

Apart from the SFM-CBF result, the $x_F$ distributions for pp data indicate central charm production with none of the structure suggestive of multiple production mechanisms that we have in the $\pi p$ data. Neither NA16 nor NA27 have
any indication of a leading $\Lambda_c$ component in the pp data.

4. Charm Pair Correlations

NA16 and NA27 are the only experiments to date which have enough $D\bar{D}$ pair events to study correlations [19]. All correlation properties investigated (effective mass, $x_F(D\bar{D})$, $p_{\perp}(D\bar{D})$, rapidity gap, $D\bar{D}$ transverse angle $\phi_{\perp}$) are in good agreement with a simple model which assumes no dynamic correlation between $D$ and $\bar{D}$. The only possible exception to this simple picture is a curious structure at $\phi_{\perp} \approx 120^\circ$ in the $\phi_{\perp}$ distribution. Far more statistics are needed before this effect achieves respectability.

5. Conclusions

From an experimental viewpoint, charm hadroproduction is very tricky indeed. Nevertheless the situation is beginning to stabilise and we are starting to make sense of the data. It is unfortunate that most of the available data has been collected with heavy nuclear targets and it is not known how to compute single nucleon cross-sections from these data. Thus, in this review, we choose to consider only single proton target data to investigate charm hadroproduction characteristics; at least that way we know what we have measured and what we are studying.

With this major simplification the main messages are the following: the total charm cross-section increases with $\sqrt{S}$, but we don't yet know how fast; D-mesons are produced with a mean $p_{\perp} \approx 1$ GeV/c; charm $x_F$ distributions indicate a variety of hadroproduction mechanisms; the small amount of $D\bar{D}$ pair data is consistent with there being no dynamic correlation between $D$ and $\bar{D}$. Probably the most important message is that more data will be on the market later this year.
Table 1

Charm Cross-sections ($x_F \geq 0.0$)

<table>
<thead>
<tr>
<th>Experiment</th>
<th>Beam</th>
<th>Charm</th>
<th>$\sqrt S$(GeV)</th>
<th>$\sigma(\mu b)$</th>
<th>Comment</th>
</tr>
</thead>
<tbody>
<tr>
<td>BNL-MPS</td>
<td>$\pi^-$</td>
<td>$D^{*-}$</td>
<td>5.5</td>
<td>$\leq 0.13$</td>
<td>All $x_F$</td>
</tr>
<tr>
<td>NA16</td>
<td>$\pi^-$</td>
<td>$D^\pm$</td>
<td>26</td>
<td>4.5$^{+2.2}_{-1.4}$</td>
<td></td>
</tr>
<tr>
<td>NA16</td>
<td>$\pi^-$</td>
<td>$D^0/\bar{D}^0$</td>
<td>26</td>
<td>7.7$^{+7.2}_{-3.5}$</td>
<td></td>
</tr>
<tr>
<td>NA27</td>
<td>$\pi^-$</td>
<td>$D^\pm$</td>
<td>26</td>
<td>5.7$\pm 1.6$</td>
<td></td>
</tr>
<tr>
<td>NA27</td>
<td>$\pi^-$</td>
<td>$D^{*\pm}$</td>
<td>26</td>
<td>5.0$^{+2.3}_{-1.8}$</td>
<td></td>
</tr>
<tr>
<td>NA27</td>
<td>$\pi^-$</td>
<td>$D^0/\bar{D}^0$</td>
<td>26</td>
<td>10.1$\pm 2.2$</td>
<td></td>
</tr>
<tr>
<td>NA27</td>
<td>$\pi^-$</td>
<td>$D^{*0}$</td>
<td>26</td>
<td>7.3$\pm 2.9$</td>
<td></td>
</tr>
<tr>
<td>NA27</td>
<td>$\pi^-$</td>
<td>$F^{\pm}$</td>
<td>26</td>
<td>$\leq 0.75\pm 0.15$</td>
<td>$\tau \geq 3 \times 10^{-13}s$</td>
</tr>
<tr>
<td>NA27</td>
<td>$\pi^-$</td>
<td>$\Lambda_c^+ / \bar{\Lambda}_c^+$</td>
<td>26</td>
<td>$\approx 4^{+5}_{-3}$</td>
<td></td>
</tr>
<tr>
<td>NA16</td>
<td>p</td>
<td>$D^\pm$</td>
<td>26</td>
<td>5.3$^{+2.4}_{-1.6}$</td>
<td></td>
</tr>
<tr>
<td>NA16</td>
<td>p</td>
<td>$D^0/\bar{D}^0$</td>
<td>26</td>
<td>10.2$^{+7.9}_{-4.3}$</td>
<td></td>
</tr>
<tr>
<td>NA27</td>
<td>p</td>
<td>$D^\pm$</td>
<td>26</td>
<td>6.2$\pm 1.7$</td>
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<tr>
<td>NA27</td>
<td>p</td>
<td>$D^0/\bar{D}^0$</td>
<td>26</td>
<td>16.2$\pm 4.2$</td>
<td></td>
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<tr>
<td>E743</td>
<td>p</td>
<td>(charm)$^\pm$</td>
<td>39</td>
<td>15$\pm 7$</td>
<td>Preliminary</td>
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<td>$\approx 100 \pm 50$</td>
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<tr>
<td>SFM</td>
<td>p</td>
<td>$\Lambda_c^+$</td>
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<td>6.5$\pm 3.2$</td>
<td>$\sigma$.BR;all $x_F$</td>
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<tr>
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<td>$D^0$</td>
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<td>63</td>
<td>$\approx 3.6 - 84.0$</td>
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References


Tagging of heavy flavours with dimuons at the CERN pp$^-$ Collider

H.-G. Moser\textsuperscript{1)}

Abstract: We report the observation of 512 dimuon events below the $Z^0$ mass with the UA1 detector at the CERN pp$^-$ collider. Most of the events are interpreted as semileptonic decays of heavy flavours, mainly b$\bar{b}$. We can calculate a cross section for b$\bar{b}$ of 1.2 $\mu$b for a $p_t$ of both b's above 5 GeV/c. Isolated dimuon events can be explained as T decays and muon pairs from the Drell–Yan process.

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1. Introduction

Muons have the unique advantage that they can be detected within hadronic activity. Therefore they can be used to tag $c^-$ and $b^-$ mesons by their semileptonic decays. In principle this can be done with single $\mu$ events, but there one has to fight against a large background from pion and kaon decays. Dimuon events have a much smaller rate, but here the background turns out to be acceptably small. In addition to muon pairs from heavy flavour decays one expects dimuons from $T$ decays and from the Drell Yan continuum, all with dimuon masses well below the $Z^0$ mass. The results on these intermediate mass dimuons from the 1983 run have already been published [1]. This presentation is based on data from the 1983, 1984 and 1985 runs.

2. Detector and Data – Selection

The UA1 detector has been described in detail elsewhere [2]. In brief, the muon momentum is measured in the central detector using a 0.7 T dipole field. It then has to pass more than 9 interaction lengths of calorimeters and iron shielding. The muon is finally detected in the muon chambers surrounding the central calorimeters. Muon events are triggered by a special hardware trigger. We collected 108 nb$^{-1}$ at $\sqrt{s} = 546$ GeV in 1983, 256 nb$^{-1}$ in 1984 and 330 nb$^{-1}$ in 1985, both at $\sqrt{s} = 630$ GeV.

For this analysis we selected dimuon events with the following cuts:

$$p_T(\mu) > 3 \text{ GeV/c for each muon}$$

$$M^{\mu\mu} > 6 \text{ GeV/c}^2$$

This yielded 732 events which were scanned in order to remove background such as cosmic ray muons, leakage and kinks (kaon decays), leaving 512 dimuon events (excluding $Z^0$ decays).
3. Background Studies

Several possible sources of background have been studied. The 9 interaction lengths of absorber make the background from punchthrough and shower leakage very small (less than $10^{-4}$ per incident hadron, from test beam measurements). The most severe background arises from decays of $\pi$'s and K's. The probability of a $\pi$ (K) to decay in flight in the CD is $0.02/p_t$ ($0.11/p_t$). Assuming that 58% (21%) of the charged particles are $\pi$'s (K's) the probability to fake a muon is $0.04/p_t$ per charged hadron. To calculate this background we used an inclusive muon sample with $p_t > 3$ GeV/c. Events having a second high $p_t$ particle were selected. Assuming this particle to be a pion or a kaon we simulated its decay into a muon. Applying our cuts we summed up the probabilities of these events to fake a dimuon. Thus we obtained a background estimate of 107 events.

4. Interpretation of the Events

Besides the $Z^0$ decay there are three main sources of dimuon events:

- The Drell-Yan mechanism
- Decays of $T,T',T''$

These processes give rise to unlike sign muon pairs where the muons are isolated (not accompanied by hadrons).

- Semileptonic decays of heavy quarks, which are produced by QCD processes (e.g. gluon fusion).

This process gives predominantly rise to unlike sign dimuons as the b decays in a $\mu^-$ and the $\bar{b}$ in a $\mu^+$ (similarly c$\bar{c}$). Like sign dimuons can be produced by the second generation decay of one quark: The b may decay first to c d $\bar{u}$ and then the c into a $\mu^+$. However, muons from second generation decays have a softer $p_t$ spectrum and are therefore suppressed by our $p_t$ cut.
There is another possible source of like sign events from first generation through $B^0 - \bar{B}^0$ mixing (analogous $K^0 - \bar{K}^0$ mixing). From calculations no significant mixing in the $B_d$ system is expected, but mixing in the $B_s$ system can be large [3].

Heavy flavours can also be produced by the decay of Intermediate Vector Bosons ("IVB") e.g.: $Z^0 \to b\bar{b}$. These events should lead to muon pairs accompanied by high $E_t$ jets; the $\mu^- - \mu^- - \text{jet} - \text{jet}$ mass should be approximately the mass of the IVB.

The muons produced in heavy flavour events are accompanied by hadrons from the fragmentation of the heavy quark and from the decay products of its semileptonic decay. The muons are therefore normally not isolated.

The most straightforward way to separate dimuons due to heavy flavour decays from Drell–Yan and $T$ events is to look at the isolation. In order to determine the isolation we sum the $E_t$ in a cone of $\Delta R < 0.7$ around the muon ($\mu$: $\Sigma E_t, \Delta R^2 = \Delta \eta^2 + \Delta \phi^2$). We call an event isolated if $(\Sigma E_t(1))^2 + (\Sigma E_t(2))^2 < 9 \text{ (GeV)}^2$. This can be shown by plotting $(\Sigma E_t(1))^2 + (\Sigma E_t(2))^2 = S$, see Fig. 1. For like sign events which have no contribution from Drell–Yan and $T$ the distribution is flat, but unlike sign events show a clear excess at $S < 9 \text{ (GeV)}^2$. Examining $W$ and $Z^0$ decays into $\mu$ we believe this cut to be 82% efficient. We observe 98 isolated unlike sign events with $S < 9 \text{ (GeV)}^2$. There is also some contribution from heavy flavour events to the isolated sample. We have tried to separate these different contributions by fitting their dimuon–mass distributions to the data. The results of this fit is shown in Figure 2. Together with acceptance calculations performed with the ISAJET Monte Carlo program [4] we can calculate cross sections for Drell–Yan and $T$:

$$\sigma_{DY}(m_{\mu\mu} > 11 \text{ GeV/c}^2) = 299 \pm 89 \text{ pb}$$

$$\sigma(pp \to T, T', T'' \to \mu^+ \mu^-) = 887 \pm 217 \text{ pb}$$

These cross sections agree well with low energy measurements of Drell–Yan and $T$ production (Fig. 3). The Drell–Yan cross section is also in very good agreement with the theoretical value of 270 pb by Altarelli et al. [5].
Figure 1: Isolation of muons: $(\Sigma E_t(1))^2 + (\Sigma E_t(2))^2 = S \Sigma E_t$ is the sum of $E_t$ in a cone of $\Delta R < 0.7$ around the muon.

Figure 2: Dimuon mass distribution for isolated events (unlike sign)
The fitted contributions from Drell-Yan, T and heavy flavour decays are indicated.

The non-isolated events (257 unlike sign events and 142 like sign events) and the 15 isolated like sign events are expected to come from heavy flavour decays. Monte Carlo studies with ISAJET and EUROJET [8] [4] have shown that most of these events come from $b\bar{b}$ production (80%). The remainder is due to $c\bar{c}$, which is suppressed because of the softer fragmentation of the $c$-quark. The muon $p_t$ spectrum is shown in fig. 4 together with the EUROJET prediction. For comparison the inclusive muon $p_t$ spectrum is also shown. The data are in good agreement with the Monte Carlo calculations. We also plot the $\mu - \mu$ - jet - jet mass distribution. The nearest jet to each muon in $\Delta R < 1$ is used to calculate the mass (if no jet is found the vector
Figure 3: Comparison of UA1 measurements with Drell–Yan (a) and T (b) cross-sections from low energy experiments. [6] [7]

sum of energy in $\Delta R < 0.7$ around the muon is used as ‘jet’). This is also compared with the Monte-Carlo calculation which describes the data fairly well. This allows us to calculate the $b\bar{b}$ cross section: Correcting for acceptance, track finding efficiency and background subtraction we find:

$$\sigma(pp \to b\bar{b}) = 1.2 \pm 0.1 \pm 0.2 \mu b$$

$$p_t > 5 \text{ GeV/c for each quark}$$

$$|\eta| < 2.0 \text{ for each quark}$$

(systematical error for luminosity and acceptance only) According to EUROJET 66% of the cross section comes from $2 \to 2$ processes, while 34% comes from $2 \to 3$ processes. ISAJET does not use exact matrix elements for higher order processes, but it includes similar processes like gluon splitting in heavy quarks which contributes to 22% of the cross section.
Figure 4: Properties of non isolated events

a) $p_T$ spectrum of muons in dimuon events and inclusive muon $p_T$ spectrum (1983 data). Also shown are the EUROJET predictions for both spectra.

b) $\mu^+\mu^- -$jet-$-$jet mass distribution for non isolated events. The curve is a Monte Carlo prediction from $cc$ and $bb$ production, normalized to the number of events.

Both Monte Carlo reproduce our data very well, with the exception of the high number of like sign events we observe. Using the non isolated events and subtracting the calculated background we get a ratio of like sign to unlike sign events of $0.46 \pm 0.07 \pm 0.06$. The systematical error includes errors of the background calculation and from possible misassociations of muon tracks to central detector tracks. The ratio we expect from second generation decays of $b-$quarks is $0.18 - 0.25$. But if we include $B_0^0 - \bar{B}_0^0$ mixing the expected ratio is $0.40$ to $0.50$. These numbers are very suggestive, nevertheless more work is needed to understand the reliability of the Monte Carlo results and their dependence of the choice of parameters (e.g. fragmentation models, the probability to create a $B_s$ meson etc.). A more detailed discussion can be found in [9].
5. Acknowledgments

I would like to thank Carlo Rubbia, Karsten Eggert and my colleagues of the UA1 experiment who made this work possible. Special thanks go to M. Jimack, R. Edgecock, N. Ellis and B. Van Eijk for their help and advice in preparing this contribution.

References

Abstract

Lowest order QCD (gluon-gluon and quark anti-quark fusion in order $\alpha^2_S$), predicts that heavy flavour production should fall off rapidly away from the central region. We calculate the process $g + q \rightarrow Q + \overline{Q} + q$ and present analytic results for the matrix element squared. This process is expected to give the dominant contribution in proton nucleon scattering along the direction of the incoming valence quark. After factorisation of the regions of collinear emission into the lowest order processes, the residual $O(\alpha^2_S)$ contribution has a small effect in the forward region. The production of heavy flavours calculated using perturbation theory is thus expected to be predominantly central. Other mechanisms, which may lead to non-central production of charmed quarks, are expected to fall off like a power of the heavy quark mass.
I. Introduction

It is common knowledge that perturbative QCD gives a poor description of the hadronic production of charmed particles at fixed target energies. For example, total cross-sections predicted by lowest order QCD lie below experimental results at fixed target energies, perhaps by more than an order of magnitude\(^1\). Lowest order QCD also predicts predominantly central production of charmed hadrons, whereas experiments may indicate significant production in the forward region\(^2\).

Before concluding that hadronic charm production is a failure of the theory, it is important to consider the mitigating circumstances. First and foremost it is questionable whether charm production is in fact described by perturbative QCD alone, because the mass of the charmed quark is not much heavier than the scale of the strong interactions. It should also be noted that experimental results are often based on a small number of events observed in a limited kinematic region. The estimation of total cross-sections from these results, requires large acceptance corrections, the size of which may depend sensitively on the model used. For data obtained from scattering on nuclear targets, comparison with theory requires an additional assumption about the atomic number dependence of the cross-sections.

It is also important to remember the imprecision of the perturbative QCD prediction itself. The 'standard' perturbative QCD formula for the inclusive charm production,

\[ H(P_1) + H(P_2) \rightarrow Q(P_3) + X \]  

is given by,

\[
\frac{E_3}{d^3 P_3} d\sigma = \sum_{i,j} \int dx_1 dx_2 \left[ \frac{E_3 d\sigma_{ij}(\alpha S(\mu^2), P_1, P_2)}{d^3 P_3} \right] f_i(x_1, \mu^2) f_j(x_2, \mu^2)
\]  

The functions \(f\) are the distribution functions of light partons (gluons, light quarks and anti-quarks) evaluated at a scale \(\mu\), which is of the order of the mass of the produced heavy quark. \(\hat{\sigma}\) is the short distance cross-section from which the mass singularities have been factored in the normal way\(^3\). Since the sensitivity to collinear emission has been removed from the short-distance cross-section, \(\hat{\sigma}\), it is calculable as a perturbation series in \(\alpha_S(\mu^2)\). The lowest order which contributes is \(O(\alpha_S^3)\). In this order there are contributions to \(\hat{\sigma}\) due to gluon gluon fusion and quark anti-quark annihilation. At fixed target energies the lowest order perturbative predictions, obtained using Eq.(1.2), depend sensitively on the input parameters, most notably the mass of the charmed quark. For example, at \(\sqrt{S}=27\) GeV, we find that the total cross-section changes by more than an order of magnitude as we vary the mass of the charmed quark between 1.2 and 1.8 GeV. Such uncertainties also afflict the predictions for the longitudinal momentum distributions of the charmed quark.
In addition to the standard formula, Eq. (1.2), it has been suggested in the literature that the following mechanisms might contribute significantly to charmed particle production.

1. Flavour excitation graphs which contribute because of the presence of charmed quarks in the wave function of the incoming hadrons. The charmed quark content of the nucleon can be calculated using perturbation theory or may be due to non-perturbative mechanisms, in which case it is said to be intrinsic.3

2. Diffractive production of a charmed quark pair from a gluon in one of the hadrons.6

3. Recombination of a produced charmed quark with a fast quark in one of the beam jets.6

4. Final state pre-binding distortion caused by the binding of charmed quarks to light quarks.6

The question which we wish to address in this paper is whether or not heavy flavours (bottom, top, etc.) will be copiously produced in the forward direction. This issue is relevant not only for present experiments but also for the design of detectors for future hadron colliders. All arguments presented here will be based on perturbation theory. Thus, strictly speaking we are considering the production of a heavy flavour, whose mass is very much bigger than the strong interaction scale ($m >> \Lambda$). Examination of the dangerous regions of phase space in low order QCD diagrams indicates that the additional mechanisms enumerated above are not relevant for the production of very massive quarks.10,11 They are either already included in the standard factorisation formula or suppressed by at least a power of the heavy quark mass. A definite confirmation of the validity of Eq. (1.2) for heavy quark production will require an all-orders proof, but the arguments of Collins, Soper and Sterman make it plausible that the QCD improved parton model provides a reliable description of the hadronic production of heavy quarks. We therefore expect that the factorisation formula, Eq. (1.2) will be useful if we can neglect terms of order $\Lambda/m$ where $\Lambda$ is the QCD scale and $m$ is the heavy quark mass. From a practical point of view it is also necessary to require that $\ln(m/\Lambda) >> 1$, so that the hard scattering cross-section is accurately represented by a limited number of terms in its perturbative expansion.

The crucial issue is whether the charmed quark is massive enough to be considered a heavy quark in the sense described above. The charm quark mass lies in the range $m = 1.2 - 1.8$ GeV so it is obviously a borderline case. A full answer to the question of whether the charmed quark can be considered a heavy quark requires theoretical control of the terms which vanish as a power of $m$ as well as the perturbative corrections, (the calculation of which is initiated here), which only vanish as a logarithm of the mass. No definitive answer to this question can be given here.

In this paper we use perturbation theory in order $g^6$ to calculate heavy quark production using the standard formula Eq. (1.2). We consider the process

$$g + q \rightarrow Q + \overline{Q} + q$$ (1.3)
The theoretical interest of this subprocess stems from the fact that it can be viewed as 'containing' both the flavour excitation of a heavy quark and the diffractive production of a heavy quark pair as subgraphs. The explicit calculation carried out here indicates that flavour excitation does not contribute in this order in perturbation theory. We find that the only singular regions correspond to the quark anti-quark and gluon-gluon fusion mechanisms. Our work can thus be regarded as a confirmation of the arguments of ref.(11). From a more practical point of view the higher order corrections provide theoretical information on the best choice of factorisation and renormalisation scale. The motivation for considering process of Eq.(1.3) (without considering all the other processes of order $g^6$) is that this process is expected to give the dominant contribution in the forward region in pN collisions. This is a consequence of the stiffness of the valence quark distribution in the proton in comparison with the anti-quark or gluon distributions. A priori, we expect that when $(1 - x_F) \approx \alpha_s$ the process of Eq.(1.3) will be competitive with the lowest order process.

In Section (IV) we present numerical results for charmed quark production in the forward region. This is not because we can demonstrate that perturbation theory, (neglecting terms which vanish as a power of the charmed quark mass), gives a reliable prediction for charmed particle production in the forward region. It is rather our intention to make the perturbative prediction as precise as possible, in order to assess the significance of charmed particle production in the forward region.

II. Heavy Flavour Production in Lowest Order

The lowest order processes which contribute to the production of a heavy quark $Q$ are the so-called fusion processes,

\[
\begin{align*}
\text{(a) } & q(p_1) + \bar{q}(p_2) \rightarrow Q(p_3) + \bar{Q}(p_4) \\
\text{(b) } & g(p_1) + g(p_2) \rightarrow Q(p_3) + \bar{Q}(p_4)
\end{align*}
\]

The four momenta of the partons are given in brackets. These processes have been thoroughly investigated in refs.4,12,13]. However in order to regulate mass singularities we will need the forms of these matrix elements in $n = 4 - 2\epsilon$ dimensions\textsuperscript{14]. We have therefore repeated the calculation of these matrix elements.

The Feynman diagram for quark anti-quark annihilation is shown in Fig. 1.

![Feynman diagram](image)

Figure 1: Lowest order diagram for the process $q + \bar{q} \rightarrow Q + \bar{Q}$.

The result for the matrix element squared, summed (averaged) over final (initial) colours and
spins can be expressed in terms of the transition probability $T_{qq}$,

$$
\sum |M^{(a)}|^2 = g^4 \mu^4 T_{qq}(\tau_1, \tau_2, \rho, \epsilon)
$$

$$
T_{qq}(\tau_1, \tau_2, \rho, \epsilon) = \frac{V}{2N^2} \left( \tau_1^2 + \tau_2^2 + \frac{1}{2} \rho - \epsilon \right)
$$

where the dependence on the $SU(N)$ colour group is shown explicitly, ($V = N^2 - 1, N = 3$).

Because this transition probability will ultimately be inserted in a QCD improved parton model formula we have chosen to express it in terms of variables which have simple behaviour under rescaling of the incoming momenta $p_1$ and $p_2$.

$$
\tau_1 = \frac{m^2 - t}{s} = \frac{p_1.p_3}{p_1.p_2}, \quad \tau_2 = \frac{m^2 - u}{s} = \frac{p_2.p_3}{p_1.p_2}, \quad \rho = \frac{4m^2}{s}
$$

and $s = (p_1 + p_2)^2, t = (p_1 - p_3)^2$ and $u = (p_2 - p_3)^2$.

The Feynman diagrams for the gluon-gluon fusion process are shown in Fig. 2. It is convenient to divide the result for the transition probability into two pieces,

$$
\sum |M^{(b)}|^2 = g^4 \mu^4 T_{gg}(\tau_1, \tau_2, \rho, \epsilon)
$$

$$
T_{gg}(\tau_1, \tau_2, \rho, \epsilon) = T_{gg}^{(1)}(\tau_1, \tau_2, \rho, \epsilon) + T_{gg}^{(2)}(\tau_1, \tau_2, \rho, \epsilon)
$$

As before these results have been averaged and summed over initial colours and spins. The average over the spin of the initial gluons has been performed by dividing by $n - 2 = 2(1 - \epsilon)$. This is in agreement with the normal convention used in the calculation of the two loop anomalous dimensions$^{15,16,17}$. The results are,

$$
T_{gg}^{(1)}(\tau_1, \tau_2, \epsilon) = \frac{1}{2V N(1-\epsilon)} \left( \frac{V}{\tau_1 r_2} - 2N^2 \right) \left( \tau_1^2 + \tau_2^2 - \epsilon \right)
$$

$$
T_{gg}^{(2)}(\tau_1, \tau_2, \rho, \epsilon) = \frac{1}{2V N(1-\epsilon)^2} \left( \frac{V}{\tau_1 r_2} - 2N^2 \right) \left( \rho - \frac{\rho^2}{4\tau_1 \tau_2} \right)
$$

In the limit $\epsilon \to 0$ the factorised form in Eqs.(2.4, 2.5) agrees with the more complicated expression given in ref.(4).
The one parton inclusive cross-sections are determined by the transition probabilities given above,

$$\frac{p^2 d\sigma_{ij}}{d^{n-1} p_3} = \mathcal{N} \frac{\alpha_s^2}{s^2} T_{ij}(\tau_1, \tau_2, \rho, \epsilon) \delta(1 - \tau_1 - \tau_2)$$  \hspace{1cm} (2.6)

\(\mathcal{N}\) is an overall normalisation factor which is equal to one in four dimensions. Using the expression for the parton cross section given in Eq.(2.6), the QCD prediction for the process,

$$H(P_1) + H(P_2) \rightarrow Q(P_3) + X$$  \hspace{1cm} (2.7)

can be cast in the characteristic parton model form,

$$\frac{E_3 d\sigma}{d^{n-1} P_3} = \frac{\mathcal{N} \alpha_s^2(\mu^2)}{S^2} \sum_{i,j} \int \frac{dx_1}{x_1} \frac{dx_2}{x_2} f_i(x_1, \mu^2) f_j(x_2, \mu^2) T_{ij}(\tau_1, \tau_2, \rho, \epsilon) \delta(1 - \tau_1 - \tau_2)$$  \hspace{1cm} (2.8)

In this equation the parton variables are expressed in terms of their hadronic counterparts as follows,

$$\tau_1 = \frac{\tau_1^H}{x_2}, \quad \tau_2 = \frac{\tau_2^H}{x_1}, \quad \rho = \frac{\rho^H}{(x_1 x_2)}$$  \hspace{1cm} (2.9)

where,

$$\tau_1^H = \frac{P_1 \cdot P_3}{P_1 \cdot P_2}, \quad \tau_2^H = \frac{P_2 \cdot P_3}{P_1 \cdot P_2}, \quad \rho^H = \frac{4m^2}{S}, \quad S = (P_1 + P_2)^2$$  \hspace{1cm} (2.10)

The phenomenological implications of these lowest order estimates for the short distance cross-sections have been investigated elsewhere in the literature. We list the salient features here,

1. They predict cross-sections which are predominantly central.
2. The rate of fall-off away from the central region is controlled by the stiffness of the distributions of gluons and light (anti-) quarks in the incoming hadrons.
3. The average transverse momentum of the heavy quarks is of the order of the mass of the produced heavy quark and the net transverse momentum of the pair of heavy quarks is small.

Parenthetically, we note that the number of charmed particles present in a jet has also been the subject of both experimental and theoretical analysis. As explained above the lowest order diagrams are not expected to contribute significantly in the region \(p_T \gg m\). In this region, charmed particles are much more likely to be the fragmentation products of a gluon jet which is produced by the normal scattering between light partons. The higher order correction calculated in this paper can also be viewed as a contribution to this process.
final colours and spins. The result is denoted by $C$.

$$\sum |M^{(C)}|^2 = g^6 \delta^6 \epsilon C(p_1, p_2, p_3, p_4, p_5)$$  \hspace{1cm} (3.3)

By examination of the graphs of Figure 3, we see that graph 1 contains a mass singularity which can be identified as a contribution to the quark anti-quark fusion process of Eq.(2.1). In a similar way diagrams 3, 4 and 5 contain singularities which contribute to the gluon-gluon fusion process. Of course it is only in a physical gauge that the mass singularities be ascribed to particular diagrams in this way. Our first step is to isolate the contribution of these singular regions in the matrix element squared $C$, calculated using the full gauge invariant set of diagrams given in Fig.3. In the limit $p_1.p_5 \rightarrow 0$, the full $n$ dimensional matrix element becomes,

$$C(p_1, p_2, p_3, p_4, p_5) \rightarrow$$

$$C^{[15]}(p_1, p_2, p_3, p_4, p_5) = \frac{1}{p_1.p_5} \left[ T_{qq}(r_1, 1 - r_1, \rho, z_1, \epsilon) \frac{1}{z_1} P_{q\bar{q}}(z_1, \epsilon) \right]$$ \hspace{1cm} (3.4)

$P_{q\bar{q}}$ is the splitting kernel\(^ {20}\) in $n$ dimensions and describes the perturbative probability of finding an anti-quark in the incoming gluon,

$$P_{q\bar{q}}(z, \epsilon) = \frac{T_R}{1 - \epsilon} \left( z^2 + (1 - z)^2 - \epsilon \right), \quad z_1 = \frac{r_2}{1 - r_1}, \quad T_R = \frac{1}{2}$$ \hspace{1cm} (3.5)

In order to investigate the limit $p_2.p_5 \rightarrow 0$, it is convenient to move to the centre of mass system of the $\overline{Q}q$ system ($\vec{p}_4 + \vec{p}_5 = 0$). In this frame we may write,

$$p_4 = \frac{(s_{45} + m^2)}{2\sqrt{s_{45}}} (1, \ldots, -\lambda \sin \theta_1 \cos \theta_2, -\lambda \cos \theta_1), \quad \lambda = \frac{s_{45} - m^2}{s_{45} + m^2}$$

$$p_5 = \frac{(s_{45} - m^2)}{2\sqrt{s_{45}}} (1, \ldots, \sin \theta_1 \cos \theta_2, \cos \theta_1)$$ \hspace{1cm} (3.6)

In Eq.(3.6) the dots represent $n - 3$ components of momenta which are determined by the mass-shell conditions $p_4^2 = m^2$, $p_5^2 = 0$. We have introduced the notation, $s_{45} = (p_4 + p_5)^2$. In the same system we may choose,

$$p_1 = \frac{(s_{45} - u)}{2\sqrt{s_{45}}} (1, \ldots, \sin \alpha, \cos \alpha)$$

$$p_2 = \frac{(s_{45} - \epsilon)}{2\sqrt{s_{45}}} (1, \ldots, 0, 1)$$ \hspace{1cm} (3.7)

In Eq.(3.7) the dots represent $n - 3$ zero momentum components. $p_3$ is determined by overall energy-momentum conservation. The limit $p_2.p_5 \rightarrow 0$ corresponds to the region of collinear
emission from the incoming quark with momentum \( p_2 \), \( (\theta_1 \to 0) \). In this limit we obtain,

\[
C(p_1, p_2, p_3, p_4, p_5) \rightarrow
\]

\[
C^{[25]}(p_1, p_2, p_3, p_4, p_5) = \frac{1}{p_2 \cdot p_5} \left[ T_{qg}(1 - \tau_2, \frac{\rho}{z_2}, \frac{\epsilon}{z_2}) \frac{1}{z_2} P_{gq}(z_2, \epsilon) \right]
\]

\[
+ T_{gq}^{[2]}(1 - \tau_2, \frac{\rho}{z_2}, \frac{\epsilon}{z_2}) \left( 2(1 - \epsilon) \cos^2 \theta_2 - 1 \right) \frac{V}{N} \frac{1 - z_2}{x_2^2}
\]

where \( P_{gq} \) is the splitting function\(^{20} \) in \( n \) dimensions.

\[
P_{gq}(z, \epsilon) = \frac{V}{2N} \left( \frac{1 + (1 - z)^2 - \epsilon z^2}{z} \right), \quad z_2 = \frac{r_1}{1 - \tau_2}
\]

Note that the second term in Eq.(3.8) vanishes after averaging over the angle \( \theta_2 \) (even in \( n \) dimensions) to reproduce the expected factorisation result for the inclusive cross-section.

By now the strategy of the calculation should be clear. The difference

\[
C(p_1, p_2, p_3, p_4, p_5) - C^{[15]}(p_1, p_2, p_3, p_4, p_5) - C^{[25]}(p_1, p_2, p_3, p_4, p_5)
\]

is perfectly finite in all regions of phase space and hence the limit \( n \to 4 \) can be taken. The contribution of the difference Eq.(3.10) to the inclusive charm production cross-section will be evaluated by numerical integration. The limit \( \epsilon \to 0 \) of \( C^{[15]} \) and \( C^{[25]} \) can be obtained from Eqs.(3.4, 3.8).
In four dimensions, using the momentum assignments of Eq. (3.2), $C$ becomes,

$$C(p_1, p_2, p_3, p_4, p_5) \rightarrow$$

$$\left[ \frac{2 \{23\}^2 + 2 \{24\}^2 + 2 \{35\}^2 + 2 \{45\}^2 + m^2(t_{25} + s_{34})}{2 t_{25} s_{34}} \right]$$

$$\times \left[ \frac{N^2 - 4}{4N^2} \right]$$

$$\left( \frac{2 \{24\}}{\{12\}\{14\}} - \frac{2 \{23\}}{\{12\}\{13\}} + \frac{2 \{35\}}{\{13\}\{15\}} - \frac{2 \{45\}}{\{14\}\{15\}} - \frac{\{25\}}{\{12\}\{15\}} - \frac{\{34\}}{\{13\}\{14\}} \right)$$

$$+ \frac{1}{4} \left( \frac{\{25\}}{\{12\}\{15\}} + \frac{\{34\}}{\{13\}\{14\}} - \frac{2 \{23\}}{\{12\}\{13\}} - \frac{2 \{24\}}{\{12\}\{14\}} - \frac{2 \{35\}}{\{13\}\{15\}} - \frac{2 \{45\}}{\{14\}\{15\}} \right)$$

$$+ \frac{N^2 - 4}{N^2} \frac{m^2}{s_{34} t_{25}} \left[ \frac{\{35\} + \{23\} - \{45\} + \{24\}}{\{14\}} \right]$$

$$+ \frac{V}{N^2} \frac{1}{s_{34} t_{25}} \left( \frac{\{12\}^2 + \{15\}^2}{\{12\}\{15\}} \right) + \frac{1}{2 s_{34}} \left( \frac{1}{\{12\}} - \frac{1}{\{15\}} - \frac{1}{\{13\}} \right)$$

$$+ \frac{m^2}{\{13\}^2} - \frac{1}{\{14\}} + \frac{4}{s_{34}} \Delta_1^2 + \Delta_2^2 + \Delta_3^2 + \Delta_4^2$$

$$+ \frac{1}{N^2} \frac{m^2}{s_{34} t_{25}} \left[ 1 + \frac{t_{25}}{\{13\}} - \frac{m^2}{\{14\}} + \frac{\{12\}^2 + \{15\}^2}{\{13\}\{14\}} \right]$$

$$- 2 \left( \frac{\{24\} + \{45\}}{t_{25}} \right) \frac{\Delta_1 + \Delta_2}{t_{25}} - 2 \left( \frac{\{23\} + \{35\}}{t_{25}} \right) \frac{\Delta_3 + \Delta_4}{t_{25}}$$

(3.11)

In this equation we have introduced the notation,

$$s_{34} = (p_3 + p_4)^2, \ t_{25} = (p_2 - p_5)^2, \ p_j \cdot p_k = \{j,k\}$$

$$\Delta_1 = \frac{p_2 \cdot p_4}{p_1 \cdot p_3} - \frac{2p_1 \cdot p_2}{s_{34}}, \ \Delta_2 = \frac{p_4 \cdot p_5}{p_1 \cdot p_3} - \frac{2p_1 \cdot p_5}{s_{34}}$$

$$\Delta_3 = \frac{p_2 \cdot p_3}{p_1 \cdot p_4} - \frac{2p_1 \cdot p_2}{s_{34}}, \ \Delta_4 = \frac{p_3 \cdot p_5}{p_1 \cdot p_4} - \frac{2p_1 \cdot p_5}{s_{34}}$$

(3.12)

The differences $\Delta_1, \Delta_2, \Delta_3$ and $\Delta_4$ vanish in the limit $t_{25} \rightarrow 0$ as a square root of $t_{25}$. This is sufficient to reduce all the apparent double poles of $t_{25}$ in Eq. (3.11) to single poles. This result, which holds also in $n$ dimensions, is already evident from the explicit result given in Eq. (3.8). In the limit $m = 0$, Eq. (3.11) is in agreement with ref. (21). Numerical results for this process including masses have been given in ref. (22). All analytic results were obtained using the algebraic manipulation program Schoonschip (23).

We must now evaluate the contribution to the cross-section of the subtraction terms $C^{[15]}$ and $C^{[25]}$. These pieces contain singularities and will be evaluated analytically in $n$ dimen-
sions. The \( n \) dimensional phase space for the process in Eq. (3.2) can be cast in the form

\[
(PS)^{(3)} = \frac{(2\pi)^2}{2 \pi^4} \int \frac{d^{n+1}p_3}{E_3} \int d\tau_{45} \, \delta(1 - \tau_1 - \tau_2 - \tau_{45})
\]

\[
\left( \frac{\tau_{45}}{\tau_{45} + \rho/4} \right) \left( \frac{\tau_{45} + \rho/4}{\tau_{45}^2} \right)^{\epsilon} \left( \frac{4\pi}{\epsilon} \right)^{\epsilon} \frac{\Gamma(1 - \epsilon)}{\Gamma(1 - 2\epsilon)} I
\]

(3.13)

In this equation \( n - 4 \) irrelevant angles have been integrated over and \( \tau_{45} = p_4.p_5/p_1.p_2. \)

The dependence on the angles \( \theta_1 \) and \( \theta_2 \), which are defined in the reference frame given in Eq. (3.6, 3.7), is contained in \( I \).

\[
I = \frac{1}{2\pi} \int_0^\pi d\theta_1 \sin^{-2\epsilon} \theta_1 \int_0^\pi d\theta_2 \sin^{-2\epsilon} \theta_2
\]

(3.14)

After angular integration the contribution of \( C^{[5]} \) and \( C^{[25]} \) to the inclusive charm production cross-section at the parton level can be cast in the factorised form,

\[
p_3^0 \frac{d^{n-1}\sigma_{ij}}{d^{n-1}p_3} = \sum_{i'j'} \int_0^1 dz_1 dz_2 \left\{ p_3^0 \frac{d^{n-1}\tilde{\sigma}_{ij'}}{d^{n-1}p_3} \right\} \Gamma_{j'i'}(z_2, \epsilon) \Gamma_{i'i}(z_1, \epsilon)
\]

(3.15)

where the short distance cross-section \( \tilde{\sigma} \) is evaluated at rescaled values of the parton-momenta,

\[
\hat{p}_1 = z_1p_1, \quad \hat{p}_2 = z_2p_2,
\]

(3.16)

For the particular gluon quark process which we are calculating Eq. (3.15) becomes,

\[
p_3^0 \frac{d^{n-1}\sigma_{gg}}{d^{n-1}p_3} = \int_0^1 dz_1 \left[ p_3^0 \frac{d^{n-1}\tilde{\sigma}_{gg}}{d^{n-1}p_3} \right] \Gamma_{gg}(z_1, \epsilon) + \int_0^1 dz_2 \left[ p_3^0 \frac{d^{n-1}\tilde{\sigma}_{gg}}{d^{n-1}p_3} \right] \Gamma_{gg}(z_2, \epsilon) + \mathcal{O}(\alpha_s^4)
\]

(3.17)

The singularities present in the subtraction terms \( C^{[5]} \) and \( C^{[25]} \) now appear as poles in \( \epsilon \) in the functions \( \Gamma \). To first order in \( \alpha_s \) the singular parts of the functions \( \Gamma \) are given by the Altarelli-Parisi functions. At this order we may define the functions \( \Gamma \) to be,

\[
\Gamma_{ii'}(z, \epsilon) = \delta_{ii'} \delta(1 - z) - \frac{\alpha_s}{2\pi} P_{ii'}(z, 0) \left( \frac{1}{\epsilon} + \ln(4\pi) - \gamma_E \right)
\]

(3.18)

The association of the \( \ln(4\pi) \) and Euler constant with the pole in \( \epsilon \) defines this to be the \( \overline{MS} \) type of mass singularity factorisation\(^{24}\). This is the factorisation scheme used everywhere in this paper. As indicated in the previous sections, the parton cross-sections \( \sigma \) are calculated by averaging over \( n - 2 \) spin states of initial gluons.

The short distance cross-sections \( \tilde{\sigma}_{gg} \) and \( \tilde{\sigma}_{gq} \) are given at order \( \alpha_s^3 \) by Eq. (2.6). The
residual term $\hat{\sigma}_{gg}$ is perfectly finite and we may take the limit $n \to 4$. In this limit we obtain,

$$p_0^2 \frac{d^3 \hat{\sigma}_{gg}}{d^3 p_0} =$$

$$\alpha_s^2 \frac{\alpha_s}{2\pi} \frac{1}{(1 - t_1)} T_{gg}(t_1, t_2, z_1, \rho/s) \left[ \frac{P_{gg}(z_1, 0)}{z_1} Y + 2T_R(1 - z_1) \right]$$

(3.19)

$$+ \alpha_s^2 \frac{\alpha_s}{2\pi} \frac{1}{(1 - t_2)} T_{gg}(t_1, t_2, z_2, \rho/s) \left[ \frac{P_{gg}(z_2, 0)}{z_2} Y + \frac{V}{2N} \right]$$

where,

$$Y = \ln \left( \frac{(1 - t_1 - t_2)^2 s}{(1 - t_1 - t_2 + \rho/4)\mu^2} \right), \quad z_1 = \frac{t_1}{1 - t_1}, \quad z_2 = \frac{t_2}{1 - t_2}$$

(3.20)

The full answer for the short distance cross-section $\hat{\sigma}_{gg}$ is obtained by adding the contribution of Eq.(3.19) to the numerically evaluated contribution coming from the finite difference Eq.(3.10).

Before presenting numerical results we make some remarks of a more theoretical nature. By explicit computation we have shown that the only singular contributions in the full matrix element for process $C$ come from regions which can be associated with quark anti-quark annihilation and gluon-gluon fusion. In particular, in extracting the singular part of the full matrix element squared (Eq.(3.11) and its $n$ dimensional generalisation) we found that the double poles in $t_{25}$ vanish. The cancellation of the double poles can be simply demonstrated on a graph-by-graph basis working in the gauge $p_1.A = 0$. Naive power counting works in this gauge as long as the mass of the recoil system is greater than zero $(p_1 + p_2 - p_0)^2 < M^2 > 0$. This reduction of the double pole to a single pole has also been considered in ref.11.

Had the double poles persisted in the full answer they would have emphasized the low momentum region in heavy quark production. This would have suggested that mechanisms not described by perturbation theory were important for heavy quark production. It would also have lead to large production in the forward region.

Even after the cancellation of the double poles, the remaining $1/t_{25}$ term still displays a logarithmic sensitivity to the low momentum region. This sensitivity is the familiar one due to collinear parton emission and is removed by factoring the low momentum region $|t_{25}| < \mu^2$, $|t_{15}| < \mu^2$ into the incoming hadron wave-functions. In the remaining cross-sections all propagators are off-shell by at least $\mu^2 \sim m^2$. It is hence a bona fide higher order term in the short distance cross-section. There is no room left for a flavour excitation contribution.

IV. Numerical Results

In this section we present numerical results for charm production cross-sections using the standard parton model formula Eq.(1.2). We have used the parton distribution functions of
Table 1: Total cross-sections for the production of charmed quarks under various assumptions for the input parameters.

<table>
<thead>
<tr>
<th>√S (GeV)</th>
<th>σ_{gg} (µb)</th>
<th>σ_{q̅q} (µb)</th>
<th>σ_{TOT} (µb)</th>
</tr>
</thead>
<tbody>
<tr>
<td>27.0</td>
<td>1.2</td>
<td>0.16</td>
<td>1.3</td>
</tr>
<tr>
<td>PP, D01, Δ=0.2 GeV, m=1.8 GeV</td>
<td>17.</td>
<td>2.7</td>
<td>19.7</td>
</tr>
<tr>
<td>PP, D02, Δ=0.4 GeV, m=1.2 GeV, m_{th}= 1.8 GeV</td>
<td>9.0</td>
<td>.7</td>
<td>9.7</td>
</tr>
<tr>
<td>62.4</td>
<td>7.6</td>
<td>0.7</td>
<td>8.3</td>
</tr>
<tr>
<td>PP, D01, Δ=0.2 GeV, m=1.8 GeV</td>
<td>42.3</td>
<td>8.1</td>
<td>50.4</td>
</tr>
<tr>
<td>PP, D02, Δ=0.4 GeV, m=1.2 GeV, m_{th}= 1.8 GeV</td>
<td>26.7</td>
<td>3.0</td>
<td>29.8</td>
</tr>
<tr>
<td>630.0</td>
<td>85.</td>
<td>4.0</td>
<td>89.</td>
</tr>
<tr>
<td>PP, D01, Δ=0.2 GeV, m=1.8 GeV</td>
<td>169.</td>
<td>31.7</td>
<td>200.</td>
</tr>
<tr>
<td>PP, D02, Δ=0.4 GeV, m=1.2 GeV, m_{th}= 1.8 GeV</td>
<td>117.</td>
<td>15.5</td>
<td>132.5</td>
</tr>
</tbody>
</table>

Duke and Owens\textsuperscript{25}. These exist in two versions.

1. soft gluon distributions, Δ=0.2 GeV,(D01).

2. hard gluon distributions, Δ=0.4 GeV,(D02).

As a preliminary test we have investigated the sensitivity of the total charm production cross-section to variations in the input parameters. From Table 1 we see that the cross-sections depend sensitively on the charm quark mass. Because of smaller quark anti-quark luminosity it contributes only about 10% of the cross-section in proton-nucleon collisions. In calculating the total cross-section with a light quark mass (m=1.2 GeV) we have also investigated the effect of imposing a physical threshold (s > 4m_{th}^2). The resultant total cross-sections are shown in Table I at √S=27.0 GeV, √S=62.4 GeV and √S=630.0 GeV. Near threshold there is a considerable sensitivity to the value of the heavy quark mass which provides the hard scale for the interaction. These large variations in the total cross-sections should be borne in mind when looking at the more differential distributions which follow. At collider energies the sensitivity to the heavy quark mass is reduced, but the parton luminosities are uncertain because the value of x ≈ 2m/√S is so small.

The Feynman x_F distribution coming from the lowest order gluon gluon fusion and quark anti-quark annihilation are shown in Fig. 4. The cross-section falls steeply with x_F; by x_F = 0.5 it has fallen by about two orders of magnitude from its central value. Fig. 5 displays the result of including the higher order correction calculated in Section III. Numerical results were obtained using the adaptive Monte Carlo integration program Vegas\textsuperscript{26}. We have set the
Detailed results for the multiplicity of heavy quarks in a gluon jet have been given in ref.(18,19). Ref.(18) also includes estimates of the first non-perturbative contribution to gluon fragmentation into a heavy quark anti-quark pair. Although the presence of heavy flavours in jets may be of great experimental interest we do not consider this kinematic region further. We rather concentrate on the bulk of the cross-section which is produced at smaller transverse momenta of the order of the heavy quark mass.

III. Higher Order Corrections To Heavy Flavour Production

In order $g^6$ there are three types of process which contribute to the inclusive production of a heavy quark,

$$\begin{align*}
(A) & \quad q(p_1) + \bar{q}(p_2) \rightarrow Q(p_3) + X \\
(B) & \quad g(p_1) + g(p_2) \rightarrow Q(p_3) + X \\
(C) & \quad g(p_1) + q(p_2) \rightarrow Q(p_3) + X
\end{align*}$$

(3.1)

Processes $A$ and $B$ are radiative corrections to processes $a$ and $b$ in Eq. (2.1). A calculation of the order $g^6$ contributions of these processes would provide valuable information about the optimum choice for the scale $\mu$. This information would be of phenomenological importance especially for charm production. In order $g^6$ the processes of Eq.(2.1) receive both real and virtual corrections which separately contain both soft and collinear singularities. The calculation of these radiative corrections remains an open and challenging problem.

The calculation of process $C$ is much simpler since this process first appears at order $g^6$. In order $g^6$ perturbation theory the relevant parton process is,

$$g(p_1) + q(p_2) \rightarrow Q(p_3) + \bar{Q}(p_4) + q(p_5), \quad p_3^2 = p_4^2 = m^2$$

(3.2)

The Feynman graphs which contribute to this process are shown in Fig. 3. We have calculated the full result for the matrix element in $n$ dimensions averaged and summed over initial and
Figure 4: The differential cross-section \( \frac{d\sigma}{dx_F} \) in order \( \alpha_s^2 \) for the production of a charmed quark at \( x_F = 2p_T/\sqrt{S} \) at \( \sqrt{S} = 62.4 \) GeV. The total (solid curve) is comprised of the gluon gluon contribution (dashed curve) and the quark anti-quark contribution (dotted curve). The charmed quark mass is taken to be \( m = 1.2 \) GeV and the parton distributions DO2 were used.

Figure 5: The differential cross-section \( \frac{d\sigma}{dx_F} \) in order \( \alpha_s^3 \) (solid curve) and including the order \( \alpha_s^3 \) contribution from Eqs. (3.10, 3.19) (dotted curve). All other parameters as in Fig. 4.
scale of the hard interaction $\mu^2 = 4m^2$. After factorisation of the collinear singularities the resultant correction is always negative. Note that the quark-gluon contribution is expected to dominate only at large $x_F$ and that that Fig. 5 is therefore expected to be reliable, including effects of order $\alpha_s^3$, only at large values of $x_F$. In the central region it is possible that the gluon gluon process Eq.(3.1B) makes a large modification of the lowest order prediction. Fig. 5 indicates that perturbative QCD effects do not give large modifications of the lowest order prediction in the forward region in proton nucleon scattering. This result is in stark contradiction with earlier work in ref. 6) where large modifications of the lowest order results were found from perturbative flavour excitation type contributions. We believe that the discrepancy is due to the fact that the authors of ref. 6) include pole terms $(1/t^2)$ in the flavour excitation diagrams. Our complete gauge-invariant calculation shows that these terms are in fact cancelled.

It is amusing to note that process Eq.(1.3) introduces an asymmetry between the production of heavy quark or a heavy anti-quarks. Forming the combination,

$$C^{(-)}(p_1, p_2, p_3, p_4, p_5) = \frac{1}{2} [C(p_1, p_2, p_3, p_4, p_5) - C(p_1, p_2, p_4, p_3, p_5)]$$

we obtain in the same notation Eq.(3.12) as before,

$$C^{(-)}(p_1, p_2, p_3, p_4, p_5) = \frac{N^2 - 4}{4N^2} \left[ \frac{2}{t_{25}s_{34}} \left( \{23\}^2 + \{24\}^2 + \{35\}^2 + \{45\}^2 + \frac{m^2}{2} (t_{25} + s_{34}) \right) \right.$$

$$\left. - \frac{\{24\}}{\{12\}} \frac{\{23\}}{\{14\}} - \frac{\{35\}}{\{12\}} \frac{\{13\}}{\{14\}} + \frac{\{35\}}{\{13\}} \frac{\{15\}}{\{14\}} - \frac{\{45\}}{\{14\}} \frac{\{15\}}{\{13\}} \right] + 4m^2 \frac{s_{34}}{t_{25}} \left( \frac{\{35\}}{\{14\}} + \frac{\{23\}}{\{14\}} - \frac{\{45\}}{\{13\}} - \frac{\{24\}}{\{13\}} \right) \right) \]

In $p\bar{p}$ collisions at large values of the transverse momentum the processes of Eq.(3.1B,C) are expected to dominate. It could be that the mechanism of Eq.(4.2) induces an observable asymmetry between charmed and anti-charmed quarks in these processes.

**Acknowledgments**

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ABSTRACT  First results on $D^0$, $D^+$ and $F^+$ production from charm photoproduction experiment E691 at Fermilab are presented. The observed charm signals, based on a small fraction of data, are remarkably clean and large. The expected final statistics of charm particles exceed by more than an order of magnitude those of previous fixed target experiments. Analysis plans for immediate and more distant future are presented.

STATUS  —  The experiment was performed with the Fermilab Tagged Photon beam, $80 \ GeV < E_\gamma < 250 \ GeV$ and $5 \ cm$ Be target. We benefited a great deal from experience with the spectrometer from previous experiment, the E516. To reduce combinatorial backgrounds, which we found severe (despite a priori favourable in photoproduction charm to non-charm ratio), we have increased the magnetic field in the analysing dipoles (momentum resolution), added six new drift chamber planes (tracking efficiency) and installed new mirrors in Cerenkov counters (particle identification). The improvement in mass resolution resulting from spectrometer upgrades exceeds a factor of two. Also very importantly, we have taken advantage of advances in technology. We have added a high resolution silicon microstrip (SMD) vertex detector (transverse resolution of $20\mu m$, longitudinal of $300\mu m$). This allows to identify charm particles decay vertices and further reduce backgrounds. We used a wide
open $E_T$ (transverse energy) trigger accepting $\sim 40\%$ of the total hadronic cross section with $\sim 80\%$ efficiency for charm. It would have been very hard to decide to take so much data without the Advanced Computer Program 2) project at Fermilab, already well under way in 1985. It is a system of parallel 32-bit microprocessors, offering very large computing power (at low cost) for problems typical in high energy physics — how to reconstruct quickly very large samples of similar events. We have recorded $9 \times 10^7 E_T$ and $1 \times 10^7$ total hadronic cross section events, also special run data to study the $A$ dependence of $J/\psi$ photoproduction cross section ($H_2$, $Be$, $Fe$ and $Pb$ targets; $J/\psi \rightarrow \mu^+\mu^-$ trigger). We expect to process $\sim 20\%$ of full sample by summer 1986. With ACP we hope to finish reconstruction late this year (without ACP it would take $\sim 3$ years).

**FIRST RESULTS** — The signals $D^0 \rightarrow K^-\pi^+$, $D^+ \rightarrow K^-\pi^+\pi^+$, $D^0 \rightarrow \phi\pi^+$, $F^+ \rightarrow \phi\pi^+$ (charge-conjugate states are implicitly included) are shown in Figures 1-4. Extrapolation to the full data sample gives the following estimates of signals: $D^* \rightarrow D^0\pi$, $D^0 \rightarrow K^-\pi^+ - 3000$ events with no SMD cuts, $2000$ events with cuts as in Fig.2; $D^0 \rightarrow K^-\pi^+$ inclusive, with cuts as in Fig.1 - 4000 events; $D^+ \rightarrow K^-\pi^+\pi^+$ with cuts as in Fig.3 - 2500 events. We expect $> 8000$ events in $D^* \rightarrow D^0\pi$ (summing over three largest $D^0$ decay modes), and about 100-200 events per decay mode for $F$ and $\Lambda_c$.

**ANALYSIS PLANS** — $J/\psi$ cross section $A$ dependence — completion by June 1986.
— lifetimes — with 20% statistics available in summer 1986 we could measure $D^0, D^+, F^+$ lifetimes with precision exceeding the measurements done so far. With full statistics errors for $D^0, D^+$ lifetimes will reach the level of few percent, for $F$ and $\Lambda_c$ we expect to reach 10% precision.
— $\bar{D}^0 - D^0$ mixing — expected to be small 3) in the Standard Model offers a window to "new physics" 4). Clean $D^*$ tag and very good statistics allow sensitivity well below 1%.
— spectroscopy — high statistics should allow study the unseen decay modes of $F$ and $\Lambda_c$, other charm baryons, ratios of branching ratios; also a study of P-wave mesons with charm quark 5).
— cross sections — we could measure charm cross sections in the $E_T$ range 80–220 GeV. With careful Monte Carlo we should measure $x_F$, $y$, $p_T$ and other production characteristics, also $\bar{c} - c$ correlation in various variables for events with both charm particles identified.

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FIG. 1  Effective mass of $K\pi$ system. A signal of $\sim 120 \ D^0 \to K^-\pi^+ + c.c$ events is based on $\sim 3\%$ of data. Cuts applied: probability of $K\pi$ hypothesis based on Cerenkov information $P(K\pi) > 30\%$; separation between the primary and $D^0$ decay vertices $\Delta Z > 6\sigma$, (keeps $\sim 60\%$ of signal suppressing background by a factor of 80).
FIG. 2  Effective mass of $K\pi$ system. Cuts as in Fig. 1, with additional requirement that a $D^0$ comes from $D^* \rightarrow D^0 \pi$ decay. Signal of $\sim 30$ events from $\sim 3\%$ of data translates into 1000 events expected from the full sample.
FIG. 3  Effective mass of $K\pi\pi$ system. A signal of $\sim$ 90 events of $D^+ \to K^-\pi^+\pi^+ + c.c$ is based on $\sim$ 4% of data. Cuts applied: probability of $K\pi\pi$ hypothesis based on Cerenkov information $P(K\pi\pi) > 30\%$; separation between the primary and $D^+$ decay vertices $\Delta Z > 12\sigma_z$ (background suppression factor $\sim 300$).
Effective mass of $\phi^0(1020)\pi$ system. A signal of $\sim 10$ events in Cabibbo favoured mode $F^+ \to \phi^0\pi^+ + c.c.$ events is clearly seen, together with the Cabibbo suppressed mode $D^+ \to \phi^0\pi^+ + c.c.$ They come from $\sim 7\%$ of data. Cerenkov probability and vertex separation cuts applied. Extrapolation to the full data sample gives $\sim 150$ events in $F \to \phi^0\pi$ mode.
RESULTS FROM AN EVOLUTIONARY CUSB

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ABSTRACT

A progress report on the first BGO calorimeter in high energy physics, CUSB-II, is given. Preliminary results on the fine splitting of the $X_0'$ states obtained in early 1986, using CUSB-II, are presented. Included in here also are the recent limits obtained from CUSB-I on axions and from CUSB-I.5 on Higgs and gluinos.

INTRODUCTION

The CUSB-I detector was a NaI-Pb glass calorimeter, located in the North Area of CESR, which operated continuously from 1979 to mid 1984, mapping out during its lifetime: six $^3S_1$(b$b\bar{b}$), six $^3P_{2,1,0}$(b$b\bar{b}$) and two (b$d$) states$^1$. Recognizing that we have exhausted the possibilities of this fine instrument [despite it having achieved the designed resolution of $3.8%/E^X$(GeV)], in 1984 we began the first step of its upgrade by removing the inner strip chamber used for tracking and putting in the latter's place a quadrant of BGO crystals. In this configuration (CUSB-I.5) we ran at the peak of the T(1S) energy for an integrated luminosity of 22 pb$^{-1}$, and obtained a whole series of limits on new particles$^2$. In December 1985 we finished installing the whole complement 360 BGO crystals which constitute the heart of the new CUSB: CUSB-II. Early this
year we had a short run on the T(3S). I will begin by presenting the most recent results from CUSB-I and CUSB-I.5, then give a brief description and status report on CUSB-II, and end with a first look at the new data on the fine splitting of the $\chi_b$ states.

1. THE RECURRENT AXION

In 1978 Weinberg\textsuperscript{3)} and Wilczek\textsuperscript{4)} (WW) pointed out that the Peccei-Quinn\textsuperscript{5)} (PQ) mechanism proposed to avoid strong CP violation requires the existence of a pseudoscalar which they called "axion". In the original implementation of the WWPQ axion the PQ symmetry is supposed to be broken by the same VEV's which break the electroweak symmetry. In such a case the coupling of the axion to all fermions is uniquely determined up to an arbitrary parameter $X$ for up like quarks and $1/X$ for down like quarks and charged leptons, where $X$ is the ratio of the VEV of the two neutral Higgs in the model. The mass of this axion is $-75(X+1/X)$keV. Such axions were completely ruled out by searches in $J/\psi$\textsuperscript{6)} and $T$\textsuperscript{7)} radiative decays, independent of $X$.

The recent observation at GSI\textsuperscript{8)} of a $\approx 400$keV positron signal in heavy ion collision has been interpreted as the production and decay of an axion of mass $2m_e+2x400=1.8$ MeV. In the original WW model this corresponds to $X=0$ or 1/24, and a lifetime of the axion of $\approx 10^{-12}$s, a result which is still incompatible with $J/\psi$ and $T$ searches. Since the scale breaking of the PQ symmetry is unknown, lifetimes as low as $10^{-13}$sec are acceptable without creating inconsistencies with other pieces of physics such as the $g-2$ value of the electron. But in such case the $J/\psi$ and $T$ searches would have failed to observe the decay of $T(J/\psi)\rightarrow a+\gamma$ where the axion escapes undetected.

A search\textsuperscript{9)} for $T$ radiatively decaying into a short lived axion was made using CUSB-I. We obtain a negative result with varying sensitivity depending on the axion lifetime, see figure 1. For lifetimes of the order of $2x10^{-13}$sec where our experiment has the least sensitivity, the experimental limit for $B(T\rightarrow a\gamma)$ is $<2.5x10^{-3}$; while for such lifetimes assuming the WW model it is predicted to be of the order of 20%, totally in contradiction with our experiment. Thus, we again exclude the existence of an axion which couples to $b$ quarks with a strength greater than 10% of its coupling to electrons and down quarks.

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{figure1.png}
\caption{B(T$\rightarrow a\gamma$) versus Mean Decay Length (CUSB).}
\end{figure}
2. LIMITS ON THE HIGGS BOSON MASS

In the SM with only one Higgs doublet, there was a unique prediction for the branching ratio \( B(T \rightarrow \gamma + H) \), the Wilczek formula\(^{10}\). This has lately been modified by including first order radiative QCD corrections which reduced the expected branching ratio by \( \sim -\) a factor of two. CUSB\(^{11}\) had previously reported, assuming the Wilczek formula, exclusion of Higgs bosons of masses less than 4 GeV. With the QCD radiative corrections, the limit would barely hold at the 60% CL. However, if there were more than one Higgs doublet, the branching ratio is modified by the ratio of the VEV's \( \langle \phi_1 \rangle / \langle \phi_2 \rangle \). Therefore we reinterpret our limit on the Higgs' mass as a function of this ratio, figure 2.

3. LIMITS ON GLUINO MASS

Interest in supersymmetry continues unabated among physicists, despite complete lack of experimental evidence for the existence of sparticles. Searches for gluinos in particular have been performed at e\(^+\)e\(^-\) colliders and in beam dumps, yielding negative results, usually expressed as lower limits on the gluino mass in terms of some other sparticles\(^{12}\). At Kyoto, Farrar\(^{13}\) pointed out that while beam dump experiments are usually quoted for limits of \( m_\tilde{g} > 2 \) to 3 GeV, these limits hold true only for a small range of \( m_\tilde{q} \). Furthermore, for the low gluino mass region \( (m_\tilde{g} < 2 \) GeV), there was virtually no range of allowed \( m_\tilde{q} \) for which the \( m_\tilde{g} \) are excluded. We have remedied in good part this situation by having performed a search for low mass bound gluino states by searching for monochromatic \( \gamma \) signals. Our negative result translates into the range of excluded \( m_\tilde{g} \) shown in figure 3.

![Figure 2. Limits on the Higgs Mass.](image1)

![Figure 3. Gluino Masses Excluded by CUSB 1.5.](image2)
4. THE CUSB-II DETECTOR

For the second generation of T physics experiments, we needed a spectrometer with improved photon resolution and perfect projective geometry. To accomplish these purposes we need more radiation lengths ($\lambda_0$), no inactive materials, while retaining the longitudinal segmentation and source imbedding features of CUSB-I, occupying the radial space $10 \text{cm} < r < 25 \text{cm}$. These are possible using bismuth germanate crystals (BGO), because despite its lower light output than NaI, its better material properties, including that of being nonhydroscopic, allows us to fit $12 \lambda_0$'s of BGO with ideal geometry into the available space.

Each element has trapezoidal cross section and subtends $10^\circ$ in $\phi$ and $45^\circ$ to $90^\circ$ (or $90^\circ$ to $135^\circ$) in $\theta$. Five such elements, of increasing size, form one $\phi$-$\theta$ sector; Each crystal is viewed at one end by a miniature photomultiplier tube (p.m.), permanently glued on. All crystals are uniform, in light response, to better than 2% across their length (some after compensation, most as delivered from the Shanghai Institute of Ceramics). Teflon and aluminized mylar wrapping (less than three mils thick) maximize light output and remove optical cross talk. Figure 4 shows one polar half of the BGO assembly (180 elements) viewed from the free crystal end. One notes 5 free standing concentric rings, constructed in a time honored (Roman Arch) method. Figure 5 shows the crystal assembly viewed from the phototube end.

![Figure 4. Crystal End View of BGO.](image1)

![Figure 5. Phototube End View of BGO.](image2)

All CUSB calibration is done in real time during data taking with sources which are imbedded between crystal layers. The online calibration is accomplished by having a dual signal path, such that two different sensitivity channels measure collision events and source signals. The complete calibration cycle for all crystals is $\leq 30 \text{min}$. We have achieved $\leq 0.1\%$ channel to channel calibration and can achieve $\approx 1-2\%$ absolute calibration as checked with Bhabha events and shower leakage calculations. See next section.
5. PRELIMINARY RESULTS.

In figure 6 we show a first comparison of the relative bhabha energy resolution between CUSB-I and CUSB-II. Indeed the BGO array provided the expected improvement of a factor of two. A fit to this preliminary BGO spectrum yields a \( \sigma_E/E = 1.2\% \) for the upper half of the peak.

5.1 \( B_{\mu\mu} \) of the T(3S)

The second excited state of the bound (b\bar{b}) system, the T(3S) was first resolved in 1980 at CESR\(^{14}\). However, the total integrated luminosity accumulated on this resonance has been small, <15 pb\(^{-1}\). As a result, some fundamental properties such its branching ratio into a muon pair (\( B_{\mu\mu} \)) was poorly known. Using CUSB-II this spring we improved the knowledge on this quantity by \( = \) a factor of 5. We obtain:

\[ B_{\mu\mu}[T(3S)] = 0.016 \pm 0.003. \]

This value is preliminary until we finish our calculation of the systematic errors. This value is in agreement with that which is obtained from scaling using \( B_{\mu\mu}[T(1S)]^{15}\).

5.2 Fine Splitting of the \( \chi_b \)'s

The \( \chi_b \) states were discovered by CUSB-I in 1982\(^{16}\) from observations of the photons (\( \gamma_1, \gamma_2 \)) resulting from transitions between the triplet S and triplet P states:

\( 3^3S_1 \rightarrow \gamma_1 2^3P_2, 1, 0 \rightarrow \gamma_1 \gamma_2 1(2) 3^3S_1 \). Refer to level diagram, Figure 7. The \( \gamma_1 \)'s were observed in our inclusive photon spectra, and both \( \gamma_1 \) and \( \gamma_2 \) were

Figure 6. The Bhabha Energy Resolution of CUSB-I and CUSB-II.

Figure 7. Bound T(b\bar{b}) Level Diagram.
observed in events where the final $^3S_1$ states decayed into a pair of muons or electrons ("exclusive" events).

With CUSB-I we did not have sufficient resolution to resolve between the $\gamma_1$'s which came from the T(3S) to the J=2 and J=1 $\chi_b$ lines in either the inclusive photon spectrum or in the "exclusive" events. Neither were we able to see the $\gamma_2$'s in the inclusive photon spectrum.

From a preliminary analysis of our run on the T(3S) this spring with CUSB-II we obtained $\approx 30$ "exclusive" events. The scatter plot of the high energy $\gamma$ versus the low energy $\gamma$, figure 8, show two distinct clustering, the upper and lower clusters corresponding to T(1S) and T(2S) final states respectively. The projection of figure 8 onto the lower energy photon axis, figure 9, demonstrates visibly that now the $\chi_b$'s J=2 and J=1 lines are fully resolved.

The inclusive photon spectrum situation is complicated by the large multiplicity contained in an T(3S) event. Energy clusters tend to overlap each other and we have to optimize a photon algorithm which can select clean isolated clusters originating from photons without losing too much efficiency. We are still working on this problem. However, it is encouraging that our first preliminary inclusive spectrum from CUSB-II, figure 10, shows definite, distinct structures that, after background subtraction (figure 11), we can identify: the first three are the $\gamma_1$'s, and the second two are $\gamma_2$'s ending in T(2S) and T(1S). The width of these lines are also as
predicted from our MonteCarlo (MC) simulation, shown in figure 12. Finally, our preliminary analysis indicates complete agreement between CUSB-II and our published CUSB-I numbers for the photon energies and El transition rates.

In conclusion, our first run with CUSB-II fully vindicates our choice of BGO for the calorimetry. We need to complete the ancillary parts such as a vertex chamber for tracking and insertion of a layer of silicon strips at the shower maximum for $\theta$ shower centroid localising. Then we expect to commence our search for the singlet states (indicated by the short dashed lines in figure 7) of the $T(b\bar{b})$ system.

Figure 10. Inclusive $\gamma$ Spectrum. Figure 11. Inclusive $\gamma$ spectrum signal. Figure 12. MC BGO Resolution.

ACKNOWLEDGEMENTS

The author wishes to emphasize that CUSB-II is a group enterprise, albeit a small group by today's high energy physics standards, and thanks all collaboration members for the successful realization of CUSB-II. In particular, Drs. R. D. Schamberger and P. M. Tuts especially deserve credit for their leadership during its construction and installation, and Mr. T. Zhao for much work on this project. CUSB-II is supported by U. S. National Science Foundation Grants Phy-8310432 and Phy-8315800.
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†CUSB-II members include: P. Franzini, P. M. Tuts, S. Youssef and T. Zhao of Columbia University; J. Lee-Franzini, T. M. Kaarsberg, D. M. J. Lovelock, M. Narain, S. Sontz, R. D. Schamberger, J. Willins and C. Yanagisawa of SUNY at Stony Brook.

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Abstract

Preliminary results are presented from a comparative study of continuum production and fragmentation properties of the five charmed particles $D^0$, $D^\pm$, $D^{*\pm}$, $F^\pm$ and $A_c^\pm$, using data collected with the CLEO detector at the Cornell Electron Storage Ring. In addition an upper limit of $10.8\%(90\%CL)$ is set on the fraction of observed $D^{**}$ produced by the decay of an excited state $D^{**0} \rightarrow D^{**} \pi^-$. 
The processes by which quarks are physically manifested as hadrons are not understood in
detail. One interaction which is relatively favorable for the study of such processes is the high
energy $e^+e^-$ annihilation, in which a quark-antiquark pair is subsequently created. Since the
associated production of quarks to dress the primary quark has a rate inversely related to the
secondary quark mass, the course of a heavy primary quark can be traced by identification
of the heavy hadrons containing the quark. Although high energies are desirable to reduce
kinematic effects due to mass, the collection of data becomes difficult due to reduction in
cross sections and background from cascade decays of heavier quarks. At the Cornell Electron
Storage Ring (CESR) the study of charm is conducted near the threshold for production of
beauty.

The data reported here were taken at energies on, above and below the $T(4S)$ resonance. Since
B mesons produced on the $T(4S)$ resonance are a source of charmed particles with momenta
less than 2.5 GeV/c, only those with higher momenta were investigated in the analyses using
$T(4S)$ data. The most recent data, 78 pb$^{-1}$ on resonance and 35 pb$^{-1}$ at continuum ener­
gies below, were taken in 1985. Improvements to the CLEO detector$^1$ in mid-1984 included
the replacement of the three-layer inner proportional chamber by a ten-layer vertex detector
with a beryllium beam pipe and the addition of electronics to enable ionization measurements
in the CLEO main drift chamber. These, and improvements in software, have resulted in a
momentum resolution $\Delta p/p \approx 0.6\% p$ and expanded capability of particle identification, signif­
icantly improving the identification of several charmed particles, notably the $\Lambda_c^+$ in the decay
to $pK^-\pi^+$ (Fig. 1). Preliminary results on charm production from a study of the charmed
particles $D^0$, $D^+$, $D^{*+}$, $F^+$ and $\Lambda_c^+$ (inclusion of charge conjugate states is implied) in 113 pb$^{-1}$
are reported here. Also reported is a search for the decay of an excited charmed state $D^{**0}$
$\rightarrow D^{*+}\pi^-$, found by ARGUS at mass 2.42 GeV/c$^2$, in a total of 259 pb$^{-1}$ of data.

**Fragmentation**

In recent years charm production has been studied extensively, principally utilizing $D^{*+}$ mesons
in $e^+e^-$ annihilations.$^3$ These measurements have shown conclusively that the fragmentation
spectrum is hard, as expected,$^4$ and that within statistical errors it conforms to a shape pre­
dicted by C. Peterson et al.$^6$ Since the original Peterson publication there have proliferated a
number of theoretical predictions of the fragmentation function.$^6$ All are in reasonable agree­
ment with the available data and differ substantially from each other only in the statistically
poor lower momentum regions and in detailed behavior at the high momentum limits.

The comparison of experimental data with theory is hindered by several uncertainties. Fore­
most is the fact that many predictions are presented as a function of a variable which is not
experimentally measurable. The light cone variable, $x^+ \equiv \frac{E_h + p_{h\parallel}}{E_h + p_{h\parallel}}$ (where $E_q$, $p_q$, $E_h$ and $p_{h\parallel}$
are the primary quark energy and momentum, the hadron energy and the component of the
hadron momentum in the direction of the original quark, respectively), is favored theoretically
because of its Lorentz invariance but is well defined only if the original quark jet axis is known.
Its experimentally measurable analogue, $x_{exp}^+ \equiv \frac{E_h + p_{h\parallel}}{E_{beam} + p_{ma}}$, ($p_{ma} \equiv \sqrt{E_{beam}^2 - M_h^2}$) and the
other commonly used variables, $x \equiv p/p_{max}$ and $z \equiv E/E_{beam}$, are good approximations only at
Figure 1 (above). Distribution of pK-π+ invariant mass, for three-track combinations with total momentum exceeding 2.3 GeV/c. Requirements include positive identification of the K, consistency in identification of the p and π and cuts on individual track momenta and center-of-mass decay angles. The presence of Λc is evident in the fit to a smooth background plus Gaussian centered at 2.286 GeV/c².

Figure 2 (right). Measured CLEO fragmentation functions in the variable x≡p/p_{max} for (a)D⁰ and D⁺, (b)D^{∗+}, (c)F⁺ and (d)Λc⁺. Fits are shown to the fragmentation model of Peterson et al.,⁵ for purposes of illustration and comparison between different particles (see text). For (b) the weighted average of results for two decay modes is plotted, and the fit uses the value of c measured by CLEO in 1983.
high momenta or in the limit that the mass is small compared to the beam energy. This is not the case at CESR energies, so there is significant latitude in the choice of variables. Comparisons are further complicated by radiative QED and QCD effects which soften all spectra, as well as detector resolution which tends to smear details. For this report the fragmentation functions have been measured in the variable x and fitted to functions predicted by several parametrizable models (Peterson, Kartvelishvili, Collins and Bowler)$^5$ for the charmed particles $D^0$, $D^+$, $D^{*+}$, $F^+$ and $\Lambda_c^+$, in the decay channels shown below with respective branching fractions and numbers of reconstructed events used.

$$D^0 \rightarrow K^- \pi^+ (0.056\pm0.005)^7$$
$$D^+ \rightarrow K^- \pi^+\pi^+ (0.116\pm0.016)^7$$
$$D^{*+} \rightarrow D^0 \pi^+ (0.60\pm0.15)^8$$
$$D^0 \rightarrow K^- \pi^+$$
$$D^0 \rightarrow K^- \pi^+\pi^+\pi^- (0.118\pm0.014)^7$$
$$F^+ \rightarrow \phi \pi^+ (0.044)^9,$$
$$\phi \rightarrow K^+K^- (0.50\pm0.01)^9$$
$$\Lambda_c^+ \rightarrow pK^-\pi^+ (0.022\pm0.010)^9$$

It is found that the statistical uncertainties are too large to enable significant differentiation between models by direct fitting. It is, however, useful to compare the five particles, which have differing combinations of charge, spin, accompanying quark and baryon number, to determine the effects on fragmentation due to these properties. To define quantitatively the hardness of spectra one may, for example, examine fits to the parameter $\epsilon$ of the Peterson function, which has an inverse relationship with the hardness of the spectrum. The values of $\epsilon$ obtained from the fits for $D^0$, $D^+$, $F^+$ and $\Lambda_c$ are $0.23\pm0.04$, $0.20\pm0.04$, $0.15\pm0.05$ and $0.30\pm0.15$, respectively. Values previously measured by CLEO include $0.14\pm0.03$ ($D^*$) and $0.21\pm0.08^{10}$ ($\Lambda_c$). The pseudoscalar ($D$) spectra are in good agreement with each other and appear to be softened relative to the vector ($D^*$) spectrum, possibly due to significant contributions from $D^*$ decays. The difference between $D^*$ and $F$ is not significant, but the $\Lambda_c$ spectrum seems to be softer. The data with fits are shown in Fig. 2.

Total Production

The measured fragmentation functions are summed to yield total production cross sections, given below with the included x region indicated. For $F$ and $\Lambda_c$ the decay branching fractions are not well known, so the product of the branching fraction and cross section is calculated.

$$\sigma_{D^0} = 0.85 \pm 0.09 \text{ nb } \ x > 0.0$$
$$\sigma_{D^+} = 0.50 \pm 0.11 \text{ nb } \ x > 0.0$$
$$\sigma_{D^{*+}} = 0.32 \pm 0.01 \text{ nb } \ x > 0.5$$
$$B(F^+ \rightarrow \phi \pi^+)\sigma_F = 6.1 \pm 0.2 \text{ pb } \ x > 0.0$$
$$B(\Lambda_c^+ \rightarrow pK^-\pi^+)\sigma_{\Lambda_c^+} = 17. \pm 7. \text{ pb } \ x > 0.0$$
The measurement of the $D^{*+}$ in two modes differing only in the decay of the $D^0$ yields the ratio of branching fractions for the two $D^0$ decay modes,
\[
\frac{B(D^0 \rightarrow K^- \pi^+ \pi^+ \pi^-)}{B(D^0 \rightarrow K^- \pi^+)} = 2.2 \pm 0.2,
\]
consistent with the value $2.1 \pm 0.1$ measured by Mark III.\(^7\)

The production of $F$ is somewhat lower than that previously observed by CLEO in the 1983 data sample $(12. \pm 3. \text{ pb})$.\(^11\)

The above value for $A_c^+$ and the published CLEO value for the decay $A_c^+ \rightarrow A^0 \pi^+ \pi^+ \pi^-$,\(^10\) can be used to derive the ratio of branching fractions:
\[
\frac{B(A^+_c \rightarrow pK^- \pi^+)}{B(A^+_c \rightarrow A^0 \pi^+ \pi^-)} = 1.3 \pm 0.4
\]

**Search for $D^{*++}$ → $D^{*+} \pi^-$**

The ARGUS collaboration has reported the identification of an excited charmed state, which I will call $D^{*0}$, with mass 2.42 GeV/c\(^2\) decaying in the mode $D^{*0} \rightarrow D^{*+} \pi^-$.\(^2\) A search for this state has been conducted at CLEO using, in addition to the recent data, 59 pb\(^{-1}\) from the 1982 $T(4S)$ data and 67 pb\(^{-1}\) from the 1983 scan above the $T(4S)$ resonance.

The $D^{*+}$ were identified in the decay mode $D^{*+} \rightarrow D^0 \pi^+$ with the $D^0$ decaying to three different modes, $K^- \pi^+(1)$, $K^- \pi^+ \pi^+ \pi^-(2)$ and $K^- \pi^+ \pi^0(3)$. In decay (3) the $\pi^0$ is not observed but the $K^- \pi^+$ mass peaks in the region $1.5 < m_{K^-} < 1.7$ GeV/c\(^2\) due to polarization of the cascade decays $K^- \rho^+$ and $K^*^- \pi^+$ and the $D^{*+} - D^0$ mass difference is not sensitive to the missed particle.\(^12\) The difference in invariant mass $\Delta m$ between each $D^*$ candidate and its offspring $D^0$ was required to have $|\Delta m - 0.1454| < 0.002$ GeV/c\(^2\) for channels (1) and (2) and $\Delta m < 0.155$ GeV/c\(^2\) for (3). The $D^0$ candidate was then required to have mass within 25 to 65 MeV/c\(^2\) (depending on the data set) of the $D^0$ mass for channels (1) and (2) and 1.54 to 1.70 GeV/c\(^2\) for (3). To obtain the number of signal and background events the $D^0$ candidate mass distribution was fitted to a smooth background plus a Gaussian. The #peak/#background falling within the cuts was inferred to be 871/256, 1176/1057 and 836/1708, respectively. The invariant mass distributions are shown in Fig. 3.

Each $D^{*+}$ candidate was paired with each of the remaining tracks in the event and the invariant mass calculated assuming the additional track to be a pion. Comparison of the mass distributions for combinations with same and opposite sign charges and momenta $p_{D^{*+}} > 2.5$ GeV/c (Fig. 4) reveals no obvious relative structure in the 2.4 GeV/c\(^2\) mass region. The correct sign mass distributions were fit to a smooth background plus a Gaussian signal centered at 2.42 GeV/c\(^2\), yielding $23 \pm 16$, $123 \pm 30$ and $19 \pm 29$ $D^{*0}$ candidates, respectively (Fig. 5).

These values set an upper limit on the production ratio
\[
R^{**} = \frac{\sigma(D^{*0} \rightarrow D^{*+} \pi^-)(p > 2.5 \text{ GeV/c})}{\sigma(D^{*+})(p > 2.5 \text{ GeV/c})} < 10.8\% (90\% CL).
\]

This limit appears to be in disagreement with the corresponding value measured by ARGUS of $24^{+8}_{-6}\%$. Based on the reported rate and the numbers of $D^{*+}$ with $x_D > 0.45$ found in the two experiments it is estimated that $208^{+71}_{-48}$ and $132^{+61}_{-46}$ $D^{*0}$'s should be detected by CLEO in
Figure 3. Distributions of difference between $D^*$ and $D^0$ candidate masses for three different decay channels.

Figure 4. Difference between $D^*\pi$ and $D^*$ mass combinations, correct sign (solid line) and incorrect sign (dashed line) charges.

Figure 5. Fits of correct (opposite) sign distribution from Fig. 4 to smooth background plus Gaussian centered at 2.42 GeV/c².

Figure 6. Distributions for direct comparison between ARGUS (data points) and CLEO (solid histogram) for 1985 CLEO data, normalized to total number of events in plot.
the channels (1)+(2) and (3), respectively. It is also difficult to interpret as signal the $4\sigma$ effect found at CLEO in channel (2) due to the correlation of signals in the different channels with the known relative branching fractions of $D^0$ decays, verified for $D^{*+}$ in the previous section. Figure 6 shows a direct comparison, with CLEO data analyzed as reported by ARGUS. The difference between the two observations is puzzling in view of the nearly trebled size of the CLEO data sample relative to that on which the ARGUS result is based.

Summary

Preliminary results from a study of the production of five charmed particles in the most recent CLEO data set of 113 pb$^{-1}$ near the $T(4S)$ resonance have been reported here. It is anticipated that this study will yield useful information on fragmentation by comparison of charmed particles of different properties of charge, spin, companion quark and baryon number. It is not clear whether improvements to detector resolution and statistical significance will yield sufficient accuracy to enable comparison of different fragmentation models in individual particles.

A search in 240 pb$^{-1}$ of data collected near the $T(4S)$ resonance yields a signal for the decay $D^{**0} \rightarrow D^{*+}\pi^-$ which is consistent with zero. This is in disagreement with observations by ARGUS.

References

RESULTS ON HYPERON PRODUCTION FROM ARGUS

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ABSTRACT. Results on hyperon production are reported from data accumulated by the ARGUS experiment, operating at around 10 GeV centre-of-mass energies in the DORIS II e+e− storage ring at DESY. Signals for both the octet states Λ, Σ0 and Ξ− and the decuplet states Σ±, Ξ∗0(1530) and Ω− have been observed. Baryon rates from direct Υ(1S) decays are found to be enhanced by a factor of 3 over the continuum.
A large sample of $e^+e^-$ annihilation events at centre-of-mass energies near 10 GeV/c$^2$ has been accumulated using the ARGUS detector at the $e^+e^-$ storage ring DORIS II. Results are presented on the production of both spin 1/2 and 3/2 hyperons. For an updated report on our observation of $D^0 \rightarrow K^0\phi$, and for results on the reconstruction of B decays from $\Upsilon(4S)$ decays, the reader should consult reference 1.

ARGUS is a 4$\pi$ solenoidal spectrometer, with good charged and neutral particle detection. Particle identification is made possible both by measurement of mean $dE/dx$ loss in the central drift chamber, and by a system of time-of-flight counters. From this information, a likelihood ratio is constructed which describes the relative probability for each of the possible particle assignments. If the reconstruction procedure requires a specific particle type, then all tracks for which this probability exceeds 5% are considered acceptable candidates. A more detailed description of the experiment can be found in reference 2.

The long lifetime of the $\Lambda$ permits the use of vertex information as a means of reducing backgrounds. The procedure used for finding vertices is a necessarily iterative process, performed by making an unconstrained fit of a common point of origin to a given set of measured tracks and rejecting those tracks which contribute large $\chi^2$. A fit is considered successful if the distances of closest approach to the common origin are less than seven standard deviations for all tracks. The algorithm is first applied in the construction of a primary vertex, using all tracks, then to secondary vertices, using tracks which have been rejected from the primary vertex. A final step is to attempt the reconstruction of secondary vertices from combinations of remaining unassigned tracks with tracks already included in a vertex.

Figure 1 shows the invariant mass distribution for $p\pi^-$ combinations from reconstructed secondary vertices. The prominent $\Lambda$ peak contains $27330 \pm 195$ entries, as determined by a fit using the sum of two gaussians to account for the dependence of the resolution on momentum and decay point. The division of the number of observed $\Lambda$'s among data taken on the various $\Upsilon$ resonances and in the nearby continuum can be found in Table 1. Subtracting the continuum and vacuum polarization contributions to the signal on the $\Upsilon(1S)$, we find that the rate of $\Lambda$ production per event from $\Upsilon(1S)$ decays is $2.9 \pm 0.4$ times that for the continuum. The enhancement of the baryon rate on the $\Upsilon(1S)$ was first observed by DASP 4), and is in agreement with results reported by CLEO 5).

For purposes of reconstructing higher mass hyperons, all $\Lambda$ candidates within 3 sigma of the nominal $\Lambda$ mass, and with momenta above 200 MeV/c, were subject to a mass constraint fit and used to form mass combinations. Shown with the open histogram in figure 2a is the mass distribution of $\Lambda\pi^-\pi^-$ combinations with momenta greater than 400 MeV/c, but with no vertex requirement for the $\pi^-$, because of the relatively long average flight distance of the $\Xi^-$. 
A signal of 1885 ± 86 events is observed at the Ξ− mass, divided among the resonance and continuum data subsets as shown in Table 1. Once again, the production rate for baryons per event from direct Υ(1S) decays is found to be enhanced over the continuum, for the Ξ− by a factor of 2.9 ± 0.9.

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<td>Υ(1S)</td>
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<td>56 ± 17</td>
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<td>2020 ± 60</td>
<td>136 ± 28</td>
<td>8 ± 9</td>
</tr>
<tr>
<td>Total</td>
<td>143.8</td>
<td>27330 ± 200</td>
<td>1885 ± 86</td>
<td>279 ± 46</td>
</tr>
</tbody>
</table>

Table 1. Observed number of baryons on- and off-resonance.

In order to study Σ*± production, Λπ± combinations were used with the requirement that the π± fit to the primary event vertex, but with no momentum cut. Figure 2a (hatched histogram) and b show the resulting mass distributions. The Ξ− signal is still visible, along with a broad enhancement in both figures corresponding to the Σ− and Σ*+ respectively. The solid lines are the result of a fit with a polynomial background, a gaussian for the Ξ− and a relativistic Breit-Wigner of width 40(35) MeV/c² for the Σ−(Σ*+). The number of Σ− and Σ*+ events was obtained by integrating the fitted Breit-Wigner between 1.35 and 1.42 GeV/c², yielding 721 ± 93 and 793 ± 101 entries respectively. Correcting for efficiency and the missing tails of the Breit-Wigner, the following ratios for decuplet-to-octet baryon production are obtained:

Σ−/Λ = 0.047 ± 0.006⁻⁰.⁰⁵⁰⁺⁰.⁴⁶⁰
Σ*+/Λ = 0.052 ± 0.007⁻⁰.⁰⁰⁴⁺⁰.⁴⁰⁰

The sum of the two ratios is well within the limit reported by TASSO 6), but is considerably smaller than the value of 0.32 ± 0.14 ± 0.08 obtained by TPC 7).

A search for the Ξ*(1530), also a decuplet state, has been made in the Ξ−π+ mass spectrum (Figure 3). For this purpose, a Ξ− was defined as a Λπ− combination lying within
±20 MeV/c² of the nominal Ξ⁻ mass, and with momentum greater than 400 MeV/c. These candidates were then subject to a mass constraint fit, before being combined with the π⁺. The prominent Ξ⁺0(1530) peak corresponds to 279 ± 46 events. Table 1 once more shows the division of the signal among the subsets of the data on- and off-resonance. In order to compute the Υ(1S) direct rate, the combined Υ(4S) and continuum signals were used for the continuum subtraction. By this means, the inclusive rate for Ξ⁺⁻(1530) production from Υ(1S) decays was found to be a factor of 3.3 ± 1.1 times the continuum rate. Averaging over all data, the ratio of Ξ⁺0(1530) to Ξ⁻ production is 0.26 ± 0.04±0.04. This is well in excess of the limit of 16% at the 90% CL reported by CLEO 8).

Evidence for Ω⁻ production can be seen in the ΛK⁻ mass distribution shown in Figure 4. The observed signal of 53 ± 16 events at the Ω⁻ mass corresponds to a Ω⁻/Λ ratio of (5.4 ± 1.8 ± 1.2) × 10⁻³. Using the present average for the Λ production rate 9), this is equivalent to an Ω⁻ rate of (1.2 ± 0.5) × 10⁻³ per event. The signal reported by TPC 7) at centre-of-mass energies of 29 GeV is equivalent to (2.7±1.8 ± 0.8) × 10⁻² per event, or more than an order of magnitude larger.

Finally, in the invariant ΛγC mass distribution (Figure 5), where γC is a converted photon reconstructed from e⁺e⁻ pairs forming secondary vertices, 53 ± 10 events are observed at the Σ⁰ mass. The product of conversion probability times reconstruction efficiency is only a few percent, but the excellent resolution for the photon energy results in a detectable signal with a RMS width of only 2.7 ± 0.5 MeV/c². The acceptance corrected ratio of Σ⁰ to Λ production is determined to be 0.33±0.11.

One way to interpret these results is in terms of models where baryons are created from combinations of quarks and di-quarks. The observed pattern of octet and decuplet production reflects the different probabilities for producing singly- or doubly-strange di-quarks in a spin 0 or spin 1 state 6). A fit to the ARGUS results, with the additional constraint provided by the present average Λ and Λ/Ξ⁻ rates 9), leads to the following values for these relative probabilities:

P(s)/P(u) = 0.310 ± 0.014
P(us)₀/P(ud)₀ = 0.058 ± 0.005
P(ud)₁/P(ud)₀ = 0.081 ± 0.013
P(us)₁/P(ud)₀ = 0.0044 ± 0.0013
P(ss)₁/P(ud)₀ = 0.0010 ± 0.0003

where P(qq)ᵢ is the probability for producing a spin j di-quark. Further insight into the production mechanism will be possible through study of baryon-baryon correlations.


3) References to a specific charged state should be interpreted as implying the charge conjugate state also.

4) C. W. Darden et al. (DASP II collaboration) Phys. Lett. 80B (1979) 419.


7) H. Aihara et al. (TPC collaboration), contributed paper to the European Physical Society Meeting on High Energy Physics, Bari (1985).


Figure 1

Figure 2a

Figure 2b

Figure 3

Figure 4

Figure 5
ABSTRACT: Studies on the construction of a high resolution muon detector to operate in high intensity, multi-TeV collider have been performed. Prime objective of this detector is to discover new, heavy particles by precise measurements of the muon pairs they decay into. A 1 TeV particle decaying into μ-pairs will be discovered with a 1.4% mass resolution.
1. INTRODUCTION

In view of various proposals to construct high energy hadron colliders in the next decade, the 20 TeV on 20 TeV Super Conducting Super Collider (SSC) in the United States, to 50 TeV on 50 TeV ELOISATRON in Italy and the Large Hadron Collider (LHC) at CERN, a design study of a new detector has been carried out.

This detector (Fig. 1 and 2), named L3+1, will consist of two calorimeter systems followed by high resolution muon chambers, which are installed inside a large magnet. The L3+1 Detector will completely cover the azimuthal angle and a polar angle region between 40° and 140°. With this detector arrangement we will operate the L3+1 Detector at highest luminosities foreseen, $L=10^{33}$ cm$^{-2}$ s$^{-1}$, where order of $10^{10}$ particles/s will be produced. Furthermore the construction concept of the L3+1 Magnet does allow easy adaptation of the inner detector system should the physics objectives change.

Figure 1: The L3+1 Detector (front view)
2. THE INNER CALORIMETER

Surrounding the beam pipe we will install a first calorimeter system made out of 20 layers of tungsten plates (3.5 mm thick) sandwiched with wire chambers operating in the proportional mode. This Inner Calorimeter has 20 radiation lengths equivalent to 0.7 interaction lengths and will be used to measure hadronic and electromagnetic showers. The first tungsten plates will be installed at a radial distance of 50 mm with respect to the beam line, the last at 220 mm.

3. THE OUTER CALORIMETER

Immediately after the first calorimeter system we will install the
Outer Calorimeter which will have a coarse sampling. After the first four layers of tungsten plates (7.0 mm thick) the Outer Calorimeter will consist of 128 layers of 7.0 mm thick uranium plates. These absorber plates again will all be interleaved with 5.0 mm thick proportional chambers. The Outer Calorimeter thus will have 264 radiation lengths equivalent to 8.8 interaction lengths. We have chosen this arrangement of absorber material in order to shield the uranium plates from the high level of radiation produced by the interaction beams.

We have concluded from our experience in building a similar calorimeter system for the L3 Experiment to achieve energy resolutions of 40%/√E for electromagnetic and 50%/√E for hadronic showers (E in GeV). An energy resolution as good as this is essential for a high resolution spectrometer as L3+1, since muons have lost a non negligible amount of energy when passing through the calorimeter system. High energy muons loose most of their energy via Bremsstrahlung and pair production. Including also effects of ionization and nuclear interactions we have calculated the energy loss in the L3+1 calorimeter system for muons of energies between 0.5 TeV and 5.0 TeV. We also quote the measurement error and list the results in Table 1.

<table>
<thead>
<tr>
<th>Muon Energy (TeV)</th>
<th>Energy loss (GeV)</th>
<th>Error (GeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.5</td>
<td>30</td>
<td>2</td>
</tr>
<tr>
<td>1.0</td>
<td>58</td>
<td>3</td>
</tr>
<tr>
<td>2.0</td>
<td>115</td>
<td>4</td>
</tr>
<tr>
<td>5.0</td>
<td>236</td>
<td>7</td>
</tr>
</tbody>
</table>

4. MUON CHAMBERS

Following the principles of construction of muon chambers as developed for the L3 Detector we intend to install three packages of drift chambers at radial distances of 263 cm, 538 cm and 813 cm. The packages will consist of 32, 64 and 32 planes of wires. As indicated in Figure 3 we will track muon trajectories over a distance of 5.50 meters. To achieve a
high momentum resolution it is essential to measure the sagitta $S$ of the muon traversing in the magnetic field with great precision. For example, a muon of momentum $p_\mu=500$ GeV/c would have a sagitta of $S=1720$ µm when traversing the L3+1 muon chamber system which is inside a 7.5 kG magnetic field.

Using the techniques developed and optimized in constructing the precision drift chambers for the L3 Experiment we will obtain a single wire resolution of $\sigma=150$ µm. Since muon tracks are measured using 32 or 64 wires in a chamber package, we determine the average track position with a precision of $\varepsilon_1=\varepsilon_3=27$ µm for the inner and outer package and $\varepsilon_2=19$ µm for the middle one. Calculations of the measurement error on the sagitta $S$, including the omission of 6% of wires effected by $\delta$ rays and an alignment of the chamber system to better than 20 µm, lead to $\Delta S=34$ µm. The error on the sagitta has a direct influence on the momentum resolution. Figure 4 displays the expected momentum resolution of muons measured with the L3+1 Detector as a function of their momentum.

Using the relation

$$\frac{\Delta M}{M} = \frac{\Delta p}{\sqrt{2}p} = \frac{\Delta S}{\sqrt{2}S}$$

for a heavy particle at rest decaying into a muon pair, we obtain a 2% momentum resolution for muons of $p_\mu=500$ GeV/c, resulting in a mass resolution of 1.4% for a 1 TeV particle.
5. **MAGNET**

The construction of the magnet will be similar to that of L3. We will weld prefabricated aluminum plates together to form a 300-turn coil. The inner radius of the coil is 8.60 m, the outer 9.60 m. The coil will weigh 2700 tons and have a length of 20.50 m. It will be operated with 26.5 kA. Surrounding the coil we will install 27500 tons of iron with a thickness of 1.50 m. The pole pieces of 1.95 meters thickness weigh 12000 tons and will be removable in order to allow access to the inner detector components. The magnetic field strength of the L3+1 Detector is 0.75 T.

6. **CONCLUSIONS**

The L3+1 Detector will enable us to measure muon pair masses of $M_{\mu\mu} = 1$ TeV with a 1.4% precision. With a detector of this simple concept, but high accuracy, we will be able to explore the existence of new particles and perform precision muon measurements up to energies of several TeV.

**ACKNOWLEDGEMENT**

I would like to thank my colleagues U. Becker, M. Chen, M. Fukushima, H. Hofer, P. Lecomte and S.C.C. Ting for many interesting and valuable discussions as well as for their contributions to this study. Special thanks to M. Harris and F. Wittgenstein for their advice on technical matters.

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The large N method, where N is the number of spatial dimensions, is a powerful, new technique for analytically determining the eigenstates of the Schrödinger equation, even for potentials which are in no sense "small" and hence not treatable by standard perturbation theory. A somewhat modified, physically motivated approach, which yields "shifted large N expansions", incorporates exactly known analytic results into large N expansions, greatly enhancing their accuracy, simplicity and range of applicability. The rate of convergence can be further improved by using supersymmetry transformations. A systematic large N approach to potential scattering problems has also been recently developed.
The main reason for widespread interest in large $N$ expansions is because they offer the possibility of analytically handling interactions which need not be small. The symbol $N$ has different meanings for different physical problems. For example, in the context of quantum chromodynamics, $N$ is the number of colors, and one examines the consequences of a SU($N$) gauge group with $N$ large, although to get physically interesting results one eventually wants $N = 3$ colors. In the large $N$ limit, there is a substantial simplification in calculations, since Feynman diagrams with loops and complicated topologies do not contribute in leading order.

In this article, we will confine our attention to non-relativistic quantum mechanics. In this context, the symbol $N$ stands for the number of spatial dimensions. There are very few potentials whose eigenstates can be found exactly, and in most cases the only Hamiltonians one can treat are those of the form $H = H_0 + \epsilon H_1$ where $H_0$ is exactly solvable and $\epsilon$ is a small expansion parameter. The eigenfunctions and eigenvalues can then be expressed as a perturbation series in $\epsilon$. Potentials without a small expansion parameter can be handled by the WKB approximation (which works best for excited states) or by cumbersome variational techniques (which work best for low-lying states). Recent work has demonstrated that large $N$ expansions provide a powerful, systematic and direct technique of solving the Schrödinger equation, even for potentials which are in no sense "small". Basically, the method consists of working in $N$ spatial dimensions, assuming $N$ to be large, and taking $1/N$ as an "artificially created" expansion parameter for doing standard perturbation theory. At the end of the calculation one sets $N = 3$ to get results for three dimensional problems. Since $1/3$ is not a very small expansion parameter, one needs to calculate many orders in perturbation theory for obtaining good accuracy. However, the accuracy and convergence of such $1/N$ expansions can be greatly enhanced by incorporating exactly known analytic results. For example, for effectively one dimensional potentials, a remarkable improvement in the accuracy of energy eigenvalues results from using the quantity $1/\tilde{k}$ as an expansion parameter, where $\tilde{k} = N + 2 \ell - a$, $\ell$ = orbital angular momentum quantum number and 'a' is a shift whose form can be physically motivated from known analytic results for the Coulomb and harmonic oscillator potentials. At the present time, such "shifted large $N$ expansions" are probably the simplest known analytic expressions which accurately describe the eigenstates of arbitrary spherically symmetric potentials.
For an arbitrary spherically symmetric potential \( V(r) \) in \( N \) dimensions, one has the following radial Schrödinger equation:

\[
\left[- \frac{\hbar^2}{2m} \frac{d^2}{dr^2} + V_{\text{eff}}(r) \right] \phi(r) = E \phi(r), \tag{1}
\]

\[
V_{\text{eff}}(r) = V(r) + \frac{(N+2\ell-1)(N+2\ell-3)}{8mr^2}. \tag{2}
\]

It is important to note that \( N \) and \( \ell \) always appear together in the form \( k = N + 2\ell \). This means that the eigenstates, which could have depended on the three quantities \( N, \ell, n \), in fact only depend on \( k \) and \( n \), where \( n = 0, 1, 2, \ldots \) is the radial quantum number corresponding to eq. (1). As mentioned before, we will obtain a systematic expansion of eigenstates in \( 1/k \), where \( k = k - a \). Of course, when \( N \) is very large, \( k \approx k \). However, for \( N = 3 \) dimensions, a properly chosen shift 'a' produces great improvements in accuracy and simplicity.

For concreteness, let us focus on the class of power law potentials \( V(r) = Ar^v \) (\( v > -2 \)). The characteristic length and energy appearing in this problem are

\[
r_c = \left( \frac{\hbar^2}{2mA} \right)^{1/(v+2)}, \quad E_c = Ar^v_c = \frac{\hbar^2}{2m} \left( \frac{2mA}{\hbar^2} \right)^{2/(v+2)}. \tag{3}
\]

It is convenient to deal with the dimensionless variables \( \xi = r/r_c \) and \( \lambda = E/E_c \) obtained by scaling out \( r_c \) and \( E_c \). In order to get a useful large \( k \) limit, we want an effective potential whose shape does not vary with \( k \) at large \( k \). This is customarily accomplished by re-scaling the radial variable. \(^2\text{3}\) Defining \( n = \xi k - 2/(v+2) \) and \( \lambda = \lambda k - 2v/(v+2) \), the radial Schrödinger equation now reads

\[
\left[- \frac{1}{\Lambda^2} \frac{d^2}{dn^2} + \frac{1}{4n^2} (1 - \frac{1-a}{2})(1 - \frac{3-a}{2}) + \eta^v \right] \phi = \bar{\lambda} \phi. \tag{4}
\]

The effective potential at large \( \lambda \) is \( V_{\text{eff}}(n) = 1/(4n^2) + \eta^v \), which has a minimum at \( n_o = (2v)^{-1/(v+2)} \). Clearly, the eigenvalue \( \bar{\lambda} \) has a leading contribution \( V_{\text{eff}}(n_o) \), corresponding to the classical result of a particle at the bottom of a potential well i.e., \( \bar{\lambda} = V_{\text{eff}}(n_o) + O(1/\lambda) \). The next order clearly consists of making a parabolic, harmonic oscillator approximation to the effective potential with an angular frequency \( \omega = \sqrt{v + 2} \). This yields harmonic oscillator energies \((n + \frac{1}{2})\hbar\omega\), where \( n = 0, 1, \ldots \) is the
radial quantum number. Higher order corrections in $1/k$ can be handled by standard perturbation theory.

Choosing the shift to be $a = 2 - (2n + 1)\frac{\sqrt{\nu} + 2}{2}$, one gets the following analytic expression for the energy levels.

$$E = \frac{\nu + 2}{\nu^2} \left( \frac{\nu + 2}{\nu} \right) \frac{2}{m} \left[ \frac{\nu^2}{k} \right] - \frac{(\nu + 1)(\nu - 2)}{144k} \{ (1 + 6n + 6n^2) + \mathcal{O}(\frac{1}{k}) \}. \ (5)$$

In three dimensions, for the ground state ($n = \ell = 0$), the energy obtained from eq. (5) as a function of $\nu$ has better than 99% accuracy for potentials in the wide range $-1.5 \leq \nu \leq 8.31$ The results are of course even better for excited states with $\ell > 0$, since $1/k$ is smaller.

Shifted large $N$ expansions can be used for arbitrary spherically symmetric potentials $V(r)$. As might be expected, the expansion contains the potential and its derivatives evaluated at the minimum of the large $N$ effective potential. The general expression can be found in Ref. (3). The accuracy of shifted expansions has been tested for many smoothly changing potentials, e.g. charmonium, Morse, anharmonic oscillator, logarithmic, Yukawa, etc. In general, the accuracy for all eigenstates is extremely good. Of course, the approach will not work so well for rapidly varying potentials which cannot be adequately described by a few terms in the Taylor series expansion of the effective potential about its minimum.

Given any one dimensional potential, supersymmetric quantum mechanics provides a simple recipe for generating a partner potential with the same energy eigenvalues (except for the ground state). For a potential $V^0(x)$, the partner potential is $V^1(x) = V^0(x) - \frac{d}{dx} \left( \psi_0'/\psi_0 \right)$, where $\psi_0(x)$ is the zero energy ground state wave function. For example, the partner potential for an infinite square well is $V(x) = \csc^2 x$. Often, for many physical problems, it is profitable to deal with partner potentials. We have been able to show that the use of supersymmetric partner potentials can be exploited to further improve both the accuracy and simplicity of large $N$ expansions. Typically, we find that just the leading term in a shifted large $N$ expansion using partner potentials yields all the energy levels correctly to three significant digits for essentially any three-dimensional spherically symmetric potential of physical interest. Basically, the improvement occurs because the supersymmetric partner of a given potential in $N$ dimensions can be effectively treated as being in $N+2$ spatial dimensions, and this property decreases the expansion parameter.
The considerable success of shifted large N expansions for handling bound states, motivates one to ask whether similar techniques could be applied to the scattering region. A first attempt used the fact that the exactly solvable delta-shell potential could be chosen as the large N unperturbed problem.\cite{8} Recently, we have been able to extend these ideas to construct the first systematic large N expansion for scattering phase shifts. Scattering lengths, computed by the large N approach, are found to be in excellent agreement with numerical results for a wide variety of potentials.\cite{9} Further work in this direction is still in progress.

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References


ON THE NATURE OF THE EMC EFFECT

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A B S T R A C T

It is shown that none of the popular models suggested to explain the EMC effect seem satisfactory. A new point of view on the effect as a simple relativistic phenomenon is presented and is shown to be favoured by NA4 data.
From the time of the discovery of the EMC effect\(^1\), a number of explanations were proposed\(^2\). However, none of them seems satisfactory. All models incorporate in some form a suppression of the hard part \((x>0.4)\) of the valence quark distribution but differ by the assumed source of compensation of this suppression.

The following compensation mechanisms have been considered:

(i) by soft part \((x<0.3)\) of valence quark distribution\(^3,4\);

(ii) by pions in nucleus\(^5\);

(iii) both by the quark and the gluon sea due to increase of the confinement radius (rescaling models\(^6,7\)).

We shall demonstrate next that none of them seems to give a good description of the experimental data.

Since the QCD evolution equations do not depend on the type of the target one can show\(^8,9\) that up to higher twist corrections

\[
T_{A}(x, Q^2) = \int_{x}^{A} d\beta T_{A}^{NS}(\beta) V_{A}(\frac{x}{\beta}, Q^2) \equiv T_{A}^{NS}(x, Q^2)
\]

\[
(\Sigma_{A} + C_{\pm} \otimes C_{A}) = T_{A}^{\pm}(\Sigma_{N} + C_{\pm} \otimes C_{N})
\]

where \(V, E, G\) are the valence (non-singlet) quark, singlet quark, and gluon distribution functions, \(C_{\pm}\) are coefficients depending on \(x\) and \(\alpha_s(Q^2)\) which can be found from the perturbative QCD and \(T_{A}^{NS}, T_{A}^{+}\) are some functions independent of \(Q^2\) characterizing the nucleus \(A\). Due to the baryon number and energy-momentum sum rule \(T_{A}^{NS}, T_{A}^{+}\) satisfy the following conditions (the index \(A\) will be omitted)

\[
\int_{x}^{A} T_{A}^{NS}(\beta) d\beta = 1, \quad \int_{x}^{A} T_{A}^{+}(\beta) d\beta = \frac{M}{A m} \approx 1
\]

where \(M\) and \(m\) are the mass of nucleus and free nucleon.

For large \(Q^2\) and \(x > 0.3\) where the sea quark distributions, \(O = \Sigma - V\), in nucleus and free nucleon are negligible, one can write down

\[
\Sigma_{A}(x, Q^2) = \int_{x}^{A} T_{A}^{NS}(\beta) \Sigma_{N}(x, Q^2) d\beta
\]

Developing \(\Sigma_{N}(x/p)\) into power series in \(\delta = 1-\beta\), it is easy to obtain

\[
R(x) = \frac{\Sigma_{A}(x)}{\Sigma_{N}(x)} \approx \delta_x - \frac{1}{2} \delta_x^2 \cdot \kappa_2 + \frac{1}{2} \delta_x^3 \cdot \kappa_3 \cdot (\kappa-1) \cdot 2^2 + \ldots
\]

where \(\kappa^k = \int_{x}^{A} \delta_x^k T_{A}^{NS}(\beta) d\beta, \Sigma_{N}(x) \sim (1-x)^k, k = 3, \xi = x/(1-x)\) and the third term disappears at \(x = 0.5\) \((\xi = 1)\).

Now, it is known from experiment that both \(V_{N}(x)\) and \(V_{A}(x)\) are positive. From (1) then follows that the function \(T_{A}^{NS}\) is also positive and one can prove that the condition
is necessary to have \( R(0.5) < 1 \) (large \( x \) depletion), i.e., it is impossible to compensate the depletion by increase of the soft part of the valence quark distribution assumed in Ref. 3). This also implies the impossibility of a screening effect for the valence quarks assumed in Ref. 4). The required depletion for \( R(\text{Fe/Fe}) \) is gained when \( \delta = 0.04 - 0.045 \).

2° Let us examine next the possibility of the compensation of the missing momentum \( \langle x_V^A \rangle - \langle x_V^N \rangle \) by quark and gluon seas. Consider the ratio

\[
R_I = \frac{\int_0^A F_{2A}(x) dx}{\int_0^A F_{2D}(x) dx} \approx \frac{\langle x_q^A \rangle}{\langle x_q^N \rangle}.
\]

We assume for simplicity that \( \Sigma(x) = F_2(x) \) and \( F_{2D}(x) = F_{2N}(x) \).

Estimation of \( R_I \) using the EMC data \(^{10}\) and the natural continuation \( F_2(x < 0.05) = F_2(x = 0.05) \) gives

\[
R_I^{\text{EMC}} - 1 = (7.1 \pm 1.0 \pm 3.0) \cdot 10^{-2}
\]

Such a big positive number is difficult to obtain in any pion model. It is easy to show that \( R_I - 1 = \delta (\langle x_q^A \rangle / \langle x_q^N \rangle - 1) \) and with \( \delta = 0.045 \) the number (7) corresponds to \( \langle x_q^A \rangle / \langle x_q^N \rangle > 1.89 \), instead of \( \approx 1 \). The number (7) contradicts also any rescaling model because the transition from free nucleon to nucleus is equivalent to increasing of \( Q^2 \). So, \( \langle x_q^A \rangle \) has to increase and \( \langle x_q^N \rangle \) has to decrease, just opposite to (7).

One can possibly argue that the EMC group somehow overestimates the small \( x \) region and that the other experiments \(^{11}\) and first of all NA4 \(^{12}\) do not show such an excess for the ratio \( N_2/D_2 \) with a small systematic error. Indeed, a calculation of \( R_I \) for the nitrogen nuclei using the NA4 data gives

\[
R_I^{\text{NA4}} - 1 = (0.7 \pm 1.7 \pm 1.0) \cdot 10^{-2}
\]

which is compatible with zero. However, much smaller excess of the data over \( R = 1 \) in the small \( x \) region leads even to a greater difficulty for pion and rescaling models because it implies a much smaller number of pions in nucleus or a much smaller difference in confinement radii of free and bound nucleons which results in a much smaller depletion for \( x > 0.3 \) in strong disagreement with the data [cf. e.g., Ref. 13].

So we see that none of these compensation schemes seems satisfactory both for EMC or NA4 data. It should also be added that the multiquark fluctuations \(^{14}\) in nucleus, the existence of which is supported by many unusual phenomena in hadron-nucleus scattering, e.g., cumulative particle production \(^{15}\), can of course be
responsible for the EMC effect. However, the structure function of the fluctuations is not constrained enough to consider such an explanation as totally satisfactory.

3° The main reason for such an uneasy situation with the explanation of the EMC effect lies, in our opinion, in the fact that it is usually assumed that the nucleus is a bound state of A nucleons. This is definitely a wrong assumption in a relativistic theory because the number of particles is not conserved in an interacting system like the nucleus. Only the baryon number, i.e., the number of nucleons minus the number of antinucleons is conserved. So, first what we need is a true relativistic description. For this purpose, the Bethe-Salpeter equation (Fig. 1) seems to us to be the most suitable one.

In the impulse approximation it leads to the equation (1) for $V_A$ and to the following equations for the singlet part

$$\left\{ \begin{array}{l}
\Sigma_A(x, Q^2) \\
G_A(x, Q^2)
\end{array} \right\} = \int_0^A T^S(p) \left\{ \begin{array}{l}
\Sigma_N(x, Q^2) \\
G_N(x, Q^2)
\end{array} \right\}$$

(9)

$$G_A(x, Q^2) = \int_0^A T^S(p) G_N(x, Q^2)$$

(10)

where

$$T^{NS} = N(p) - \bar{N}(p), \quad T^S = N(p) + \bar{N}(p)$$

(11)

$$(\bar{N})N$$ being the densities of (anti-) nucleon in the nucleus with a given momentum fraction $\beta$. It is equivalent to Eqs. (2) obtained from QCD with an additional ansatz

$$T^+ = T^- \equiv T^S$$

(12)

The latter could be a consequence of the absence of free quarks and gluons in the nucleus due to colour confinement and quark-hadron duality. Therefore, we believe that Eqs. (9) to (11) have a more general validity.

Using these equations it is easy to see that $\bar{N}(p) = 0$ immediately results in

$$\bar{N} = \int(1-\beta)T^{NS}(\beta)d\beta = 0$$

and the absence of a large $x$ depletion in $R$ as it was shown in Section 1. If, however,

$$\int_0^A \bar{N}(p)d\rho = \Delta \neq 0, \quad \int_0^A N(p)d\rho = 1 + \Delta > 1$$

(13)

one can obtain from (11) and (3) the excess over $R = 1$ at $x = 0$

$$R(x = 0) = \int_0^A T^S(p)d\rho = 1 + 2\Delta > 1$$

(14)

Using the development of (9) in power series of $(1-\beta)$ similar to Eq. (4), one obtains

$$R(x) \approx 1 + 2\Delta (1 - K x^2) + \cdots$$

(15)
which guarantees also a depletion at $\xi > 1/k$ $(x>1/(k+1))$. So, these two features of the EMC effect are a simple consequence of $\Delta \neq 0$ and are closely linked with each other. For instance, if $k = 3$ and $R(0.5) = -0.1$, then the excess will be $R(0)-1 = 0.05$ independent of the details of $T^S(\beta)$. This is one of the characteristic predictions of our approach.

Notice also that in the approximation (15) the position of the intercept $R = 1$, $x = 1/(k+1)$ does not depend on the type of the nucleus, which is also one of the distinctive features of the EMC effects.

We shall try to connect next the number of antinucleons $\Delta$ with the average nucleon energy level in the nucleus. For this purpose, let us assume that the inter-nuclear forces are $C$-independent, i.e., that the average energies per particle for nucleons and antinucleons are equal. Defining now the average nucleon energy level as an energy of evacuation of one nucleon, $\int_0^A \rho N(\beta) d\beta = 1-\overline{\epsilon}/m$ one can obtain

$$\overline{E} = \frac{m\Delta}{1+2\Delta} \approx m\Delta$$

i.e., the depletion $R(0.5)-1 = 0.1$ for Fe corresponds to $\overline{\epsilon} = 23$ MeV which agrees with standard nuclear physics.

In order to obtain a more detailed behaviour of $R(\alpha)$ we have assumed for $T^S$ a simplest form of "shifted" Fermi step function

$$T^S(\beta) = \frac{4}{3\beta_0} \left( \frac{m}{K_F} \right)^3 \left\{ \begin{array}{ll} (\frac{K_F}{m})^2 - (\beta_0 - \beta)^2 & \text{for } |\beta_0 - \beta| < \frac{K_F}{m} \\ 0 & \text{otherwise} \end{array} \right.$$  

where $\beta_0 = (1+2\Delta)^{-1}$ [or $\overline{\epsilon}$] and $K_F/m$ are considered as parameters to be determined from the data. Figure 2a demonstrates the degree of agreement with the NA4 data for the ratio $R(N_2/D_2)$ supplemented by the SLAC data for $R(C^{12}/D_2)$ with $x > 0.65$. Figure 2b demonstrates the same for the ratio $R(Fe/D)$. The behaviour in the region $x < 0.2$ can be considered as the prediction. Figure 3 presents the A-dependence of the parameters for the whole set of nuclei measured at SLAC. (Due to a possible correction of SLAC data for small $x$ due to the A dependence of $\sigma_L/\sigma_T$ ratio only the points with $x > 0.3$ were included into the fit.) One can see a slight A-dependence of both parameters and an indication on a somewhat higher value of $\overline{\epsilon}$ for even-even nuclei. Notice also the good coincidence of the parameters for the iron data of SLAC and NA4 and a higher value of $\overline{\epsilon}$ for nitrogen. The latter could be a consequence of a small, $\approx 1-2\%$, systematic shift of the data. The results of the calculation of Fe/D ratio for the sea quarks and the gluons, and the corresponding data of CDHS and of EMC are presented in Fig. 4. A disagreement with EMC data for the gluons does not bother us very much.
because it seems difficult to agree such a big excess of gluons in nucleus with a
decrease of the gluon average momentum resulting from (7).

In conclusion, it is a pleasure for me to thank T. Ericson, H. Pirner,
P. and G. Todorov, L. Van Hove and especially J. Noble for fruitful and
stimulating discussions and the TH Division of CERN for kind hospitality.

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FIGURE CAPTIONS

Fig. 1 The Bethe-Salpeter equation for nuclear structure function.

Fig. 2 Fit of NA4 (•) - SLAC (A) data for (a) the ratio R(N2/D2) and
(b) R(Fe/D) by expressions (9) and (18).

Fig. 3 The A-dependence of parameters E/m (•) and K_F/m (A) obtained from SLAC
and NA4 (A = 14,56) data.

Fig. 4 The comparison of the ratio Fe/D for (a) sea quarks and (b) gluons with
the data of CDHS and EMC groups.
We study how the properties of toponium states depend on the scale parameter $\Lambda_{\overline{MS}}$. We show that hyperfine and fine splittings are sensitive to the value of $\Lambda_{\overline{MS}}$, and hence the allowed range of $\Lambda_{\overline{MS}}$ may be restricted to be around 200 MeV from the data of already known quarkonia, $\psi$ and $b\bar{b}$. 
What I would like to talk about today is how the properties of toponium states depend on the QCD scale parameter \( \Lambda_{\overline{\text{MS}}} \).

As you already know, heavy quark systems have a very clean spectroscopy, and in the very near future, the \( e^+e^- \) colliding machine such as TRISTAN, SLC and LEP will find toponium states. So, it is worthwhile to study properties of \( \bar{t}t \) system using an interquark potential obtained from \( c\bar{c} \) and \( b\bar{b} \) systems.

**\( \bar{Q}Q \) potential**

Let me start from the \( \bar{Q}Q \) potentials which are consistent with QCD and reproduce \( c\bar{c} \) and \( b\bar{b} \) properties using the Schrödinger equation. As you all know, Buchmüller and Tye\(^2\) has proposed an elegant potential which accommodates the desired properties for the \( \bar{Q}Q \) force. Their potential, however, implies a fairly large \( \Lambda_{\overline{\text{MS}}} \) of about 500 MeV, which is related to the Regge slope parameter \( (\alpha' \sim 1 \text{ GeV}^{-2}) \). It is not possible to incorporate a small \( \Lambda_{\overline{\text{MS}}} \) (down to 300 MeV) in this scheme to reproduce correct \( c\bar{c} \) and \( b\bar{b} \) spectra.

To overcome such problems, let us use the following potential (in the coordinate space);

\[
V(r) = V_{\text{AF}}(r) + ar
\]

where

\[
V_{\text{AF}}(r) = -\frac{16\pi}{25} \frac{1}{r f(r)} \left( 1 + \frac{2\gamma_E + \frac{53}{75}}{f(r)} - \frac{462 \ln f(r)}{625 f(r)} \right)
\]

with \( f(r) = \ln \left[ 1/(\Lambda_{\overline{\text{MS}}} r)^2 + b \right] \).

The asymptotic freedom potential \( V_{\text{AF}}(r) \) above satisfies the two-loop perturbative calculations for \( r \to 0 \), and becomes zero for \( r \to \infty \). We can easily prove that such an additional constant does not change the short-range asymptotic behavior up to the order of \( 1/\ln^2 (1/\Lambda_{\overline{\text{MS}}} r^2) \).

Let us now find parameters \( \Lambda_{\overline{\text{MS}}} \), \( a \) and \( b \). Suppose we fix our parameters by minimizing \( \chi^2 = \sum (m_{\text{theory}} - m_{\text{exp}})^2 \), with \( i = J/\psi, \psi(2S), \psi(4S), T(1S), T(2S), T(3S) \).

We can find an excellent fit by taking the following parameter

\[
\Lambda_{\overline{\text{MS}}} = 300 \text{ MeV}, \ a = 0.1414 \text{ GeV}^2, \ b = 1.9.
\]

Please note that \( \Lambda_{\overline{\text{MS}}} = 300 \text{ MeV} \) is more reasonable than \( \Lambda_{\overline{\text{MS}}} = 500 \text{ MeV} \) used in ref.\(^2\). In order to relax the possible range of \( \Lambda_{\overline{\text{MS}}} \), keeping the confining potential, we can add a term which disturbs neither the short-range asymptotic behavior nor the
confining part of the potential,
\[ V(r) = V_{AF}(r) + d r e^{-\alpha r} + \alpha r. \]

Then, we can change \( \Lambda_{\overline{\text{MS}}} \) from 100 MeV to 500 MeV, keeping the good fit to the \( c\bar{c} \) and \( b\bar{b} \) spectra\(^1\).

**Toponium spectroscopy and \( \Lambda_{\overline{\text{MS}}} \)**

We are in a position to predict the \( t\bar{t} \) spectra, assuming the flavor independence of the quarkonium potential. There are some theoretical ambiguities, however.

The reason is the followings:

(i) The short-range behavior \( (r \leq 0.1 \text{ fm}) \) can not be determined from the \( c\bar{c} \) and \( b\bar{b} \) spectra alone.

(ii) Such a short-range behavior becomes especially important for the \( 1S \) \( t\bar{t} \) state.

(iii) The QCD predicts the short-range behavior but the scale parameter \( \Lambda_{\overline{\text{MS}}} \) has to be determined experimentally.

Let us recalculate the \( t\bar{t} \) spectra in the similar line as ref.\(^2\) with the following improvements, i.e., for each value of \( \Lambda_{\overline{\text{MS}}} \) we have made the \( \chi^2 \)-fit to the \( c\bar{c} \) and \( b\bar{b} \) spectra. In this sense our parameters are chosen without any prejudice for each \( \Lambda_{\overline{\text{MS}}} \).

It turned out that the mass difference \( E(2S) - E(1S) \) tends to become larger with increasing \( \Lambda_{\overline{\text{MS}}} \). This tendency becomes larger for larger values of \( m_t \). (For a detail, see ref.\(^1\).)

**Spin-dependent forces and \( \Lambda_{\overline{\text{MS}}} \)**

Let me show that the fine and hyperfine splittings are very sensitive to \( \Lambda_{\overline{\text{MS}}} \). So, \( \Lambda_{\overline{\text{MS}}} \) may be restricted only from the data of \( c\bar{c} \) and \( b\bar{b} \).

The QCD radiative corrections to the spin-dependent forces have already been calculated (see Gupta et al\(^3\)). Using this scheme, let us compute the \( \Lambda_{\overline{\text{MS}}} \) dependence of spin-dependent forces.

The interactions are given by\(^3\)

\[
H_{SS} = \frac{8\pi \alpha_s}{9m^2} \delta^2 \cdot \delta_2 \left[ \left( 1 - \frac{\alpha_s}{12\pi} (26 + 9 \ln 2) \right) \delta(r) - \ldots \right]
\]

\[
= \frac{3\hat{r} \cdot \delta_1 \cdot \delta_2}{m^2}
\]

\[
H_T = \frac{3\hat{r} \cdot \delta_1 \cdot \delta_2 \cdot \hat{r} - \delta_1 \cdot \delta_2}{r^3} \left[ \ldots \right]
\]

\[
= B(3\delta_1 \cdot \delta_2 \cdot \hat{r} - \delta_1 \cdot \delta_2)
\]

\[(3)\]
Here $\mu$ is the renormalization scale. Let us choose $\mu$ so as to minimize the effect of higher-order terms.

Assuming $n_f = 4$ and using non-relativistic wave function (from our $Q\bar{Q}$ potential), we can compute expectation values of $H_{LS}$, $H_{LLS}$ and $H_T$ (see table).

Table

<table>
<thead>
<tr>
<th>$\Lambda_{\overline{\text{MS}}}$(MeV)</th>
<th>$\exp$</th>
<th>$c\bar{c}$</th>
<th>S-state</th>
<th>1P-state</th>
<th>A</th>
<th>B</th>
<th>$\mu$(GeV)</th>
<th>$\alpha_s$</th>
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<tr>
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<td>162</td>
<td>112</td>
<td>25.6</td>
<td>3.85</td>
<td>0.939</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

We can conclude with the following remarks.

(i) The fine and hyperfine splittings are very sensitive for the short-range potentials and rapidly increase as $\Lambda_{\overline{\text{MS}}}$ increases.

(ii) For $\Lambda_{\overline{\text{MS}}} \sim 200$ MeV, we can reproduce fine and hyperfine splittings of $c\bar{c}$ and $b\bar{b}$ states.
As you may notice from this table, $\Lambda_{MS} = 400 \text{ MeV}$ and $500 \text{ MeV}$ are completely ruled out.

**Toponium Prediction**

So, taking the preferred value $\Lambda_{MS} = 200 \text{ MeV}$, we can predict

$E(2S) - E(1S) = 670 \pm 15 \text{ MeV}$

$E(3S) - E(1S) = 995 \pm 15 \text{ MeV}$

$E(4S) - E(1S) = 1220 \pm 15 \text{ MeV}$

for $m_t = 40 \text{ GeV}$.

$\Gamma_{ee} (1S) = 4.9 \pm 0.1 \text{ keV}$

$\Gamma_{ee} (2S) = 1.66 \pm 0.02 \text{ keV}$

for $m_t = 40 \text{ GeV}$.

**References**


Bounds from $B^0 - \bar{B}^0$ mixing on charged-Higgs-boson masses and couplings in two-Higgs-doublet models are presented. These bounds are comparable to those obtained, with additional assumptions, from the neutral-$K$-system. The effects of the neutral Higgs bosons of these models on the spectrum and wave function of toponium is discussed. These effects could, in the future, lead to limits on, or the discovery of, these Higgs bosons.
1. Introduction

The Higgs sector remains the most elusive (and to some, unsatisfactory) feature of the standard model. It has often been suggested that it should be enlarged, or replaced altogether by bound states dynamically generated by a new strong interaction.\cite{1} Staying within "conventional" Higgs structures, there is no reason not to consider multiple Higgs doublets. In fact, many currently interesting theories, such as SUSY, left-right symmetric models, and superstring theories, require more than one doublet. Moreover, extra doublets can "decouple" the CP violation parameters $\epsilon$ and $\epsilon'$, which could prove useful if, with future measurements, the standard model is unable to account simultaneously for both values.

I will consider models with two Higgs doublets, although much of what I will discuss can be generalized to include more doublets. The new particles are two charged and two neutral bosons; an additional parameter is the vacuum expectation value (VEV) of the new doublet---or, equivalently, the ratio of the VEV's of the two doublets, if we fix an appropriate combination to be that of the standard model. Changing this VEV ratio changes the strength of the physical Higgs couplings and hence the size of the effects of the additional bosons; current physics, through the experimental absence of these effects, places limits on allowable values of the VEV ratio.

One first requires that flavor-changing neutral currents (FCNC) be absent at tree level. This can be done by imposing a discrete symmetry that forbids certain Higgs couplings. One scheme\cite{10} requires one Higgs doublet to couple only to up-type quarks (i.e., $u$, $c$, and $t$) and the other only to down-type quarks. Thus, for each set of quarks, a single Higgs doublet is responsible for both mass matrix and neutral Higgs couplings, so, as in the standard model, the two matrices diagonalize simultaneously and FCNC are absent at tree level. Another scheme\cite{11} allows only one Higgs doublet to couple to quarks at all, so that again the mass and coupling matrices diagonalize simultaneously.

In this talk I would like to discuss bounds on masses and couplings (VEV ratios) of charged Higgs bosons that follow from their effects on neutral $B$ meson mixing. I will compare these bounds to those derived from the $K_S^0 - K_L^0$ difference,\cite{4} and to those derived, with additional assumptions, from $CP$-violating effects in the $K$ system.\cite{5} I will then consider the effects of neutral-Higgs boson exchange on toponium physics. The Higgs exchange adds an attractive term to the interquark potential, which, for allowed values of the relevant parameters, can have dramatic effects on the spectrum and wave functions of toponium. However, distinguishing these effects from the variations of different, but theoretically acceptable, potentials, can present a problem.

This talk is based on work done with Gregory Athanasiu and Fred Gilman.\cite{6}
2. Limits from $B^0 - \bar{B}^0$ mixing

There are three box diagrams contributing in lowest order to $B^0 - \bar{B}^0$ mixing:

The first is the standard model contribution. The other two can only occur in a model with more than one Higgs doublet, as $H$ is the physical, charged Higgs. The $t$ quark contribution dominates the expression for the mass difference, since it is weighted by Kobayashi-Maskawa (KM) angle factors whose magnitudes are similar to those for the charm quark, while $m_t^2 >> m_c^2$. Thus we expect much tighter bounds than those found in the $K$-meson system; additionally, the freedom in choosing matrix elements, and in KM angle related factors is considerably smaller than in the $K$-meson system.

CLEO, at the $e^+e^-$ storage ring CESR, observes $B_d^0$ and $\bar{B}_d^0$ mesons pair produced near threshold, i.e., without other particles. Their decay amplitudes are therefore coherent, and the like sign to opposite sign dilepton ratio is equal to the "wrong"-sign lepton to "right"-sign lepton ratio for a single $B$ meson. This can be written as follows (neglecting effects of possible $CP$ violation)

$$r = \frac{N(l^+l^+) + N(l^-l^-)}{N(l^+l^-) + N(l^-l^+)} = \frac{\Gamma(B^0 \to l^- + \cdots)}{\Gamma(B^0 \to l^+ + \cdots)} = \frac{(\Delta M/\Gamma)^2}{2 + (\Delta M/\Gamma)^2}$$

where $\Delta M = M_S - M_L$ and $\Gamma = (\Gamma_L + \Gamma_S)/2$. CLEO's published upper limit on the mixing corresponds to

$$r < 0.30$$

which translates to the bound

$$|\Delta M/\Gamma| < 0.93.$$  

This bound uses the assumption $\tau_{B^0} = \tau_{B^+}$. Recently reported data could be interpreted as improving the bound, or as loosening the lifetime constraint.

Neglecting the $H - W$ diagram, and approximating the loop integrals, we find

$$\Delta M = \frac{G_F^2 f_B^2 m_B B s_1^2 s_2^2 m_l^2}{6\pi^2} \left[ 1 + \frac{1}{4} \left(\frac{\xi}{\eta}\right)^4 \frac{m_l^2}{M_H^2} \right],$$

where $\xi/\eta$ is the VEV of the Higgs doublet coupling to the up-type quarks divided by that of the doublet coupling to down-type quarks. Here $M_B$ is the $B$ meson mass, $s_1$ is the sine
of the first KM angle, and \( m_t \) is the t quark mass; \( f_B \) is defined analogously to the pion and kaon decay constants, \( f_\pi \) and \( f_K \); \( B_B \) is the bag factor for the \( B \) meson, and \( s_2 \) is the sine of the second KM angle. The first four parameters are fairly well-determined; we take \( M_B = 5.3 \) GeV, \( a_1 = .23 \), \( f_B = f_K = .16 \) GeV and \( m_t = 45 \) GeV (\( m_t \) could be larger, but this would only make our bound better, and it cannot be much smaller; we absorb any uncertainty in \( f_B \) into \( B_B \)).

In Fig. 1 I show our limit for various values of the bag factor and \( s_2 \).

![Fig. 1. Limits on \((\xi/\eta)^2 \) versus charged-Higgs-boson mass.](image)

As "reasonable" parameters we pick \( B_B = 1 \) and \( s_2 = 0.06 \). The dashed line is the above, approximate calculation, while the solid line is the limit resulting if we evaluate the loop integrals exactly, and include the Higgs-W cross term. I also show our limits for the conservative values \( B_B = 1/3 \) and \( s_2 = 0.04 \), and for the "optimistic" values \( B_B = 3/2 \) and \( s_2 = 0.08 \)—or equivalently, for improved experimental limits on \( B^0 - \bar{B}^0 \) mixing. For comparison, I show two previously calculated limits: the first, labeled \( ASW,^{[4]} \) is the limit from \( K \bar{K} \) mixing in the four quark model, and the second, labeled \( AG,^{[6]} \) is the limit determined by considering \( CP \) violation in the neutral \( K \) system. While this second bound is comparable to ours, it requires the additional assumption that the primary contribution to the \( CP \) violation parameter \( \epsilon \) be from the \( W - W \) diagram, rather than from those involving the Higgs, which may not be true.

With the unitarity constraint that the Higgs mass be less than of order 1 TeV, we have an Higgs-mass-independent bound of

\[
\frac{\xi}{\eta} \lesssim 10 - 15. \tag{2.5}
\]
3. Effects of allowed two-Higgs models on toponium physics

The neutral-Higgs ($H_0$) exchange contributes to the toponium potential, with the $H_0$ coupling enhanced by the ratio $\xi/\eta$ (I ignore possible mixing effects between the different neutral Higgs).

![Diagram]

The new term is an attractive Yukawa, in momentum space

$$- \left( \frac{\xi g m_t}{\eta 2 M_W} \right)^2 \frac{1}{m_H^2 + q^2} \quad \text{or} \quad - \left( \frac{\xi g m_t}{\eta 2 M_W} \right)^2 e^{-m_H} \frac{4\pi r}{4\pi r} \quad (3.1)$$

in coordinate space. This has the effect of increasing the wavefunction at the origin, since it pulls in the wavefunctions, and of lowering energy levels (increasing binding energies). It also increases the level spacings, since it affects the lowest lying states the most. The number of states below threshold could change, but not significantly, since states above the 3S are almost unaffected (this will be an unobservable effect, since with the expected resolution of SLC or LEP, we only hope to see the first 2 to 5 states out of the 11 to 13 states below threshold). Other quarkonia are, in principle, affected, though negligibly, due to their light mass.

Let us now consider the $2S/1P$ splitting. A theorem due to Martin states that if $\Delta V(r) > 0$ (true for all proposed quarkonia potentials), the nS state lies above the (n-1)P state, while if $\Delta V(r) < 0$ for all $r$ such that $dV/dr > 0$ (true for the Higgs potential), the nS state lies below the corresponding P state. Here we have a qualitative signature of the presence of the Higgs. However, the theorem requires a given condition on $\Delta V(r)$ to hold for all $r$. (The condition $dV/dr > 0$ holds for all $r$, for both potentials.) What happens when the Higgs dominates only near the origin? We might guess that relevant energy levels will be inverted if the Higgs term dominates below some relevant radius, perhaps that of the 2S or 1P. As $M_H$ increases, the range of the Higgs potential decreases and we need a larger value of $\xi/\eta$ to keep $\Delta V < 0$. This does give a qualitative picture of what happens. We find, numerically, the value of $\xi/\eta$ at which $E_{2S} = E_{1P}$, shown in Fig. 2 for two different potential models. The dashed line indicates the charged-Higgs-mass independent bound of the previous section. Level inversion occurs for points in parameter space above the curves shown.
Fig. 2. Minimum value of $\xi/\eta$ for which level inversion occurs.

We can make a semi-quantitative analysis of the wavefunction change by examining the singular part of the potentials. This goes from $-c/r$, where $c$ is some potential-model-dependent constant, to

$$-\left( c + \left( \frac{\xi g m_t}{\eta 2 M_W} \right)^2 \frac{1}{4\pi} \right) \frac{1}{r} .$$

(3.2)

But $|\psi(0)|^2 \propto (c m_t)^3$ for a Coulomb potential, so we expect the dependence

$$|\psi(0)|^{2/3} = |\psi(0)|_{\xi/\eta=0}^{2/3} \left[ 1 + a (\xi/\eta)^2 \right] ,$$

(3.3)

where the constant $a$ is deduced from Eq. (3.2). Numerically, we find this behaviour for small $\xi/\eta$ (5 to 10), although $a$ is smaller than calculated from Eq. (3.2), because of the screening effect of the factor $e^{-\alpha M_W r}$.

Table 1 shows the effect of the Higgs term for various potentials, Higgs masses, and VEV ratios. The Higgs can have striking effects; note, however, the similarity of the Cornell potential without a Higgs term to the Richardson potential with such a term. We have illustrated this problem by picking potentials that are not as physically well motivated as the Richardson potential. We would get similar, though less striking, effects by considering a QCD-inspired potential where one is free to vary $\Lambda_{\overline{MS}}$.

Figure 3 shows $R(e^+e^- \rightarrow \mu^+\mu^-)$, for toponium interfering with the $Z$, smeared with a beam width of 40 MeV, and $m_t = 47.5$ GeV. Note the qualitative similarity between the second and third figures.
Table 1. Calculated parameters of toponium.

$m_t = 50$ GeV (all units GeV to appropriate powers).

<table>
<thead>
<tr>
<th>Potential</th>
<th>$M_H$</th>
<th>$\xi/\eta$</th>
<th>$\langle r \rangle_{1s}$</th>
<th>$E_{2s} - E_{1s}$</th>
<th>$E_{2s} - E_{1p}$</th>
<th>$\Psi(0)_{1s}$</th>
<th>$\Psi(0)<em>{2s}/\Psi(0)</em>{1s}$</th>
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</table>

![Fig. 3](image)

**Fig. 3.** Effects of varying quarkonium potential.

4. Conclusions

In summary, we have seen that experimental limits on $B^0 - \bar{B}^0$ mixing yield strong limits on 2-Higgs models. For "reasonable" parameters, we have the bound

$$\left( \frac{\xi}{\eta} \right)^2 < 4.1 \frac{M_H}{m_t}.$$  \hspace{1cm} (4.1)

With unitarity, this yields an overall bound of $\xi/\eta \lesssim 10 - 15$. The enhanced neutral Higgs couplings allowed by this bound could strongly influence toponium spectroscopy. However, care must be taken in distinguishing this effect from uncertainties in potentials.
REFERENCES

NEW RESULTS ON $J/\psi$ DECAYS FROM DM2

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ABSTRACT :

$8.6 \times 10^6$ $J/\psi$ decays are analyzed. New results are presented on $q$, $S^*$, $f, f', D, n, n'$ production in hadronic decays. $n_c \rightarrow p\bar{p}$, $K\bar{K}$ are measured. $\xi$, $KK$, $\phi\phi$ are discussed. Results are presented on hadronic $J/\psi$ decays with strong apparent $SU(3)$ violation.

Talk given by A. FALVARD and G. SZKLARZ
1. INTRODUCTION

For 10 years, radiative decays of the $J/\psi$ have been intensively used to get strong production of $C = +1$ mesons. Particularly interesting are the states $E/\pi$ and $\theta$ whose gluonium content is still actively discussed. A new approach to the quark/glue content of mesons is to produce these particles in hadronic processes: $J/\psi \rightarrow (\rho, \phi, \psi) + X$. Indeed these decays are assumed to proceed via a simple mechanism where the flavours of vector and $X$ mesons are identical (Fig. 1.a). Though this scheme gives a consistent description of the various $J/\psi$ vector + Pseudoscalar decays, there is no strong theoretical argument to favor this diagramm 1). Nevertheless, experimental evidence that $J/\psi \rightarrow \omega f \gg \phi f$ and $J/\psi \rightarrow \phi f' \gg \omega f'$ indicates that the diagrams like Fig. 1.b which are doubly OZI suppressed are depreciated.

From this point of view, the observation of glueball candidates in hadronic decays is related to the quark content of these particles: $s\bar{s}$ is correlated to $\phi$ and $u\bar{u}$ or $d\bar{d}$ to $\rho, \omega$ as far as the vector mesons are ideally mixed.

This talk reports on $\theta$ production in $J/\psi \rightarrow \omega \theta, \phi \theta$ and on more standard mesons $S^*$, $f, f', D$. Of particular interest is the decay $J/\psi \rightarrow \omega S^*$ which is possibly observed with a rate of the same order of magnitude that $J/\psi \rightarrow \phi S^*$.

New results are presented on the $\eta_c \rightarrow \bar{p}p, K\bar{K}\pi$ decays which are important to get the $\eta_c \rightarrow \gamma\gamma$ partial width from R704(ISR) and $\gamma\gamma$ experiments. Angular analysis of the $\phi \phi$ system observed in radiative $J/\psi$ decays and the final point on non observation of $\xi^0 \rightarrow K\bar{K}$ are discussed.

At last, baryonic decays of the $J/\psi$ with an apparent strong SU(3) violation are presented.

The analysis is performed using the 8.6 millions of $J/\psi$ produced by the DCI in Orsay. This is determined from the $\rho \pi$ decay mode using the averaged branching ratio previously measured by other experiments.
2. NEW MEASUREMENTS IN J/Ψ → VECTOR + PSEUDOSCALAR

Among the J/Ψ → (ρ, ω, φ) + (π°, η, η′) decays, only ρη' and φπ° are not known from previous experiments. In fact, replacing the three gluons intermediate state by a virtual photon in Fig. 1.a, the φπ° mode is forbidden.

Analysing the J/Ψ → γρρ channel, (21±5) events are found from the interfering channels J/Ψ → ρη', ωη'(η'→νγ;ρ,ω+π+π-)

(Fig. 2). Assuming the two decay amplitudes to be real one relatively to the other, (18 ± 4) events are produced by the process J/Ψ → ρη'. This measurement is important since it allows to test the consistency of the dynamical hypothesis (Fig. 1.a) made to describe the decays J/Ψ→ν+p).

![Graph](image)

**Figure 2 - ργ mass in the decay J/Ψ → γρρ**

<table>
<thead>
<tr>
<th>Channels</th>
<th>Pseudoscalar decay mode</th>
<th>Number of events</th>
<th>Average branching ratios (x 10⁴)</th>
</tr>
</thead>
<tbody>
<tr>
<td>φη</td>
<td>π⁺π⁻π°</td>
<td>152 ± 12</td>
<td>7.2 ± .3 ± .1</td>
</tr>
<tr>
<td></td>
<td>π⁺π⁻γ</td>
<td>15 ± 4</td>
<td></td>
</tr>
<tr>
<td></td>
<td>γγ</td>
<td>163 ± 13</td>
<td></td>
</tr>
<tr>
<td>φη'</td>
<td>π⁺π⁻η</td>
<td>85 ± 9</td>
<td>4.0 ± .25 ± .1</td>
</tr>
<tr>
<td></td>
<td>ργ</td>
<td>92 ± 10</td>
<td></td>
</tr>
<tr>
<td>ρη'</td>
<td>ργ</td>
<td>18 ± 4</td>
<td>.78 ± .17 ± .12 *</td>
</tr>
</tbody>
</table>

**Table 1 - Summary about J/Ψ + Vector + Pseudoscalar**

* Assuming the phase between the amplitude A(ρη') and A(ωη') equal to zero

This theoretical scheme predicts (pᵥ = momentum of the vector):

$$R = \frac{|A(ωη')|^2 \times |A(ωπ°)|^2}{|A(ρη')|^2 \times |A(ρπ°)|^2} = 1; |A(V+P)|^2 = \frac{B \cdot R \cdot (J/Ψ + V+P)}{pᵥ^3}$$
Mixing the DM2 and MARKIII results \(3\), we get: \(R = 0.77 \pm 0.4\) which is in good agreement with theoretical expectation.

Moreover, new determinations of the branching ratios of the decays \(J/\psi \to \phi \eta, \phi \eta'\) have been made with large numbers of events (Table 1).

A surprising point is to observe a clear signal with \(D(1285)\) production in the decay \(J/\psi \to \phi \eta \pi^+ \pi^-\) (Fig. 3).

\[ R = 0.77 \pm 0.4 \]

\[ \text{which is in good agreement with theoretical expectation.} \]

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A surprising point is to observe a clear signal with \(D(1285)\) production in the decay \(J/\psi \to \phi \eta \pi^+ \pi^-\) (Fig. 3).

![Figure 3 - \(\eta \pi^+ \pi^-\) mass produced in front of \(\phi\)](image)

![Figure 4 - \(\eta \pi^+ \pi^-\) mass produced in \(\gamma \eta \pi^+ \pi^-\). A \(\delta\)-cut is applied on \(\eta \pi^+ \pi^-\) mass.](image)

A quite large branching ratio is obtained for this decay:

\[ \text{BR}(J/\psi \to \phi '' D'') \times \text{BR}(''D'' \to \eta \pi^+ \pi^-) = (1.77 \pm 0.4 \pm 0.25) \times 10^{-4} \]

There is no signal in the mass region around 1.395 GeV/c\(^2\) where a bump is found in the radiative channel \(J/\psi \to \gamma \eta \pi^+ \pi^-\). (Figure 4)

3. \(S^*, f, f', 0\) Production in Hadronic Decays of the \(J/\psi\)

The proposal that the scalar \(S^*\) and \(0\) observed at the \(KK\) threshold could be a \(KK\) molecule \(4\) makes interesting the measurement of \(J/\psi \to (\rho , \omega , \phi) + (S^*, 0)\). Moreover, the decays \(J/\psi \to \text{vector} + X\) are a natural place to look for glueball candidates, \(\pi\) and \(\eta\) produced in radiative decays. It constitutes a good method to observe the \(\pi \pi\) system produced in even wave below 1 GeV/c\(^2\) without \(\rho\) contamination. This is strongly correlated to the production of a possible \(0^{++}\) meson predicted in every dynamical description of strong interaction. Here are presented the results obtained about these topics from \(J/\psi \to (\phi , \omega) + (\pi^+ \pi^- , K^+ K^- , K^0 \bar{K}^0 , \phi)\).
a) $J/\psi + \phi \pi^+\pi^-$, $\omega \pi^+\pi^-$

Figures 5.a and 5.b give the $\pi\pi$ mass spectrum recoiling in front of $\phi$ and $\omega$ respectively. It shows a clear contribution of $J/\psi + \phi S^*$ which has been fitted by a FLATTE distribution 5). The following parameters are obtained:

\[ F(m) = \frac{m_0 \Gamma_m}{(m_0^2 - m^2) m_0 \left( \Gamma_m + \Gamma_K \right)} \]

\[ \Gamma_m = q_\alpha k_\alpha \] (k_K is analytically continued below the KK threshold)

\[ m_0 = 0.942 \pm 0.04 \text{ GeV/c}^2; \quad g_\pi = 0.26 \pm 0.05 \]

\[ g_K = 0.2 \text{ fixed} \]

Measured branching ratios are summarized in table 2. In the same reaction, a large bump which is not consistent with one BREIT-WIGNER distribution could be the mixing of $\phi f$ and $\phi f'$ or $\phi c(1.3)$. At last, in the region 1.7 - 1.9 GeV/c^2 could appear another structure (\theta ?). The analysis is in progress to test the consistency of these hypothesis.

The $\omega \pi^+\pi^-$ decay shows also a very rich $\pi\pi$ mass spectrum (Fig. 5.b). Apart a large $\omega f$ contribution, two facts have to be pointed out:

* the $\pi\pi$ bump produced below 1 GeV/c^2 is not a pure $0^{++}$ state. This is shown in figure 6 which gives the angle between the $\pi^-$ direction and the $X = \pi^+\pi^-$ system direction in its rest frame. This angular distribution is less and less isotropic as the $\pi\pi$ mass increases.

* An excess of events is seen just below the KK threshold which could be the evidence of the $\omega S^*$ contribution. If we believe in this
hypothesis, the branching ratio of the decay \( \psi \rightarrow \omega S^* \) is of the same order that \( \psi \rightarrow S^* \).

Table 2 summarizes the various branching ratios obtained in \( \psi \rightarrow \pi^\pm \pi^- \) and \( \omega^+ \pi^- \).

<table>
<thead>
<tr>
<th>Channel</th>
<th>B.R. ((\times 10^{-2}))</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \phi \pi^+ \pi^- )</td>
<td>((7.5 \pm 1.6) \times 10^{-2})</td>
</tr>
<tr>
<td>( \phi S^* .(S^* \rightarrow \pi \pi) )</td>
<td>((2.4 \pm 5) \times 10^{-2})</td>
</tr>
<tr>
<td>( 0.85 &lt; M_{\pi \pi} &lt; 1.15 )</td>
<td></td>
</tr>
<tr>
<td>( \omega \pi^+ \pi^- )</td>
<td>(0.66 \pm 0.06)</td>
</tr>
<tr>
<td>( \omega f )</td>
<td>(0.40 \pm 0.06)</td>
</tr>
<tr>
<td>( \rho A_2 )</td>
<td>(0.86 \pm 0.13)</td>
</tr>
</tbody>
</table>

Interesting structures appear in the \( \pi \pi \) mass spectrum (Figure 7-b). A more detailed analysis must be performed on this subject where exotic states where suggested 6).

b) \( \psi \rightarrow \phi K^\pm K^- \)

The decay \( \psi \rightarrow \phi K^+K^- \) is very clean since it is only kinematically contaminated by \( \psi \rightarrow \pi^+ \pi^- p\bar{p} \) which is easily rejected by time of flight informations. Moreover, the channel \( \psi \rightarrow 2(K^+K) \) is small and gives little combinatorial background. The same situation occurs for \( \psi \rightarrow \phi K^0 S K^0 S \) which is strongly kinematically constrained. Figures 8(a)-(b) give the KK mass spectra for \( K^+K^- \) and \( K^0 S K^0 S \) final states respectively.

The KK mass spectrum is dominated by \( f' \) production but the channel \( \phi K^+K^- \) which has a large statistics reveals the contribution of another structure \( X \) on the high mass side of the \( f' \). Two BREIT-WIGNER distributions added incoherently give the following parameters for this state \( X \):
m_X = (1648 ± 10) MeV/c² ; \Gamma_X = (13±± 28) MeV/c²

with the following branching ratios:

B.R. \((J/\psi \to \phi f') \times BR(f' \to KK) = (4.6±0.5) \times 10^{-4}\)

B.R. \((J/\psi \to \phi X) \times BR(X \to KK) = (2.7±0.3) \times 10^{-4}\)

But this mass distribution is also consistent with a \(0\) production with \(M_0 \sim 1700\) MeV/c² interfering strongly with the \(f'\).

Unfortunately, the statistics is small in this decay mode.

Moreover, the large production of \(KK\) pairs at low mass is not inconsistent with the \(S^*\) excitation observed in \(J/\psi \to \phi \pi^+\pi^-\).

The analysis of \(J/\psi \to \omega K^+K^-\) and \(K_0^S K_0^S\) is very preliminary. In particular, the analysis of \(\omega K^+K^-\) is not performed on the total statistics available. But first results (Fig. 9) highlight a large activity in the \(0\) region. Assuming the bump around \((1.65 \sim 1.8)\)GeV/c² in the \(K_0^S K_0^S\) spectrum to come from the \(0\) excitation, we get:

\[ BR(J/\psi \to \omega 0) \times BR(\omega \to KK) = (4. \pm 0.8 \pm 1.3) \times 10^{-4} \]

To summarize on \(KK\) production, we emphasize that a large activity is found in the \(\omega\) region from hadronic decays \(J/\psi \to \phi KK, \omega KK\). In both channels, we get a consistency of the observed signals with \(f' + 0\).
Nevertheless, the $\phi K^+K^-$ mode could indicate the presence of a new state with mass $(1648 \pm 10)$ MeV/c$^2$ and width $(133 \pm 28)$ MeV/c$^2$.

3. NEW RESULTS IN RADIATIVE DECAYS

a) $J/\psi \rightarrow \gamma \eta_c$

Recent experimental limit on $pp \rightarrow \eta_c \rightarrow \gamma\gamma$ and measurement of $\gamma\gamma \rightarrow \eta_c \rightarrow K_0S K^+\pi^-$ use the branching ratios of $(\eta_c \rightarrow pp), (\eta_c \rightarrow K_0S K^+\pi^-)$. In fact, the B.R. of the channel $J/\psi \rightarrow \gamma \eta_c \rightarrow pp$ is the smallest one observed in $J/\psi$ decay and large uncertainty remains on this measurement. So we made new measurements of the decays:

$J/\psi \rightarrow \gamma \eta_c, (\eta_c \rightarrow pp, KK^*)$

We get the following results:

$$BR(J/\psi \rightarrow \gamma \eta_c) \times BR(\eta_c \rightarrow pp) = (1.32 \pm 3 \pm 1.3) \times 10^{-5}$$

$$BR(\eta_c \rightarrow K_0S K^+\pi^-) = (2.3 \pm 4 \pm 0.6) \times 10^{-4}$$

$$BR(\eta_c \rightarrow K^+K^-\pi^0) = (1.46 \pm 3 \pm 0.22) \times 10^{-4}$$

which are consistent with measurements previously published by other experiments.

b) $J/\psi \rightarrow \gamma \phi \phi$

Fig 10 gives the $\phi \phi$ mass observed in the decay $J/\psi \rightarrow \gamma \phi \phi$. The angular analysis of the mass region around 2.2 GeV/c$^2$ is done and shows clearly (Fig. 11) that the enhancement observed in this range is not consistent with a positive parity state. A $\zeta(2.2) \rightarrow KK^*$ should be a positive parity particle. So the enhancement is not consistent with a $\zeta \rightarrow \phi \phi$ decay.

Figure 10 - $\phi \phi$ mass

Figure 11 - Angle between the decay planes of both $\phi$
c) $J/\psi \rightarrow \gamma \xi(2.2)(\xi + K\bar{K})$

Since the International Conference in Bari 8), a lot of work has been devoted to check the analysis of the decay $J/\psi \rightarrow \gamma K^+K^-$ in order to get a final result about observation or non observation of a $\xi(2.2)$ signal. The result of this work is to confirm the non observation of the $\xi(2.2)$ in both the $K^+K^-$ and $K_0S\bar{K}_0S$ channels:

$$B.R.(J/\psi \rightarrow \gamma \xi) \times BR(\xi \rightarrow K^+K^-) < 2.4 \times 10^{-5} \text{ (95 \% C.L.)}$$
$$B.R.(J/\psi \rightarrow \gamma \xi) \times BR(\xi \rightarrow K_0\bar{K}_0) < 2.0 \times 10^{-5} \text{ (95 \% C.L.)}$$

4. OBSERVATION OF SU(3) VIOLATIONS IN J/ψ DECAYS

Baryonic decays of the $J/\psi$ have been used to measure decays $1 \rightarrow 8 \times \Omega$, $8 \times 10$ like $J/\psi \rightarrow \Xi^+(1380)\Xi^-(1200)+C.C., \Xi^-(1320)\Xi^+(1530)+C.C., \Xi^+(1320)\Xi^-(1530)+C.C.$. Authors 9) have pointed out the interest to measure relative rates of $\Xi\Xi^*$ and $\Xi\Xi^*$ to test the mechanism of SU(3) violation. Large branching ratios have been observed in these channels. Preliminary results are summarized in Table 3.

<table>
<thead>
<tr>
<th>Channel</th>
<th>B.R.(x10^4)</th>
<th>Channel</th>
<th>B.R.(x10^4)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\Xi^- (1380)\Xi^+(1200)+C.C.$</td>
<td>3. ±3±.75</td>
<td>$\Xi^- (1380)\Xi^+(1380)+C.C.$</td>
<td>10 ±6±1.2</td>
</tr>
<tr>
<td>$\Xi^+(1380)\Xi^- (1200)+C.C.$</td>
<td>3.4±.4±.75</td>
<td>$\Xi^+(1380)\Xi^- (1380)+C.C.$</td>
<td>11.9±.7±1.4</td>
</tr>
<tr>
<td>$\Xi^- (1320)\Xi^+(1530)+C.C.$</td>
<td>6.9±.8±1.5</td>
<td>$\Xi^- (1320)\Xi^+(1320)+C.C.$</td>
<td>7 ±6±.8</td>
</tr>
<tr>
<td>$\Xi^+(1320)\Xi^- (1530)+C.C.$</td>
<td>4.5±.8±1.</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

Table 3 - B.R. for $J/\psi \rightarrow 8\times \Omega, 10\times \Xi$

Table 4 - B.R. for $J/\psi \rightarrow 8 \times \Xi$

These processes are not negligible in front of decays allowed by SU(3) symmetry (Table 4).

Figures 12(a) (b) (c) and Table 3 illustrate the results obtained in these channels.
Figure 12

(a) $\Lambda\pi^\pm$ mass observed in front of a missing mass of $\Sigma(1200)$. A clear $\Sigma^*(1380)$ is observed on a background from $\Lambda\pi^-\Sigma(1200)$

(b) Momentum of the $\Lambda\pi^-\pi^+$ system measured in the region of $\Xi^0$. The peak is from $\Xi^0\Xi^0$ (C.C. is summed)

(c) Missing mass of in front of $\Xi^-$

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SOME RECENT RESULTS OF J/ψ DECAYS FROM MARK III*

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Abstract

The Mark III Collaboration presents a preliminary analysis of a systematic study of direct J/ψ decays into the vector \( (J^P_C = 1^-) \) and the tensor \( (J^P_C = 2^{++}) \) mesons based on a data sample of \( 5.8 \times 10^6 \) J/ψ's. Some vector-scalar decays are discussed. Among the baryonic modes, observation of some SU(3) violating decays are reported.

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Introduction

The decays of $J/\psi$ has proven to be a rich laboratory for finding new and exciting states e.g. $\xi, \phi, \xi$ in the last decade. A popular searching ground for glueballs has been the radiative decays. On the other hand, pure hadronic decays of $J/\psi$ into mesons and baryons have been used as a factory to study light quark spectroscopy in 'clean surroundings' where the initial state is well defined ($J^{PC} = 1^{--}$). Decays with a $\phi$ or $\omega$ help one understand the quark content of the meson recoiling against them. Most of the recent results presented here are on hadronic decays.

Figure 1 shows the principal decay mechanisms of the $J/\psi$ in order of relative strength. Further similar or higher order diagrams exist, but are not shown. As an interesting consequence of quark correlations arising from $1(a)$ and $1(b)$, we have analyzed processes of the type $J/\psi \rightarrow (1^{--}) + (2^{++})$, as well as some of the type $J/\psi \rightarrow (1^{--}) + (0^{++})$, to complement the previous study of $J/\psi \rightarrow (1^{--}) + (0^{--})$. Particularly for the scalar nonet questions of glueball candidates or four quark bound states have become intriguing.

Of special interest for quark correlation studies are the recoils against $\phi$ and $\omega$. According to the quark line diagram in Figs. 1(a,b), $\phi$ being pure $s\bar{s}$ and $\omega$ being pure $u\bar{u}$ and $d\bar{d}$, project out respectively the strange and the non-strange quark content in the recoil system. Some of these decay modes are described below.

$\pi^+\pi^-$ Recoil from $\omega$ and $\phi$

The final state studied was the decay $J/\psi \rightarrow \pi^+ \pi^- \pi^+ \pi^- \pi^0$. Figure 2(a) shows the 5C constrained $\pi^+\pi^-\pi^0$ mass spectrum with four combinations per event. A clear $\omega$ peak over the combinatorial background is evident. Figure 3(a) shows the final $\pi^+\pi^-$ spectrum recoiling from the $\omega$. The peak at $\sim 1270$ MeV/c$^2$ is identified with the $f$, the $I = 0$ member of the $2^{++}$ nonet, containing mostly non-strange quarks. The branching ratio of $J/\psi$ into $\omega f$ was measured as $B(J/\psi \rightarrow \omega f) \times B(f \rightarrow \pi^+\pi^-) = (27.7 \pm 1.4 \pm 7.0) \times 10^{-4}$. The broad enhancement at $\sim 500$ MeV/c$^2$ has been seen in previous experiments, but is not yet understood. The structure near 1000 MeV/c$^2$ also needs further investigation. The inclusive branching ratio of $J/\psi \rightarrow \omega\pi^+\pi^-$ is shown in Table 1.

The final state $K^+K^-\pi^+\pi^-$ was analyzed to study the $\pi^+\pi^-$ system recoiling against a $\phi$. Figure 2(b) shows the $K^+K^-$ spectrum in the $\phi$ region. The $\phi$ signal is almost background free. Figure 3(b) shows the $\pi^+\pi^-$ mass spectrum. A clear peak at $\sim 980$ MeV/c$^2$ corresponding to the $S^*$ is seen. The spin-parity of the $S^*$ is known to be $0^{++}$. However, whether it is a simple $q\bar{q}$ resonance, two states close to each other$^2$ (a pole in $\pi\pi$ and another in $K\bar{K}$), or a four quark state$^3$ has been debated. In the $q\bar{q}$ scheme it is usually taken as primarily an $s\bar{s}$ bound state, but being below the $K\bar{K}$ threshold, decays mainly to $\pi\pi$. However, once the invariant mass is above the $K\bar{K}$ threshold, it decays mostly to $K\bar{K}$ and hence the sharp fall off on the high mass side of the $S^*$ in the $\pi\pi$ spectrum. Fitting the $S^*$ spectrum to the standard coupled channel Flatté parametrization$^4$, the branching ratio of $J/\psi$ into $\phi S^*$ was $Br(J/\psi \rightarrow \phi S^*) \times Br(S^* \rightarrow \pi^+\pi^-) = (2.3 \pm 0.3 \pm 0.6) \times 10^{-4}$. The structures in the mass region of 1300 to 1550 MeV/c$^2$ have been speculated to be $f$ and/or $\epsilon$. 

Fig. 1. (a) Three gluon annihilation. (b) electromagnetic decay proceeding via $c\bar{c}$ annihilation into one photon. (c) Electromagnetic decay into a final state of one photon and two gluon color singlet. (d) Doubly disconnected diagrams.
and the structure at \( \sim 1700 \text{ MeV}/c^2 \) to be the \( \theta \). The measurement of the inclusive branching ratio of \( J/\psi \rightarrow \phi\pi^+\pi^- \) is included in Table 1.

<table>
<thead>
<tr>
<th>( J/\psi ) Decay Modes</th>
<th>MK3</th>
<th>DM2</th>
<th>Particle Data Group</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \omega f )</td>
<td>49.3 ± 2.5</td>
<td>40 ± 6</td>
<td>23 ± 8</td>
</tr>
<tr>
<td>( \omega f' )</td>
<td>&lt; 1.2</td>
<td>--</td>
<td>&lt; 1.6</td>
</tr>
<tr>
<td>( \phi f' )</td>
<td>( \cdot f' \rightarrow KK )</td>
<td>6.4 ± 0.6</td>
<td>4.6 ± 0.5</td>
</tr>
<tr>
<td>( \phi f )</td>
<td>--</td>
<td>--</td>
<td>&lt; 3.7</td>
</tr>
<tr>
<td>( K^+(892)K^+(1430) + cc )</td>
<td>56 ± 4 ± 8.4</td>
<td>--</td>
<td>67 ± 26</td>
</tr>
<tr>
<td>( \rho A_2 )</td>
<td>118 ± 8 ± 29</td>
<td>86 ± 3 ± 13</td>
<td>84 ± 45</td>
</tr>
<tr>
<td>( \phi S^* )</td>
<td>( \cdot S^* \rightarrow \pi^+\pi^- )</td>
<td>--</td>
<td>2.6 ± 0.6</td>
</tr>
<tr>
<td>( \omega \pi^+\pi^- )</td>
<td>78 ± 1 ± 16</td>
<td>66 ± 10 ± 6</td>
<td>68 ± 19</td>
</tr>
<tr>
<td>( \omega KK )</td>
<td>17.2 ± 0.8 ± 3.4</td>
<td>--</td>
<td>16 ± 10</td>
</tr>
<tr>
<td>( \phi \pi^+\pi^- )</td>
<td>9 ± 0.4 ± 2.3</td>
<td>7.5 ± 0.3 ± 1.5</td>
<td>21 ± 9</td>
</tr>
</tbody>
</table>

\( K^+K^- \) Recoil from \( \omega \) and \( \phi \)

The \( K^+K^- \) recoil spectra against the \( \omega \) and the \( \phi \) are presented in Fig. 2 of the contribution by L. Köpke.\(^5\)

The recoil spectrum from the \( \omega \) was studied both in \( K^+K^- \) and \( K_sK_s \) modes. No striking signal corresponding to any of the standard \( q\bar{q} \) resonances is apparent in the \( K^+K^- \) spectrum. However, a clear peak is seen at \( 1731 \pm 10 \pm 10 \text{ MeV}/c^2 \) with a width of \( 110\pm10 \pm 15 \text{ MeV}/c^2 \), which are consistent with the parameters of the \( \theta \). The branching ratio of \( J/\psi \) into this decay mode was \( Br(J/\psi \rightarrow \omega X(1731)) \times Br(X \rightarrow KK) = (4.5_{-1.1}^{+1.2} \pm 1.0) \times 10^{-4} \). Ignoring any possible interference, an upper limit on \( \omega f' \) production was calculated at 90% C.L. as \( Br(J/\psi \rightarrow \omega f') \times Br(f' \rightarrow KK) < 1.2 \times 10^{-4} \). The \( K_sK_s \) mass spectrum, with much less statistics, reproduced the major features of the charged mode. The inclusive branching ratio of \( J/\psi \) into \( \omega KK \) is included in Table 1.

The \( K^+K^- \) mass spectrum recoiling against the \( \phi \) shows a clear enhancement at the \( f'(1525) \) mass. This is the \( I = 0 \), primarily \( s\bar{s} \) member of the \( 2^{++} \) nonet. However, a high mass shoulder to the \( f' \) is evident. The branching ratio for \( f' \) was measured as \( Br(J/\psi \rightarrow \phi f') \times Br(f' \rightarrow KK) = (6.4 \pm 0.6 \pm 1.6) \times 10^{-4} \). The details of this channel are described by L. Köpke.\(^5\)

**Table 1. Compilation of the \( J/\psi \) decay branching ratios**

![Fig. 2](https://via.placeholder.com/150)
Comparison with the Radiative Decay Modes: $\gamma\pi^+\pi^-$ and $\gamma K^+K^-$

The radiative decays of the $J/\psi$ have been very rich in new physics. The $\gamma$, $\phi$ and $\omega$ all have the same $J^{PC}$ quantum numbers. While the $\omega$ and the $\phi$ project out the non-strange and the strange quark respectively in the recoiling $q\bar{q}$ system (Figs. 1(a,b)), the two-gluon system that decays into the observed final state in the radiative decays (Fig. 1(c)) is flavor independent. Consequently, the radiative decays are independent of the quark flavor. Figures 3(a), (b) and (c) present the $\pi^+\pi^-$ spectra recoiling from $\omega$, $\phi$ and $\gamma$ respectively. Figure 3(a) has a clear $f'$ and Fig. 3(b) has a clear $S^*$ peak in accord with the flavor correlation expected from Figs. 1(a) and (b). Figure 3(c) shows a large $f'$ peak, a shoulder, probably from the $f'$ and a $\theta$ peak. The lower peak at $\sim 800$ MeV/c$^2$ originates from $\rho^0\pi^0$ feed through, where a photon from the $\pi^0$ is undetected.

The $K^+K^-$ recoil spectra against the $\gamma$, $\omega$, and $\phi$ are presented by L. Köpke. The radiative decay shows clear $f'$ and $\theta$ production, where copious $f'$ production is visible in the recoil spectrum against the $\phi$. An interesting point to note is that while the continuum process of $\phi\pi\pi$ production proceeds through a double $OZI$ suppressed diagram e.g., Fig. 1(d), the $\omega K\bar{K}$ mode can proceed through either a double $OZI$ suppressed diagram or a sequential decay mechanism. The subject of $\theta$ production in association with $\phi$ and $\omega$ has been discussed in the talk by L. Köpke and is not covered here.

Recoils from $K^*(892)$ and $\rho$

To check the quark correlations further, recoils against the isodoublet, $K^*(892)$, and the isovector $\rho$ were analyzed. In the case of the $K^*(892)$, the final state considered was $K^+\pi^-K^-\pi^+$. Figure 4(a) shows the plot of $K^-\pi^+$ vs. $K^+\pi^-$. A band due to $K^*(892)^0$ ($K^*(892)^0$) production is apparent in $K^-\pi^+$ ($K^+\pi^-$). Figure 4(b) shows the $K^+\pi^-$ spectrum recoiling from the $K^*(892)^0$. A large peak at $\sim 1430$ MeV/c$^2$ corresponding to the production of $K^*(1432)^0$, the isodoublet partner of the $2^+$ nonet is seen. The production branching ratio for this mode was measured as $\text{Br}(J/\psi \rightarrow K^*(892)^0 K^*(1430)^0) + cc = (56 \pm 4 \pm 8.4) \times 10^{-4}$. A small but clear peak is evident in Fig. 4(b) due to $K^*(892)^0$ production, which is forbidden by invariance of the strong interaction under SU(3), and therefore points to the breaking of SU(3) or the presence of a substantial electromagnetic amplitude in $J/\psi$ decays (Fig. 1(b)).

The recoil spectrum against the $\rho$ was studied in the $\eta\pi^+\pi^-\pi^0$ final state, where the $\rho^0$ ($\rho^+$) decayed into $\pi^+\pi^-\pi^0$ ($\pi^0\pi^\pm$) and the resonances were searched for in the $\eta\pi^0$ ($\eta\pi^\pm$) decay mode. Figure 5(a) shows the recoil $\eta\pi$ spectrum from all three charged states of the $\rho$. Two peaks corresponding to $\delta$ and $A_2$...
production are seen above the background at $\sim 980$ and at $1320$ MeV/c\(^2\) respectively. Figure 5(b) shows the same spectrum corresponding to the sidebands of the $\rho$. The correlation of $A_2$ production in association with $\rho$ is clear; however, the interpretation of $\delta$ production is being pursued. The branching ratio for $\rho A_2$ production was calculated to be $Br(J/\psi \rightarrow \rho A_2) = (118 \pm 8 \pm 29) \times 10^{-4}$.

![Figure 4](image4.png)

![Figure 5](image5.png)

**Fig. 4.**

**Fig. 5.**

**Conclusion of Vector-Tensor and Vector-Scalar Decays**

Table 1 summarizes the measured decay modes of $J/\psi$ into the vector-tensor and vector-scalar channels. The flavor correlations assumed by the dominance of diagrams 1(a) and (b) are clearly seen. However the differences in measured branching ratios for associated production of $\omega$ and $\phi$ seem to point to the presence of a large amount of SU(3) violation. This was seen and measured,\(^1\) along with the strong and the electromagnetic amplitudes, in the earlier systematic study of $J/\psi$ decay into the vector-pseudoscalar nonets, which led to a measurement of the quark contents of the $\eta$ and the $\eta'$. In the case of the scalar mesons, a systematic study has begun, and promises to be interesting.

**Some Baryonic Decays of $J/\psi$**

The following decay modes have been analyzed to search for possible SU(3) violating decays:

- $J/\psi \rightarrow \Xi^- \Xi(1530)^+ + cc$ \hspace{1cm} (1)
- $J/\psi \rightarrow \Xi^0 \Xi(1530)^0 + cc$ \hspace{1cm} (2)
- along with the “elastic” process \hspace{1cm} $J/\psi \rightarrow \Xi^- \Xi^+$. \hspace{1cm} (3)

The two decay modes of the $\Xi(1530)^+$ were analyzed separately. Figure 6(a) shows
the missing mass spectrum recoiling against the $\Xi^-$, where the $\Xi(1530)^+$ decays through $\Xi^0$. The elastic peak at $1320$ MeV/c$^2$ is very prominent, while a peak at $1535$ MeV/c$^2$ from the $\Xi(1530)^+$ production is also evident.

For the decay mode where the $\Xi(1530)^+$ decays through $\Xi^0$, the missing mass off the reconstructed $\Xi^-$ and the bachelor $\pi^+$ (not belonging to the $\Lambda$) combination was plotted as shown in Fig. 6(b). A peak at $1315$ MeV/c$^2$ due to $\Xi^0$, corresponding to $\Xi(1530)^+$ decaying into $\Xi^0$ and $\pi^+$, is seen. The low mass peak at $1120$ MeV is due to $\Lambda$ from the elastic process, with the $\Xi^+$ decaying into $\Lambda\pi^+$. The charge conjugate modes were analyzed separately. The weighted average of the branching ratios derived from the two decay modes of the $\Xi(1530)^+$ was $Br(J/\psi \rightarrow \Xi^-\Xi(1530)^+ + cc) = (8.8 \pm 0.7 \pm 2.2) \times 10^{-4}$, while the branching ratio of the elastic reaction was $Br(J/\psi \rightarrow \Xi^-\Xi^+) = (8.6 \pm 0.5 \pm 2.0) \times 10^{-4}$. Since $J/\psi$ is an SU(3) singlet, the decay into $\Xi\Xi^*$, i.e into $B_8\bar{B}_{10}$ (octet-decuplet), is SU(3)$_{\text{strong}}$ violating. One obtains the ratio of the SU(3) violating to SU(3) allowed decay modes as $\frac{1}{2} \times Br(J/\psi \rightarrow \Xi^-\Xi(1530)^+ + cc)/[Br(J/\psi \rightarrow \Xi^-\Xi^+)] = 0.50 \pm 0.05 \pm 0.2$. In the neutral mode (2), at most a small signal was observed. In addition, Monte Carlo studies showed possible feedthrough from other channels. As a conservative approach, we therefore quote an upper limit: $Br(J/\psi \rightarrow \Xi^0\Xi(1530)^0 + cc) < 4.1 \times 10^{-4}$ at $90\%$ C.L.

A strong interaction decay respects isospin invariance. The $J/\psi$, being an isosinglet, should have equal branching ratios for reactions (1) and (2) if the decay proceeds by the strong interactions. The large difference in the two rates therefore shows the presence of an electromagnetic amplitude. Hence the origin of the observed branching ratios of (1) and (2) might arise from an interference between the electromagnetic and the strong amplitudes. Further studies are being performed.

Acknowledgements

I wish to thank Professor F. Gilman for useful discussions.

References

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A COMPARISON OF RADIATIVE AND HADRONIC J/ψ DECAYS

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ABSTRACT

Recent results from the Mark III Collaboration on radiative and hadronic J/ψ decays are presented. The large data sample of $5.8 \times 10^6$ observed J/ψ events allows for the observation of rare phenomena, as well as systematic studies of related decays. The status of the $\xi(2230)$ and searches for additional $\xi$ decay modes are discussed. The complete set of allowed $J/ψ \rightarrow \gamma 1-1-$ decays is used to study the $\eta_c$ decay mechanism. A comparison of $J/ψ \rightarrow \{\gamma, \omega, \phi\} X$ decays is made. The final states $X = \{K\bar{K}, K\bar{K}\pi, \eta\pi\pi\}$ are studied, and conclusion on the $\phi(1720)$ and $\omega(1460)$ resonances are drawn.

INTRODUCTION

Interest has lately been centered on radiative $J'/\gamma$ decays since this field is considered to be an excellent "hunting ground" for gluonic matter. Indeed, several new phenomena such as the $\zeta(1460)$, $\theta(1720)$, $\xi(2230)$, and pseudo scalar structures in the $\rho\rho$ and $\omega\omega$ final states have been discovered. However, the observation of a state in the radiative $J'/\gamma$ decay does not necessarily mean that it is a glueball; after all, well-known $q\bar{q}$ resonances have also been seen. It seems that only a systematic comparison of the gluon enriched radiative decays with those produced by other mechanisms can disentangle the situation. Especially interesting are flavor dependent channels of the type $J'/\gamma \rightarrow \omega X$ and $J'/\gamma \rightarrow \phi X$, as one expects a quark correlation between the vector meson and the recoiling system.

STATUS OF THE $\xi$

The observation of the $\xi(2230)$ particle was one of the true surprises in the study of radiative $J'/\gamma$ decays.\cite{1} Enhancements were seen in the $K^+K^-$ (Fig. 2a) and $K\pi\pi$ final states with a significance of 4.6 and 3.2 s.d., respectively. The resonance parameters were measured as $m_\xi = 2230 \pm 6 \pm 14$ MeV, $\Gamma = 26^{+22}_{-10} \pm 17$ MeV, and $B(J'/\gamma \rightarrow \gamma\xi) \cdot B(\xi \rightarrow K^+K^-) = (4.2^{+1.7}_{-1.4} \pm 0.8) \times 10^{-5}$ in the $K^+K^-$ mode and $m_\xi = 2232 \pm 7 \pm 7$ MeV, $\Gamma = 18^{+23}_{-15} \pm 10$ MeV and $B(J'/\gamma \rightarrow \gamma\xi) \cdot B(\xi \rightarrow K\pi\pi) = (3.1^{+1.6}_{-1.3} \pm 0.7) \times 10^{-5}$ in the $K\pi\pi$ mode.

Since the DM2 Collaboration did not confirm the $\xi$ signal, the Mark III Collaboration intensified its search for other decay modes. One of the more exotic $\xi$ decays proposed is the OZI violating decay $\xi \rightarrow \omega\phi$, which might be expected if the $\xi$ were a hybrid state.\cite{2}

$J'/\gamma \rightarrow \gamma\omega\phi$ AND A SYSTEMATIC STUDY OF $\eta_c$ DECAYS INTO TWO VECTOR MESONS

The $J'/\gamma \rightarrow \gamma\omega\phi$ decay was observed in the $K^+K^-\pi^+\pi^-\gamma\gamma$ final state with $B(J'/\gamma \rightarrow \gamma\omega\phi) = (1.40 \pm 0.25 \pm 0.28) \times 10^{-4}$. The $\omega\phi$ mass spectrum in Fig. 1 does not show significant structure leading to the upper limits $B(J'/\gamma \rightarrow \gamma\xi) \cdot B(\xi \rightarrow \omega\phi) < 5.9 \times 10^{-5}$ and $B(J'/\gamma \rightarrow \gamma\eta_c) \cdot B(\eta_c \rightarrow \omega\phi) < 1.3 \times 10^{-5}$ at 90% C.L.

With this measurement, all radiative decays of the $J'/\gamma$ into pairs of vector mesons were measured by the Mark III Collaboration. According to ref. 3, this set can be used to determine the relative contributions of the three mechanisms relevant to $\eta_c \rightarrow 1^-1^-1^-$ decays. The observed decay pattern indicates that the $\eta_c \rightarrow 1^-1^-1^-$ decay rate increases with the number of strange quarks in the final state, a SU(3)-breaking pattern very different from the one observed in $J'/\gamma \rightarrow 1^-0^+$ decays.\cite{4} Since the $\eta_c$ mainly hadronizes via two gluon exchange, the study of its decays should also further our understanding of gluonia decays.

Fig. 1: $\omega\phi$ mass spectrum from $J'/\gamma \rightarrow \gamma\omega\phi$ reaction. Background estimate shown in lower plot.
As explained in the introduction, a comparison between corresponding radiative and hadronic decays might be very helpful to learn more about the "glueball" candidates $\theta$ and $\iota$, as well as other phenomena observed in radiative $J/\psi$ decays.

**Study of the $K\bar{K}$ Final State**

In this section we are mainly interested in the study of the $\theta(1720)$, which is clearly observed in the $J/\psi \rightarrow \gamma K^+K^-$ decay (Fig. 2 a).

$J/\psi \rightarrow \omega K\bar{K}$: This channel was studied in the $K^+K^-$ and the $K_sK_s$ final states. The doubly OZI violating channel is produced with $B(J/\psi \rightarrow \omega K\bar{K}) = (17.2 \pm 0.8 \pm 3.4) \times 10^{-4}$. The $K^+K^-$ mass spectrum recoiling against the $\omega$ shows clear evidence for a structure at $1731 \pm 10 \pm 10$ MeV with a width of $100_{-45}^{+45} \pm 15$ MeV (Fig. 2 b). The enhancement, which is consistent with the $\theta$ parameters, is produced with

$$B(J/\psi \rightarrow \omega f') \cdot B(f' \rightarrow K\bar{K}) < 1.2 \times 10^{-4}$$

at 90\% C.L.. The $K_sK_s$ mass spectrum, with much less statistics, confirms the features of the $K^+K^-$ channel.

$J/\psi \rightarrow \phi K^+K^-$: The $K^+K^-$ mass distribution recoiling against a $\phi$ is shown in Fig. 2 c. A clear enhancement can be seen at the nominal $f'(1520)$ mass, however, the peak also shows a statistically significant high mass shoulder. When parametrized with two non-interfering Breit-Wigner amplitudes, the fit assigns a mass of $1671_{-17}^{+15} \pm 10$ MeV and a width of $126_{-40}^{+60} \pm 15$ MeV to the high mass structure. The fit yields the following branching ratios:

$$B(J/\psi \rightarrow \phi f') \cdot B(f' \rightarrow K\bar{K}) = (6.4 \pm 0.6 \pm 1.6) \times 10^{-4}$$

and

$$B(J/\psi \rightarrow \phi \theta(1671)) \cdot B(\theta \rightarrow K\bar{K}) = (3.4_{-0.8}^{+1.0} \pm 0.9) \times 10^{-4}.$$  

The fitted mass of the higher mass structure is lower than expected for the $\theta(1720)$. Conversely, a fit allowing for coherent production of $f'$ and $\theta$ can accommodate a standard $\theta$ with good probability. Coherent and incoherent fit give consistent results for the $f'$, while the branching ratio for the higher mass structure turns out to be $\approx 2.5$ smaller if interference is assumed.

**Conclusions on the $K\bar{K}$ Data:** Structures consistent with the $\theta$ are observed in the $K\bar{K}$ spectra from the $J/\psi \rightarrow \omega K\bar{K}$ and $J/\psi \rightarrow \phi K^+K^-$ reactions. In the following discussion
we will assume that it is indeed the $\theta$ which shows up in the three channels presented in Fig. 2, and that the production of "pure" gluonic matter is suppressed in hadronic $J/\psi$ decays. Assuming a correlation of the $\omega$, $\phi$ quark contents with those of the recoiling system, one must then conclude that the $\theta$ contains a sizeable amount of $u$, $d$, and $s$ quarks. This does not rule out the "glueball" interpretation of the $\theta$, as mixing of bound gluonic states with $q\bar{q}$ states is expected. The appearance of the $\theta$ in reactions together with a $\phi$ and an $\omega$ means that it is not ideally mixed like the other tensor mesons, and suggests that the $\theta$ is not an ordinary $q\bar{q}$ state.

Study of the $K\bar{K}\pi$ Final State

In this study we are mainly interested in the study of the $E/\iota$ mass region. The $\iota(1460)0^{-+}$ state, which is produced with the largest branching ratio in radiative $J/\psi$ decays, is a prime "glueball" candidate (see Fig. 3 a). The $E(1420)$, has long been known from hadronic interactions. First seen in $p\bar{p}$ annihilations at rest, its spin parity was determined as $0^-$. Later measurements, using other reactions, indicated $J^P = 1^+$, while the latest measurement of Chung et al. again points to spin 0. This result has led to the speculation that $E$ and $\iota$ are one and the same object. With the large data set of $J/\psi$ decays available we have the chance to address the $E/\iota$ question within one experiment.

$J/\psi \rightarrow \omega K\bar{K}\pi$: This reaction was studied using the $\omega K^\pm K_{s}\pi^\mp$ and $\omega K^+K^-\pi^0$ final states. Both $K\bar{K}\pi$ mass spectra show clear evidence for a state in the $E/\iota$ mass range with consistent resonance parameters and branching ratios. From the combined set of both channels we obtained the mass spectrum displayed in Fig. 3 b and the resonance parameters listed in Table 1. Due to background, the study of the $K\bar{K}\pi$ substructure is difficult. Still, a comparison with the $\iota$ reveals differences, i.e. a less pronounced peaking of the $K\bar{K}$ submass at $K\bar{K}$ threshold. The data suggest non zero spin.

$J/\psi \rightarrow \phi K\bar{K}\pi$: The reaction $J/\psi \rightarrow \phi K^\pm K_{s}\pi^\mp$ was studied using the $K^+ K^-$ and $K_sK_l$ decay modes of the $\phi$. Figure 3 c shows the added $K^\pm K_{s}\pi^\mp$ mass spectra, with no indication for the production of a resonance in the $E/\iota$ mass range. The resulting upper limit is listed in Table 1.

Conclusions on the $K\bar{K}\pi$ Data: To summarize, a resonance at 1444 MeV with $24 < \Gamma < 84$ MeV at 90%
C.L. was found in the $J/\psi \to \omega K\bar{K}\pi$ decay. The width is hardly compatible with that of the $\iota$. The next best guess is that the structure is due to the $E$ meson. The absence of the $J/\psi \to \phi E$ decay indicates that the $E$ cannot be the $ss$ state in an ideally mixed $1^{++}$ nonet.

**Study of the $\eta\pi\pi$ Final State**

In this section we are mainly concerned with the study the $D$ and $E$ mesons. In addition, $\eta\pi\pi$ is a favorable final state to look for the radial excitations of the $\eta$ and $\eta'$.

$J/\psi \to \gamma\eta\pi^+\pi^-$: This decay was studied in two decay modes of the $\eta$, $\eta \to \gamma\gamma$ and $\eta \to \pi^+\pi^-\pi^0$. Figure 4 a shows the $\eta\pi^+\pi^-$ mass spectrum after adding both channels and demanding that at least one $\eta\pi^\pm$ mass combination is consistent with the nominal $\delta(980)$ mass within 50 MeV. The spectrum looks very complicated, with several resonances contributing to it. Two states around 1285 MeV and 1390 MeV, however, can be isolated. The first structure is consistent with the $D(1283)$ parameters. No signal is seen at the nominal $\iota$ mass. The broad structures at higher mass are interesting, but an interpretation has to wait until a spin parity analysis has been completed.

$J/\psi \to \omega\eta\pi^+\pi^-$: This reaction was observed in the $4\pi^\pm 4\gamma$ final state. The $\eta'$ and two states consistent with $D(1283)$ and $E(1420)$ appear in the $\eta\pi^+\pi^-$ mass spectrum recoiling against an $\omega$. The values for mass, width and $J/\psi$ branching ratios can be found in Table 1. A study of the $\eta\pi^+\pi^-$ system reveals that the resonances at 1283 MeV and 1421 MeV are correlated with a $\delta(980)$ in the $\eta\pi^\pm$ subsystem. Therefore the background is much reduced, if one demands that at least one $\eta\pi^\pm$ submass is consistent with the nominal $\delta$ mass within 50 MeV (Fig. 4 b).

$J/\psi \to \phi\eta\pi^+\pi^-$: This reaction was studied in the $K^+K^-\pi^+\pi^-\gamma\gamma$ final state. The $\eta\pi^+\pi^-$ mass distribution recoiling against the $\phi$ shows clear evidence for $J/\psi \to \phi\eta'$ and a narrow structure in the 1280 MeV mass region (see Fig. 4 c). The enhancement appears to be correlated with a $\delta(980)$ in the $\eta\pi^\pm$ subsystem.
Conclusions on the $\eta\pi\pi$ Data: In all three channels, a narrow structure is evident in the 1280 MeV mass range. The natural candidate is the $D(1283)$, a well established resonance which is classified as the $u\bar{u} + d\bar{d}$ state in an ideally mixed axial vector nonet. The observation of the "$D$" recoiling against a $\phi$ with a third of the rate of $J/\psi \rightarrow \omega "D"$, however, is surprising and contrary to the assumption that the quark contents of the two final state particles are correlated. Therefore either this premise is wrong, the $1^{++}$ nonet is not ideally mixed, or what we observe is not the $D$. Indeed, there is another candidate, the radial excitation of the $\eta$, the $\eta(1275)$. A determination of the spin-parity of the observed structures will be helpful to clear up this confusing picture.

The second interesting feature in the radiative decay, the peak at 1395 MeV, does not have its equivalent in the hadronic channels. A peak with similar width, but a mass $\approx 25$ MeV higher, appears in the $\omega$ recoil mass spectrum. Whether the mass shift can be explained by interference still has to be investigated. Similar to the $K\bar{K}\pi$ final state, there is little evidence for this structure in the $\eta\pi^+\pi^-$ mass spectrum recoiling against the $\phi$. Consequently, if the structures are due to the $E$ meson, it cannot be a pure $s\bar{s}$ state, and $B("E" \rightarrow \eta\pi\pi)/B("E" \rightarrow K\bar{K}\pi) \approx 1.3$.

Interestingly enough, the broad structures observed in $J/\psi \rightarrow \gamma\eta\pi^+\pi^-$ between 1.5 and 2.0 GeV are not obvious in the hadronic channels, making them candidates for something "new".

**Table 1.**

Resonances in the $\eta\pi^+\pi^-$ and $K\bar{K}\pi$ systems; the product branching ratios $B(J/\psi \rightarrow \{\gamma, \phi, \omega\}R) \cdot B(R \rightarrow \{K\bar{K}\pi, \eta\pi\pi\})$ are corrected for unobserved decay modes.

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<td>$\iota$</td>
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**References**

Recent Results on Two-Photon Processes from DESY

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Abstract

In this talk some recent results on two-photon scattering processes from DESY are presented, including results on the production of the charmonium state $\eta_c(2980)$, on the exclusive production of proton-antiproton pairs, on jet formation by two quasi-real photons, and on $\Lambda_{MS}$ determinations from the structure function of the photon.
1 Introduction

Scattering of two photons can be observed at e⁺e⁻ storage rings in the reaction (see diagram):

\[ e^+e^- \rightarrow e^+e^-X \]

The photons producing the final state X are predominantly quasi-real with the scattered electrons going down the beam pipe. However, there is also a tail of large Q² photons which can be tagged by detecting the scattered electrons at some finite angle.

The analyses of the two-photon production of the ηₑ, of pp pairs, and of high pₑ jets as described below have been done in the 'notag' mode, i.e. with quasi-real photons. A typical application of tagging is the study of the photon structure functions where one of the photons has a high Q² probing the structure of the usually quasi-real 'target photon'. Most recent results on the photon structure function F₀ will be presented.

Results on two-photon processes from the Crystal Ball experiment will be reported in another contribution to this meeting [1].

2 Observation of the Charmonium State ηₑ

The PLUTO collaboration for the first time measured the two-photon production of the charmonium state ηₑ [2]. In a data sample corresponding to an integrated luminosity of 45 pb⁻¹ they searched for the reaction

\[ γγ \rightarrow ηₑ \rightarrow K_S^0K^±π^±. \] (1)

The K_S^0's were observed via their π⁺π⁻ decay mode and identified by requiring that the π⁺π⁻ came from a secondary vertex separated from the primary e⁺e⁻ collision point. The charged kaons and pions in reaction (1) were not identified. Therefore, both the K⁺π⁻ and the K⁻π⁺ assignments were tried leading to two entries per event in the KKπ invariant mass plot of Fig.1. In the ηₑ region, the difference in the KKπ invariant mass of these two combinations is smaller than the mass resolution of about 100 MeV.

The PLUTO group found 7 events in the ηₑ region distributed as expected from the detector resolution. The background in the ηₑ region is assumed to be negligible. From the observed 7 events the PLUTO group obtained for the product of the γγ width times the branching ratio into K_S^0K^±π±:

\[ Γ(ηₑ \rightarrow γγ) \cdot B(ηₑ \rightarrow K_S^0K^±π^±) = (0.5 \pm 0.2 \pm 0.15 \pm 0.1) keV \]

The TASSO collaboration searched for the two-photon production of the ηₑ in the same decay mode and with similar analysis methods. In particular the K_S^0's were also identified by requiring a secondary vertex for the decay pions. The vertex finding procedure is still being improved and thus the following results are not yet final.

The K_S^0K^±π± invariant mass spectrum obtained by TASSO is shown in Fig.2. The mass resolution in the ηₑ region, σ ≈ 60 MeV, is better than in the PLUTO experiment. However, the region around the ηₑ is not free of background. From a fit to the mass spectrum 6.6 ± 3.3 ηₑ events above background are found yielding the preliminary result:

\[ Γ(ηₑ \rightarrow γγ) \cdot B(ηₑ \rightarrow K_S^0K^±π^±) = (1.2 \pm 0.6 \pm 0.4) keV \]
Within the errors this result is in agreement with the PLUTO measurement.

Using isospin invariance the $K^0 K^+\pi^+$ branching ratio is 1/3 of the $K\bar{K}\pi$ branching ratio which has been measured by the Mark III group to be [3]:

$$B(\eta_c \rightarrow K\bar{K}\pi) = (6.1 \pm 2.2)\%.$$  \hfill (2)

This value has been obtained using the branching ratio $B(J/\psi \rightarrow \gamma\eta_c) = (1.27 \pm 0.36)\%$ measured with the Crystal Ball [4].

With the Mark III value for $B(\eta_c \rightarrow K\bar{K}\pi)$ one obtains for the $\gamma\gamma$ width of the $\eta_c$ from the two experiments:

$$\Gamma(\eta_c \rightarrow \gamma\gamma) = (25 \pm 14)\text{ keV} \quad (\text{PLUTO})$$
$$\Gamma(\eta_c \rightarrow \gamma\gamma) = (59 \pm 41)\text{ keV} \quad (\text{TASSO})$$

An ISR experiment searching for the reaction $p\bar{p} \rightarrow \eta_c \rightarrow \gamma\gamma$, obtained the upper limit [5]:

$$\Gamma(\eta_c \rightarrow \gamma\gamma) < 7\text{ keV} \quad (95\% \text{ c.l.}).$$

A theoretical estimate can be obtained by relating the $\gamma\gamma$ width of the $\eta_c$ to the leptonic width of the $J/\psi$. Assuming equal wave functions for the spin singlet and triplet charmonium states yields

$$\Gamma(\eta_c \rightarrow \gamma\gamma) = 3 e^2 \cdot \Gamma(J/\psi \rightarrow e^+e^-)$$

With the measured leptonic width of the $J/\psi$ this gives $\Gamma(\eta_c \rightarrow \gamma\gamma) \approx 6\text{ keV}$. QCD sum rule calculations indicate that this value could be even smaller [6]. Thus the $\eta_c$ rates observed by the PLUTO and TASSO groups seem to be rather large. For example, using the ISR upper limit and the branching ratio (2) at face value TASSO should see less than 0.8 events of reaction (1).

### 3 Two-Photon Production of Proton-Antiproton Pairs

One of the most complicated processes which have been calculated by perturbative QCD methods is the exclusive production of hadron pairs by two photons [7]. Measurements of charged pion and kaon pair production (pions and kaons not separated) by the PLUTO and Mark II groups are in good agreement with these calculations for $\gamma\gamma$ invariant masses,
W_{\gamma\gamma} above about 2 GeV [8]. The TASSO group [9] found that the cross section for proton-antiproton pair production,

$$\gamma\gamma \rightarrow p\bar{p},$$

in a similar W_{\gamma\gamma} range (2 to 3 GeV) is much larger than predicted by the QCD calculations of Farrar et al. [10].

The TASSO measurement has now been confirmed by the JADE group which analysed p\bar{p} production in a W_{\gamma\gamma} range from 2.0 to 2.6 GeV [11]. The protons and antiprotons have been identified by the energy loss in the JADE jet chamber. Good separation of protons from pions and kaons is possible for momenta up to 1 GeV. In a data sample corresponding to an integrated luminosity of about 84 pb^{-1} 41 p\bar{p} events have been found. The cross section of reaction (3) in the c.m.s. angular range |\cos \Theta'| < 0.6 is shown as a function of W_{\gamma\gamma} in Fig.3. The data are compared to a QCD curve [10] which lies much below the data. The JADE group emphasizes an apparent peaking of the cross section around 2.25 GeV as indicative for a possible resonance formation. However, the TASSO measurement, while statistically in agreement with the JADE data, does not exhibit this shape.

The differential cross section obtained by JADE averaging over the whole W_{\gamma\gamma} range is shown in Fig.4. The cross section seems to be enhanced around 90°. By comparing the region below and above |\cos \Theta'| = 0.4 the JADE group finds that the \chi^2 probability for the angular distribution being isotropic is less than 0.1%. The TASSO group also observed a slight preference for large angles, but their result is much less significant, partly because they cover a smaller angular range. In a preliminary study the TPC/\gamma\gamma collaboration found in a similar W_{\gamma\gamma} region also some enhancement at large angles while for larger W_{\gamma\gamma} they observed rather a peaking in the forward direction [8].

An enhancement in the angular distribution around 90° is suggestive for resonance production as it can hardly be explained by another production mechanism. E.g., the QED Born term for the coupling of two photons to proton-antiproton pairs yields a flat angular distribution near threshold and develops a forward peak with increasing energy. This is also expected for any non-resonant hadronic production mechanism and, in particular, is also predicted by the QCD calculations.

The angular distribution observed by JADE can be reproduced by J=2 states. For the lowest orbital angular momenta in the \gamma\gamma initial state one can assume that the \gamma\gamma helicities are dominantly \lambda=\pm 2 for J^P = 2^+ states and \lambda=0 for J^P = 2^- states. The corresponding
angular distributions are shown as the dotted \((J^P = 2^+, \lambda = 2)\) and the dashed-dotted \((J^P = 2^-, \lambda = 0)\) curves in Fig.4. For \(J^P = 2^-\) the \(p\bar{p}\) state can occur in spin states \(S=0\) and \(S=1\). The shown curve is for \(S=0; S=1\) is not in agreement with the data.

In summary we can conclude that around 2.5 GeV \(p\bar{p}\) production is definitely not yet in the continuum region described by perturbative QCD. Possible resonance formation can be tested analysing the angular distribution. However, this task requires better statistics than available today.

4 Jet Production by Two Quasi-Real Photons

In the Born approximation hadron production by two photons proceeds via quark pair production according to the diagram below. The coupling of the photons to quarks is in this case purely electromagnetic and the ratio of hadron production to muon pair production, \(R_{\gamma\gamma}\), is given by the charges of the quarks with different flavour \(f\) and colour \(c\):

\[
R_{\gamma\gamma} = \frac{1}{3} \sum_f \left( \sum_c e_f^2 \right)^2.
\]

The standard fractionally charged quark (FCQ) model and the Han-Nambu model of integrally charged quarks yield \(R_{\gamma\gamma}\) values which differ by about a factor of 3 (\(R_{\gamma\gamma}\) is 34/27 for the FCQ and 10/3 for the ICQ model including \(u,d,s,c\) quarks). In a gauge invariant formulation of the ICQ model \(R_{\gamma\gamma}\) can be even larger due to additional couplings to charged gluons [12]. In this model with broken SU(3)\(_{\text{color}}\) symmetry, gluons acquire masses resulting in a form factor suppression of all non-color-singlet contributions to \(R_{\gamma\gamma}\). Thus the difference between the FCQ and ICQ models will be damped for virtual photons by a factor proportional to

\[
\frac{M_{\text{pluon}}^2}{M_{\text{pluon}}^2 + Q^2}
\]

Hence scattering of two quasi-real photons is best suited to distinguish between these models.

In contrast to one-photon annihilation processes the two-photon Born diagram describes hadron production only for large momentum transfers along the virtual quark line corresponding to large transverse momenta of the quarks and the fragmentation products. Recent calculations indicate that the QCD corrections to the Born term (K-factor) are small [13], [14]. Furthermore, also the contributions of higher order QCD processes are predicted to be small for transverse momenta of hadrons above about 1 to 2 GeV. The more surprising was the observation by the TASSO group that the yield of high-\(p_t\) hadrons produced by quasi-real photons (no-tag experiment) was about 3 to 4 times larger than expected from the Born diagram [15]. An excess of two-jet events with \(p_t^{'\text{re}}\) values above about 1.5 GeV was also observed by the PLUTO group [16]. This excess was found to decrease with increasing \(Q^2\) of one of the photons so that at \(Q^2\) above about 12 GeV\(^2\) consistency with the Born term was reached.

The observed \(Q^2\) and \(p_t\) dependence can be explained by the gauged ICQ model. However, before accepting this non-standard model of quarks and gluons, more experimental tests can be made. The PLUTO group recently investigated the topology of events produced by two quasi-real photons (no-tag experiment) [17]. Using a thrust algorithm they divided each event into two jets. In Fig.5 the ratio of the observed number of jets to the one expected
from the Born term, $\tilde{R}_{\gamma\gamma}$, is plotted versus the jet transverse momentum, $p_t\text{jet}$. The large excess at small $p_t\text{jet}$ is explained by a VMD model which is, however, unable to describe the excess persisting above about 2 GeV where the Born term is expected to dominate. An inspection of the thrust distribution shows that the average thrust of large-$p_t$ jets is smaller than expected from the Born term. Including only events with a thrust larger than 0.9 (Fig.6) yields $\tilde{R}_{\gamma\gamma} = 1.2 \pm 0.3$ for $2 < p_t\text{jet} < 6$ GeV. Thus the 'naive' Han-Nambu model of integer quark charges ($\tilde{R}_{\gamma\gamma} = 2.65$) is excluded by more than 4 standard deviations. Using the same data the PLUTO group had already published an upper limit for the gluon mass [18]:

$$M_{\text{gluon}} < 5 \text{ MeV} \quad (95\% \text{ c.l.}).$$

As reported at this meeting, experimental results on high energy Compton scattering and on the production of two prompt photons in pion-nucleon scattering also strongly support the model of fractionally charged quarks.

Thus the ICQ model is not likely to be the right explanation for the excess of large-$p_t$ jets (or hadrons) which, according to the PLUTO finding, are produced in less jetty events. Aurenche et al. [14] suggest that a modified VMD contribution, as derived from photoproduction data, could make a significant contribution even at large $p_t$. The $p_t$ distribution of inclusive hadrons measured by TASSO may just be explainable by adding such a VMD estimate on top of the Born term and the QCD corrections [19].

The PLUTO group finds that the event topology and the $p_t$ distributions can be qualitatively described by adding contributions from a QCD process which leads to four-jet final states (two high-$p_t$ jets and two beam pipe jets). It has to be studied further if a quantitative description can be achieved by summing all leading QCD diagrams.

## 5 The Structure Function of the Photon

The hadronic structure of a photon can be probed in collisions of a high-$Q^2$ photon with a quasi-real 'target photon'. In such an experiment the virtual photon is tagged by detecting the scattered electron at a large angle. The $\gamma\gamma$ invariant mass, $W_{\gamma\gamma}$, is determined from the hadronic system observed in the detector.

The structure function $F_2^\gamma(x,Q^2)$ of the photon has been calculated perturbatively in leading and next-to-leading order QCD [20]. Imposing the so-called 'asymptotic boundary
condition, i.e., requiring $F_2^\gamma$ to approach the Born term for $\gamma\gamma \rightarrow q\bar{q} \rightarrow \text{hadrons}$ at large $Q^2$, allows to calculate $F_2^\gamma$ absolutely. The pointlike, perturbatively calculable part of the structure function yields the $Q^2$ dependence

$$F_2^\gamma \sim \ln \frac{Q^2}{\Lambda^2}$$

with $\Lambda$ being the QCD scale parameter. The hadronic component of the photon cannot be calculated perturbatively and is usually estimated by the VMD model. Several groups at DESY determined $\Lambda$ employing relation (4) as obtained by higher order QCD calculations.

It should be mentioned that the validity of this method of determining $\Lambda$ is theoretically controversial [21]. The problems arise from the separation of the hadronic and the pointlike piece of the photon. Using calculations in which these pieces are not separated the experimental sensitivity to $\Lambda$ is lost.

The latest results on $F_2^\gamma(x, Q^2)$ have been obtained by the TASSO [22] and PLUTO [19] groups. The TASSO group determined $F_2^\gamma(x, Q^2)$ at an average $Q^2$ of 23 GeV$^2$. Figure 7 shows the measured structure function with the charm contribution (estimated by the Born term) subtracted. The charm contribution is assumed to be not sensitive to $\Lambda$ since in this case the relevant scale may be set by the charm quark mass rather than $\Lambda$. The subtracted structure function is compared to the sum of a higher order QCD calculation with $\Lambda_{\overline{MS}} = 150$ MeV and a hadronic piece estimated by VMD. Above about $x=0.3$ the VMD contribution is small and the structure function is sensitive to $\Lambda$. The measured structure function is equally well described if the pointlike QCD part is replaced by the quark-parton model prediction using constituent quark masses (pure QED couplings).

The PLUTO group determined the structure function $F_2^\gamma(x, Q^2)$ in a wide range of $Q^2$ values [19]. Figure 8 shows the measured $Q^2$ dependence averaged over the $x$ range from 0.3 to 0.8. Also shown are data points from TASSO [22] and JADE [23]. The measurements are compared to higher order QCD calculations for $\Lambda_{\overline{MS}} = 183$ MeV. Note that the $\Lambda$ value is determined by the absolute height rather than the slope of the curve.

Taking the average of all $\Lambda_{\overline{MS}}$ determinations from $F_2^\gamma$ by DESY groups yields (statistical and systematic errors have been combined in quadrature) [19]:

$$\Lambda_{\overline{MS}} = (193 \pm 43) \text{ MeV}.$$
References

[5] R704 Coll., C. Baglin et al., paper contributed to the Kyoto Conf. (see [19])
[17] PLUTO Coll., Ch. Berger et al., to be published
A MEASUREMENT OF $\Gamma_{\gamma\gamma}(\eta')$ USING THE CRYSTAL BALL DETECTOR

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ABSTRACT

The reaction $e^+e^-\rightarrow e^+e^-\eta'$ has been observed using the Crystal Ball detector at the DORIS II storage ring at DESY. The $\eta'$ has been detected via the decay chain $\eta'\rightarrow \eta\pi^0\pi^0\rightarrow 6\gamma$. The radiative width $\Gamma_{\gamma\gamma}(\eta')$ has been determined to be $4.10 \pm 0.31 \pm 0.80$ keV. A search has been made for other states decaying into $\eta\pi^0\pi^0$.

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The study of the partial widths of mesons into two photons, $\Gamma_{YY}$, can be used to provide information on the quark content of these mesons. I report here on a new measurement of $\Gamma_{YY}$ for the $\eta'$ meson. The data were taken with the Crystal Ball detector at the DORIS II storage ring at DESY. The reaction studied was $\gamma\gamma \rightarrow \eta' \eta \pi^0 \pi^0 \gamma$. The Crystal Ball detector with its spherical array of NaI(Tl) shower counters is well suited to measure such a reaction.

Analysis criteria were applied to a data sample of 154 pb$^{-1}$ to select events with six photons which formed an $\eta$ and two $\pi^0$'s. The mass of the system was calculated, and is shown in Figure 1. We have 185.3 ± 13.9 events at the $\eta'$ mass, with essentially no background. This corresponds to a partial width $\Gamma_{YY}(\eta') = 4.10 \pm 0.31 \pm 0.80$ keV (preliminary). This value agrees well with the world average$^1$ $(4.3 \pm 0.30$ keV).

From the measurements of $\Gamma_{YY}$ one can deduce the SU(3) mixing angle for the pseudoscalar nonet with the following formulae:

\[
\frac{\Gamma_{YY}(\eta)}{m^3_\eta} = \frac{\Gamma_{YY}(\pi^0)}{m^3_{\pi^0}} \left( \frac{1}{\sqrt{3}} \cos \theta_p - r_p \sin \theta_p \right)^2
\]

\[
\frac{\Gamma_{YY}(\eta')}{m^3_\eta} = \frac{\Gamma_{YY}(\pi^0)}{m^3_{\pi^0}} \left( \frac{1}{\sqrt{3}} \sin \theta_p + r_p \cos \theta_p \right)^2
\]

$r_p$ is the ratio of the decay constants for the singlet and octet members of the nonet. A value equal to one is the special condition called "nonet symmetry".

Our measurement, plus the world average$^1$ values for $\Gamma_{YY}(\eta)$ and $\Gamma_{YY}(\pi^0)$ yield:

$$\theta_p = -20.0^\circ \pm 2.0^\circ$$

$$r_p = 0.94 \pm 0.06.$$ 

The standard Gell-Mann-Okubo mass formula gives a value for $\theta_p$ which is a factor of two smaller (-10$^\circ$). However, recent calculations$^2$ of first-order corrections have shown that although the corrections to the mixing formulae in Equation 1 are small, the corrections to the mass formula can be large, and a consistent picture is possible with $\theta_p \approx 20^\circ$.

We have also searched for higher mass states decaying into $\eta \pi^0 \pi^0$. An example of such a state would be a radially excited $\eta$ or $\eta'$, which are expected$^3,4$ to be in the 1100-2000 MeV/c$^2$ mass range. Furthermore, they should have large branching ratios into $\eta \pi \pi \pi^0$.

As can be seen in Figure 1, no structure is seen, and we have set 90\%
confidence level upper limits for $\Gamma_{\gamma\gamma}(X)B(X\to\eta\pi\pi) < 0.3$ keV (preliminary) for masses less than 1500 MeV/c$^2$.

This upper limit is especially interesting for the $\eta(1275)^5,6)$. If this meson is a radially excited $\eta$, its two photon width is expected to be 2-3 keV$^7,8)$. Thus, unless the branching ratio into $\eta\pi\pi$ is small, it is difficult to interpret this state as a radially excited pseudoscalar.

I would like to thank J. Tran Thanh Van and the organizing committee for their excellent organization of this conference, and for the wisdom to locate it in such pleasant surroundings.

REFERENCES

1. H. Kolanoski, invited talk at the 1985 International Symposium on Lepton Photon Interactions at High Energies, Kyoto, Japan, Aug 19-24, 1985, BONN-HE-85-34. The world average value for $\Gamma_{\gamma\gamma}(\eta)$ is $0.56 \pm 0.04$ keV, where only the results from $e^+e^-$ have been used. The value for the $\pi^0$ is $7.33 \pm 0.20$ eV.


8. M. Peskin, private communication.

FIGURE CAPTION

1. The mass distribution of the $\eta\pi^0\pi^0$ system.
FIG 1
EVIDENCE FOR AN EXTRA I=0 SCALAR NEAR 1 GEV

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presented by D. Morgan

ABSTRACT

Incorporating new data on pp \rightarrow pp\pi\pi(K\bar{K}) a coupled channel analysis of results on I = 0 S wave \pi\pi, K\bar{K} final state interactions is performed extending up to 1.6 GeV. The resonance content of our solution, remarkably stable between fits, comprises 4 states, 2 more than required by the naive quark model. In detail, we find 2 wide resonances \epsilon(1000) and \epsilon'(1420) and 2 narrow objects S_1(996) and S_2(988) which conspire to reproduce the S^* effect. S_1, which couples like an SU(3) singlet, forms a natural glueball candidate. Other reactions like \gamma\gamma \rightarrow \pi\pi (K\bar{K}) and heavy flavour decays are amenable to similar analysis but as yet lack the precision to probe fine details.
What can we learn about the meson spectrum from double Pomeron exchange (D P E) reactions and from similar production and decay processes like γγ + mesons and heavy flavour decay? Such 'production' data is a valuable adjunct to conventional meson-meson scattering in the quest for glueballs and other novelties suggested by QCD. According to various model calculations, these additions to the (qq̅) spectrum of the naive quark model should occur from about 1 GeV and the search for candidates has come to be a central objective of spectroscopy. Among glueballs, the lightest specimens should be a scalar \(\pi_0(0^{++})\), a pseudo-scalar \(\eta(0^{-+})\) and a tensor \(f_2(2^{++})\). Highly plausible candidates already exist for the latter two in the shape of \(\eta(1460)\) and \(\phi(1690)\); this intensifies the interest in establishing an \(\pi_0\) signal in the appropriate channels, \(\pi\pi, K\bar{K}\) and perhaps \(\eta\eta\). Despite the likelihood of mixing, we can still aim to establish the existence of extra \(I=0\) states and study the associated decay patterns. Such information is obviously not sufficient to establish a new glueball (for instance, we need to demonstrate the absence of \(I\neq0\) companions) but it is a necessary start. To this end, we shall present evidence for a new \(I=0\) scalar near 1 GeV.

The new data that we study comes from the D P E reaction illustrated in the figure with the centrally produced system a \(\pi^+\pi^-\) or \(K^+K^-\) pair. This type of reaction has long been commended as a good way to make glueballs from the supposed association of the Pomeron with multi-gluon exchange in QCD. Quantum numbers available are \(J^{PC}= 0^{++}, 2^{++}, \ldots\), all for \(I=0\), so that \(\pi_0\) and \(f_2\) should be excited. A number of experiments have been performed, in particular a high statistics study (3.10^6 events) of pp + pp \(\pi^+\pi^- (K^+K^-)\) by the AFS collaboration at the ISR with \(\sqrt{s} = 63\) GeV. Very small deflections and energy losses are required for the through-going protons; as a result, a seemingly very clear \(P\, P\) signal is extracted. The AFS group have decomposed the data into angular components so that we have available the cross-section for S-wave production.

The above 'special gluon mechanism' strategy does not obviously confer any advantage where only a few channels are open (\(\pi\pi, K\bar{K}\)), since unitarity then ties the amplitudes for \(P\, P + \pi\pi\) (\(KK\)) rather closely to the corresponding 'hadronic' amplitudes \(T(\pi\pi + \pi\pi)\) etc. Although at first sight discouraging, this linkage, can be exploited to make a joint fit of the new \(P\, P + \pi\pi\) (\(K^+K^-\)) data with its excellent precision along with classic \(\pi\pi + \pi\pi(K\bar{K})\)
information $^7$,$^8$). This is our approach.

The ingredients and steps of our fit are:

(i) The AFS collaboration data from ref (4). Note, we are in a region of
dimeson mass where $\pi\pi$ (henceforth, channel 1) and $K\bar{K}$ (channel 2) dominate.
(ii) By removing phase-space and effective luminosity factors (well-described
by standard Regge systematics - see refs (3) and (5)), we can convert the
published spectra into effective $P\bar{P} + \pi\pi$ ($K\bar{K}$) S-wave cross-sections (cf.
the conventional discussion of $\gamma\gamma$ processes). We can introduce associated
amplitudes

$$F_1 \equiv F(P\bar{P} + \pi\pi), F_2 \equiv F(P\bar{P} + K\bar{K}).$$ (1)

(iii) By unitarity, the above production amplitudes are related to the ordinary
hadronic scattering amplitudes $T_{ij}$ (for $\pi\pi + \pi\pi$, $K\bar{K}$ etc.) by the formulæ

$$F_1 = \alpha_1 T_{11} + \alpha_2 T_{21}, F_2 = \alpha_1 T_{12} + \alpha_2 T_{22}$$ (2)

with $\alpha_i$ smooth, real functions of energy. We may think of them as 'intrinsic'
production amplitudes characteristic of the process in question - here $P\bar{P} + x$, at other times $\gamma\gamma + x, \psi' + \psi x$ etc.
(iv) The final ingredient is classic information on $\pi\pi + \pi\pi$ ($T_{11}$)$^7$ and $\pi\pi + K\bar{K}$ ($T_{12}$)$^8$.

Our procedure is to parametrize $T_{ij}$ via the associated $K$ matrix (or its
inverse $K^{-1}$), allowing both polynomial and pole contributions and to assign
quadratic forms to the $\alpha_i(s)$- a maximum of $28$ parameters. The quality of the
resulting fits, unusually good ($\chi^2/\text{NDF} = 1.1$) for a relatively heterogeneous
data set, is illustrated in Figs 1 and 2.

Once having fitted the parameters in the formula for $K$ or $K^{-1}$, it is a
mechanical task to locate the poles of the corresponding $T$ - matrix. One has
to be careful to specify on which sheet of the energy plane the pole in
question is located (see figure).

For all our many solutions, $T$ has 7
complex poles, 4 denoted ($A,B,D,G$) on sheet II
and 3 labelled ($C,E,F$) on sheet III. Their
positions on the energy plane are shown in Fig.
3 ( sheet II poles ; sheet III poles $x$). Note
how most of the sheet II and sheet III poles
pair up and associate to resonances (c,f.
disc uss ion in refs (2,3)). Locations are remarkable stable from one solution to another, regardless of the parametrization used. These 7 poles are our major new results.

As seen from Fig 3, our set of poles comprises a triplet (A,B,C) very close to $K\bar{K}$ threshold which conspire to produce the $S^*$ effect and 2 associated pairs (D,E) and (G,F) controlling the long range variations with energy. The whole system corresponds to 4 resonances $S_1(996), S_2(988), \epsilon(1000)$ and $\epsilon'(1420)$, as illustrated in Fig. 3. Resonance parameters may be extracted according to a simple recipe explained in ref (3) and are listed in Table 1:

<table>
<thead>
<tr>
<th>Resonance</th>
<th>Poles</th>
<th>$E_R$ (GeV)</th>
<th>$g_\pi$ (GeV)</th>
<th>$g_K$ (GeV)</th>
<th>$g_\pi/g_K$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$S_1(996)$</td>
<td>A,C</td>
<td>$0.996 - 0.024i$</td>
<td>0.25</td>
<td>0.32</td>
<td>0.8</td>
</tr>
<tr>
<td>$S_2(988)$</td>
<td>B</td>
<td>0.988</td>
<td>0.01</td>
<td>0.28</td>
<td>0.04</td>
</tr>
<tr>
<td>$\epsilon(1000)$</td>
<td>D,E</td>
<td>$1.00 - 0.40i$</td>
<td>0.53</td>
<td>0.33</td>
<td>1.6</td>
</tr>
<tr>
<td>$\epsilon'(1420)$</td>
<td>G,F</td>
<td>$1.42 - 0.22i$</td>
<td>0.7</td>
<td>0.2</td>
<td>3.5</td>
</tr>
</tbody>
</table>

Table 1: I = 0 S-wave resonances below 1.0 GeV from our fits

Relative to the PDG Tables 9) and the naive quark model (NQM), we have 2 extra I = 0 scalars. Where PDG lists $S(975)$ and $\epsilon(1300)$, our spectrum resolves $\epsilon(1300)$ into $\epsilon(1000)$ and $\epsilon'(1420)$, modifies $S(975)$ to $S_1(996)$ and inserts a brand new $K\bar{K}$ bound state, $S_2(988)$10). The latter which is the special new feature of our solution is overwhelmingly favoured by current experimental information 3).

What flavour (more generally parton) content should be ascribed to our new spectrum? Mass splittings are a poor guide where such broad states feature. We are therefore left with the ratio of decay amplitudes to $\pi\pi$ and $K\bar{K}$ as the sole criterion. Comparing various idealized possibilities with the $g_\pi/g_K$ ratios of Table 1 shows $\epsilon(1000)$ and $\epsilon'(1420)$ to be consistent with a ($uu + dd$) composition, $S_2$ with an (ss) identification (albeit with a reduced coupling) and $S_1$ with an SU(3) singlet description. A possible overall assignment would therefore be: for $\epsilon$ and $S_2$ to form the regular ground-state (qq) compounds of the NQM, $\epsilon'$ a radial excitation (to go along with Erkin et al's $S'(1770)$8)11), if this is confirmed) and $S_1$ a glueball.

The lynch-pin of our scenario is the new $K\bar{K}$ bound state B ($S_2(988)$); securing further direct evidence for it is a top priority. One very characteristic consequence of B is the occurrence of a sharp downwards hook in the phase of $\pi\pi + K\bar{K}$ (arg $T$ ) from threshold. Hints of this
are seen in the existing data (Fig 1(b)) but need better confirmation. The same goes for the spike immediately above threshold of $\sigma(\pi^{-}\pi^{-} + K^0\bar{K}^0)^{(3/4)}$. In both these cases, it would be useful to have comparable data for $K^+K^-$ and $K^0\bar{K}^0$. The latter channel would be a good place to explore virtual $K\bar{K}$ scattering close to the expected peak associated with $B$ using reactions like $K^-p + A K^0\bar{K}^0$. One additional piece of information already present in the AFS data but not used in our fit is their S-D interference signal. Comparison with our prediction (Fig 2(b)) is very satisfactory as regards general trends; however, the data is not sufficiently fine-grained to probe the detailed structure.

The DPE reactions have been central to our analysis; because of their exceptional precision, they are the only production reactions that we have allowed to influence resonance parameters. However, once the hadronic amplitudes, $T_{ij}$, are determined, other analogous reactions like $\gamma\gamma + \pi\pi$ and various heavy flavour decays become legitimate targets for fitting using formula (2) with new characteristic $\alpha_i$'s as fit parameters. The results already have interesting bearing on the flavour preferences, $(\alpha_1^x: \alpha_2^x)$, that different production mechanisms ($x$) reveal. In all cases, the pressing need is for improved statistics; where the spectrum extends above 1 GeV, partial wave separation is also required.

In conclusion, our central result is that a global fit of existing $\pi\pi$ and $K\bar{K}$ data unambiguously resolves the $S^*$ structure into 3 poles comprising one resonance (a plausible glueball candidate) and a $K\bar{K}$ bound state.

References
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Figure Captions
Fig. 1. Fit to $\pi^+ + \pi^0, K\bar{K}$ data: (a)$\delta_0^0$, (b) $1/4[1-\eta^2]$; (c) arg $T(\pi^+ + K\bar{K})$.
Fig. 2. AFS data on $p + p p^+ \pi^- + \pi^+$ : (a) mass spectrum; (b) S-D interference.
Fig. 3. $T$-matrix poles on E-plane: *(sheet II); x(sheet III); resonance assignments.
Fig. 1.

Fig. 2.
Fig. 3.
Some new results from the GAMS collaboration are presented. The existence and properties of G(1590) are confirmed in measurements made with 100 GeV/c pions. Evidence for a structure around 2.2 GeV/c^2 decaying into ηη' is presented.
1. **INTRODUCTION**

The GAMS collaboration* has collected data on \( \pi^- p \) interactions in two distinct experiments, namely at 38 GeV/c (IHEP) and 100 GeV/c (CERN) incident pion momentum. Both setups are similar and are essentially based on lead-glass calorimeters of the GAMS type [1,2]. Fig. 1a shows the CERN experimental setup (NA12 experiment).

Data have been collected for exclusive reactions leading to neutral final states, i.e.

\[
\pi^- p \rightarrow M^* + n \\
\rightarrow m \gamma
\]

where \( M^* \) decays into photons which are detected in GAMS. Events with a multiplicity of up to 20 photons have been recorded. Fig. 1b shows the lego representation of a typical event with 8 \( \gamma \) produced in the decay \( M^* + \eta \).

The results presented here concern:

i) the \( \eta \) decay channel, the emphasis being on the confirmation at 100 GeV/c of the existence of \( G (1590) \) with \( J^P = 0^+ \), discovered by the GAMS collaboration at IHEP [3];

ii) the possible existence of a structure at 2.2 GeV in the \( \eta \eta' \) decay channel;

iii) the observation of the \( \pi^+ \pi^- \) decay channel in the \( K/\Lambda \) mass region.

2. **NEW EVIDENCE FOR \( G (1590) \rightarrow \eta \eta \) AT 100 GeV/c**

A first selection of \( \eta \eta \) pairs has been made for 4 \( \gamma \) events \( (\eta_1 \rightarrow 2\gamma, \eta_2 \rightarrow 2\gamma) \) and 8 \( \gamma \) events \( (\eta_1 \rightarrow 2\gamma, \eta_2 \rightarrow 3\pi^+ \rightarrow 6\gamma) \) by requiring the reconstructed mass of \( \eta \) candidates to be in the range 509 to 584 MeV/c\(^2\).

A 3C-fit of the selected 4 \( \gamma \) events, including the mass of the recoiling neutron, at a 97% confidence level has produced a sample of 12000 \( \eta \eta \) events (to be compared with \( 1.2 \times 10^6 \pi^+ \pi^- \) and \( 8 \times 10^4 \eta \eta \) identified 4 \( \gamma \) events).

Similarly, a 6C-fit of the selected 8 \( \gamma \) events, taking into account the 3 additional constraints on the mass of the reconstructed \( \pi^+ \pi^- \), has produced a sample of 8000 \( \eta \eta \) events.

The angular distributions of the $\eta$ in the Gottfried-Jackson system of the $\eta\eta$ pairs are shown on fig. 2a (150 MeV mass bins). Angular distributions have been fitted in 50 MeV mass bins with the following expression

$$\frac{d\sigma}{d\Omega} \propto |H_0 + H_-|^2 + |H_+|^2$$

in a maximization of the event by event likelihood function. The only relevant terms in the analysis are

$$H_0 = S P_{00} (\cos\theta) + \sqrt{5} D_0 P_{20} (\cos\theta) + \sqrt{9} G_0 P_{40} (\cos\theta)$$

$$H_- = \sqrt{5/6} D_- P_{21} (\cos\theta) \cos\phi$$

where $S$, $D_0$, $G_0$, $D_-$ are complex amplitudes and $P_{lm} (\cos\theta)$ are Legendre polynomials. $H_+$ as well as $l > 4$ waves have been found to be negligible. The $G$ wave contributes only for $M_{\eta\eta} > 1.7$ GeV/c$^2$.

Fig. 2b shows the result of the fit. Two solutions are found above 1.7 GeV/c$^2$ but no ambiguity exists below 1.7 GeV/c$^2$ where a clear peak shows up in the $S$-wave both in $4\gamma$ and $8\gamma$ events. It corresponds to $G(1590)$, with the following parameters

$$M_G = (1580 \pm 30) \text{ MeV/c}^2$$

$$\Gamma_G = (280 \pm 40) \text{ MeV/c}^2$$

$$\sigma \cdot \text{BR}(G+\eta\eta) = (3.8 \pm 0.7) \text{ nb}$$

Other structures appearing in fig. 2b have been identified as corresponding to $c(1300)$ in the $S$-wave, $f(1270)$ and $f'(1525)$ in the $D$-wave and $h(2240)$ in the $G$-wave. The interpretation is less straightforward for the high mass region in the $S$- and $D$ waves. More details will be found in the forthcoming paper [4].

3. EVIDENCE FOR A HIGH MASS STRUCTURE IN $\eta\eta'$ DECAY

The $\eta\eta'$ channel is another potentially interesting channel which is easily identified among the 4 $\gamma$ events. The data collected at IHEP with 38 GeV incident $\pi^-$ mesons have shown evidence for the decay of $G(1580)$ into $\eta\eta'$ with a B.R. substantially larger than for $\eta\eta$ [5]. This was in fact one of the properties that supported the interpretation of the $G$ meson as a possible glueball [6]. However, other interpretations of the $G$ have been proposed [7]. This analysis has not been done for the data collected at CERN at 100 GeV as the number of $G$ events expected in the sample is rather small.
A 3 to 4 standard deviations significant bump is observed at 38 GeV both for uncorrected and corrected for acceptance events (top of fig. 3) near a mass of 2.2 GeV. The same plot for the 100 GeV events also shows an effect with about the same level of statistical significance (bottom of fig. 3) for exactly the same mass bins. The fact that an accumulation of events appears in two different samples of data taken with different, although similar setups seems to exclude a purely statistical origin for these peaks. The width of the peak, about 120 MeV at 100 GeV and perhaps slightly more at 38 GeV, is comparable to the instrumental resolution for a mass of 2 GeV, so nothing can be said about the real width of this object but it could be narrow. In that case one might be tempted to consider that it is a new decay channel of the $\eta(2230)$ seen in Mark III experiments in the $\bar{K}K$ decay channel.

Fig. 4 shows the angular distributions in the Gottfried-Jackson system of the $\eta\eta'$ pairs for three 120 MeV wide mass intervals, one centered on the peak and the two neighbouring ones.

The acceptances have been calculated by Monte-Carlo methods, and as can be seen by comparing the histograms shown on the left and on the right of fig. 3, the global acceptances do not vary significantly on the mass range of interest here. However, at 38 GeV, the efficiency for events near $\cos \theta_{GJ} = -1$ (forward going $\eta'$) drops to zero near 2.2 GeV.

The angular distribution of events in the peak, with both forward and backward enhancements, is quite different from that of the background. Spin zero is clearly excluded and the presence of events at $\cos \theta_{GJ} = 0$ at 38 and 100 GeV also excludes odd spins. An even spin of 2 or higher is favoured.

If this effect is indeed the $\eta(2230)$, it would be the first time that a meson not fitting in the standard $q\bar{q}$ nonets has been observed both in radiative $J/\psi$ decay and in hadron collisions. Until now, the possible gluonium or hybrid mesons were observed either at $e^+e^-$ colliders (G, $\xi_T$ states) or at proton accelerators (G, $\xi_T$ states), but never at both.

4. THE $\eta\pi^+\pi^-$ CHANNEL

One important question left open in the field of light meson spectroscopy is the $E/\gamma$ puzzle. As fig. 5 shows, the invariant mass plot for the $\eta\pi^+\pi^-$ events selected among the 6 $\gamma$ events shows very clearly two peaks: one at 1280 MeV (D meson) and the second at 1420 MeV (E meson?). The Dalitz plots
of the $\eta\pi\pi$ system for all mass bands between $\sim 1$ GeV and $\sim 2$ GeV show an accumulation of events near 1 GeV for the $\eta\pi$ subsystem. So we believe that there is a large contribution from the $\delta\pi$ channel to the decay of both observed mesons. A Zemach type analysis of the Dalitz plots density (i.e. not using the information coming from the production of the $\eta\pi\pi$ system) has been attempted. The channels involved in the analysis are, in addition to phase space, the $\delta\pi$ and the $\epsilon\pi$ channels where $\epsilon$ is used for the $I=0$ $\pi\pi$ scattering amplitude. The D meson is a well known $1^-$ object, but we cannot use it for calibration purposes as there is apparently also a $0^+$ meson, the $\eta(1285)$ [8] at the same mass. Although the width of $\eta(1285)$ is larger than that of the D, our resolution is not good enough to identify the peak that we see in $\eta\pi^+\pi^-$ as a unique object. In our analysis of the "D" meson mass region we could obtain either a dominant $1^-$ or $0^+$ wave by changing only slightly the parameters of the $\delta$. So it is clear that we are not yet in a position to give a reliable result. A more detailed analysis has to be performed and we hope to have soon a good measurement of the spin of the mass bump near the $E$ which could be an important contribution to the problem mentioned in the beginning.

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Ando et al., presented at the Lepton and Photon Symposium – Kyoto Aug. 85.
FIGURE CAPTIONS

Fig. 1a) Experimental setup. Aside GAMS, it includes: beam hodoscopes (H0 to H2); beam definition scintillators (S0 to S3); a system of veto counters surrounding the target, including scintillator counters (AC, AH1, AH2, AGS) for the rejection of charged particles; lead-scintillator sandwich counters SW for the rejection of charged and neutral particles emitted outside the acceptance of GAMS; lead-glass blocks (VP) for the rejection of $\gamma$ emitted at large angle, coming for instance from the decay of a $N^*$ or $\Lambda$ formed in reaction (1).

1b) Typical 8 $\gamma$ event.

Fig. 2a) $\cos\theta$ distribution of $\eta\eta$ events (4$\gamma$ and 8$\gamma$ decay) for successive mass bins (indicated in GeV in the upper part of each figure). Dashed line curves are Monte-Carlo evaluations of the efficiency.

2b) Amplitudes and solutions for 4$\gamma$ and 8$\gamma$ multiplicities for $\eta\eta$ decays.

Fig. 3 Invariant mass spectra of $\eta\eta'$ system measured at 38 GeV/c (top) and at 100 GeV/c (bottom). The data show on the left are raw data while those on the right have been corrected for efficiency $\epsilon$.

Fig. 4 $\cos G_{\eta\eta}$ of $\eta\eta'$ systems for three 120 MeV/c$^2$ wide mass bins around 2.2 GeV/c$^2$ (corrected for efficiency).

Fig. 5 $\eta\pi^+\pi^-$ invariant mass plot for 6$\gamma$ events at 100 GeV/c incident pion momentum and for $-t_{\eta\pi^+\pi^-} > 0.1 \text{ (GeV/c)}^2$. 
PRELIMINARY RESULTS $\eta \eta'$

$38 \text{ GeV/c}$

$100 \text{ GeV/c}$

FIG 3
PRELIMINARY RESULTS

\[ \Delta M = 2.03 - 2.15 \]

\[ \Delta M = 2.15 - 3.27 \]

38 GeV

\[ \Delta M = 2.27 - 2.39 \]

100 GeV

FIG. 4
FIG. 5

$\eta \pi^0 \pi^0$

100 GeV/c
EFFECTS OF \( \delta \) PARAMETERIZATION IN K\( \pi \) DALITZ PLOT ANALYSIS

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ABSTRACT

We have performed Dalitz plot analyses of data from reaction \( \pi^- p + K^+ K S^- n \) at 8 Gev/c taken at the BNL Multi-particle Spectrometer varying the parameterization of the \( \delta \). We find that a \( 0^{++} \) state can be accommodated under the \( 1^{++} \) D(1285), but it is not required. We also find that final state interactions as parameterized by Frank et al.\(^2\) cannot reproduce the data in the E(1420) region.
The analysis of high statistics data from reactions:

1) \( \pi^- p + K^+ K_S^- n \) at 8 GeV/c (15,000 events)
2) \( \bar{p} p + K^+ K_S^- + X \) at 6.5 GeV/c (12,000 events)

taken by the experiment 771 collaboration at the Multiparticle Spectrometer at BNL showed that the D(1285) is a \( J^{PG} = 1^{++} \) meson decaying mostly to \( \delta \pi \) while the E(1420) is a \( J^{PG} = 0^{-+} \) meson with a substantial \( \delta \pi \) decay mode and some \( K^*R \). The results on the E-meson are in contradiction with two previous experiments that concluded the E is a \( 1^{++} \) meson decaying almost exclusively to \( K^*K \) (Dionisi et al., Armstrong et al.).

A more recent experiment at KEK (Ando et al.) on reaction:

3) \( \pi^- p + n \pi^+ n \) at 8 GeV/c

confirms the results of experiment 771 but observe an additional \( 0^{-+} \) state \( \eta_R(1275) \) in the same mass region as the D(1285). This state was first claimed in a previous analysis of a similar reaction by Stanton et al. but did not seem to be prominent in the experiment 771 data. To determine whether there is any contradiction between these experiments, the Dalitz plot analysis of reaction (1) was redone varying the parameters of the \( \delta \). The results of this reanalysis are presented here. It was found that using parameters that give the best fit in the D(1285) an \( \eta_R(1275) \) is not needed, but one with as large a production cross section as reported by Ando et al. can be accommodated without significantly worsening the fits.

It was also suggested by Frank et al. that near the KR threshold (i.e., the E-region) the \( \delta \pi \) decay mode is purely a final state interaction effect and that the E may in fact have a 100% \( K^*R \) decay mode (this suggestion was actually made for the \( i(1440) \) but should be even more applicable to the E). Their parameterization was tried in the E-region and found that it fits the data badly, without affecting the conclusions about the spin-parity content.

The amplitudes for \( \delta \pi \) decay in a Dalitz plot analysis have the form \( T_\delta \) for \( 0^{-+} \) and \( \bar{p} T_\delta \) for \( 1^{++} \), the difference between them being only the break-up momentum \( p \) for D + \( \delta \pi \). In a coupled channel system with a resonance one can write:

\[
T_\delta = \frac{g_1^2 g_2^2}{g_1^2 + g_2^2} [m_R^2 - m - i m_R (a_1 q_1 + a_2 q_2) \frac{m_T}{m}] (1)
\]

where \( m \) = mass, \( q_1 = \pi \) break-up-momentum, \( q_2 = K^*R \) break-up momentum, \( m_T \) = m at KR threshold, and \( a_1, a_2 \) are unitless constants related to the coupling constants. Isobar model assumption states that one must use
the elastic amplitudes for a given decay mode. If one fits published data for $\delta + \pi \pi$ to a simple Breit-Wigner form one finds $m = 981 \pm 6$ MeV and $\Gamma = 55 \pm 5$ MeV. Define $r = a_2/a_1$ and $\Gamma_1 = a_1 q_1 m \Gamma / m_R$. Using equation (1) the $\pi \pi$ mass spectrum can be reproduced well by varying $m_R$, $r$, $a_1$ ($\Gamma_1$) in a correlated way, i.e., the following values give a fit equivalent to that of the simple Breit-Wigner:

1) $m_R = 978$ MeV, $r = 0.5$, $a_1 = 0.212$ ($\Gamma_1 = 68$ MeV)
2) $m_R = 962$ MeV, $r = 1.5$, $a_1 = 0.387$ ($\Gamma_1 = 122$ MeV)
3) $m_R = 898$ MeV, $r = 2.5$, $a_1 = 1.02$ ($\Gamma_1 = 289$ MeV)

The three sets of values were tried in the D-region (1.25 - 1.33 GeV). The Dalitz plot was fitted dividing the data into 4 20-MeV bins and found that parameter set (1) gave the best fit. The difference in log of the likelihood between sets (1) and (2) was 16 while between (1) and (3) was 33. Set (1) gives the largest amount of $0^{-+}$ wave in the D-region, but it is still substantially smaller than that expected if the KEK results are correct. However, if we force the ratio $0^{-+}/1^{++}$ to be the same as that observed in KEK experiment, we find that the log of the likelihood only gets worse by 5 units. Given the uncertainties in the parameterization this difference cannot be considered significant, so our data can accommodate an $m_R$ (1275) compatible with the state claimed in $\pi p + \pi^+ + \pi^- + n$ reactions but cannot confirm it.

It has been suggested by Frank et al. that the observation of a $\delta \pi$ decay mode near $K^*K$ threshold may in fact be due to final state interactions between the $K^+$ and $K_S$ even though the state may have mostly a $K^*K$ decay mode. They proposed a relatively simple parameterization for final state interaction effects which seems to successfully reproduce the Dalitz plot observed in $i(1440)$ to $K\pi\pi$ observed in radiative decays. The basic ingredient is to replace the width $\Gamma$ in the Breit-Wigner describing the $K^*$ with a distorted $\Gamma_d = \Gamma \sqrt{d_{KK}}$ where $d_{KK}$ is a function of the $K\bar{K}$ breakup momentum, a binding energy and a potential energy (see ref. 2). We have used $\Gamma_d$ in the amplitudes for $0^{-+}$ and $1^{++}$ and removed any direct $\delta \pi$ decay mode amplitudes. These results are shown in Fig. 1. Qualitatively they are not very different, the $1^{++}$ wave still shows a rapid rise after threshold and the $0^{-+}$ wave peaks in the E region. However, the fits are substantially worse throughout:

<table>
<thead>
<tr>
<th>$m$</th>
<th>1.39</th>
<th>1.41</th>
<th>1.43</th>
<th>1.45</th>
</tr>
</thead>
<tbody>
<tr>
<td>difference in log, $L$</td>
<td>18</td>
<td>15</td>
<td>50</td>
<td>40</td>
</tr>
</tbody>
</table>

(20 MeV bins)
They are particularly poor above 1.42 where the K* bands are most distinctive in our data, i.e., the parameterization of Frank et al., fails to reproduce correctly the K* in the 1.42 to 1.46 GeV KR mass region.

In conclusion, we have found that it is possible to accommodate an \( \eta_{\Pi}(1275) \) in the Experiment 771 data but without a firmer knowledge of the \( \delta + \text{KR} \) parameterization, one cannot confirm or disprove its existence with a Dalitz plot type analysis. The question is being studied now with a full partial wave analysis. We have also found that final state interactions as parameterized by Frank et al. to replace a \( \delta \pi \) decay mode cannot reproduce the data in the E-region.

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NEW TESTS FOR POTENTIAL MODELS

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ABSTRACT

Potential models have been applied to the light quark systems with a great deal of success. The fact that the confining potential transforms as a Lorentz scalar and appears to be linear in r allows a unification of the light and heavy meson hyperfine splittings. The new determinations of the two-photon decays of \eta and \eta', the limit on that of the \eta(1440) and the possible confirmation of \eta(1275) pose stringent tests for all potential models.

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INTRODUCTION

It is more than twenty years since the first classification was made of baryons and mesons as hyperfine multiplets related by M1 transitions. Not long after the early successes of potential model descriptions of the J/ψ system, the modern approach to the light baryons and mesons, based on QCD ideas, was set down. Now we know that the dynamics and spectra of the light quark systems, both baryons and mesons, show agreement with a number of features expected from QCD; even the non-relativistic approximations, obviously applicable to the J/ψ and γ systems, seem to be in good agreement with the light quark systematics.

In this talk I want to concentrate on some new positive results of these ideas and to emphasize that we may already be in a position to severely test many of the existing models.

THE LIGHT QUARK SCALE AND THE HEAVY MESONS

The new results of the past year are based on our understanding of the form of potentials derived from systematic studies of baryon spectroscopy and of the heavy meson systems. The important features are that the confining potential is flavour-independent, linear and transforms as a Lorentz scalar whereas, at short distances the spin-dependent forces are those one would expect from single gluon exchange in the QCD model. For S-wave qq mesons, which is what I shall only consider in this talk, this description will suffice. The extent to which this description has been tested and to what extent deviations from it are allowed have been described elsewhere. The fact that the long range confining potential varies as r is usually taken as a prediction of lattice calculations. Recently, it has been pointed out that in Monte Carlo studies of lattice gauge theories, there can also be spin-spin interactions consistent with a short range δ-function-like potential, similar to the single gluon exchange potential of QCD.

Last year at this meeting, an interesting regularity among the vector and pseudoscalar mesons was presented viz., \( m_v^2 - m_p^2 = 0.56 \text{ GeV}^2 \) for all \( q_i \bar{q}_j \) (i ≠ j) mesons. This regularity holds for the \( \pi-\rho, K-K^*, D-D^*, \) and \( B-B^* \) mesons, i.e., from the lightest to the heaviest of the known mesons! This regularity can be understood in view of the properties of the potentials described above. For a flavour-independent Lorentz scalar potential, many of the effects of binding can...
be included in the effective, or constituent, quark masses\(^9\)). That is, the constituent quark masses already contain the effects of the non-perturbative confining forces. The multiplet splittings can be thought of as arising from the spin-dependent interactions. Recently, these have been classified in a general way\(^3\)-\(^6\); for S-wave mesons only the hyperfine splittings are involved.

The masses of a vector \((V_{ij})\) and pseudoscalar \((P_{ij})\) meson \(q_iq_j\) \((i \neq j)\) are thus mainly described by the sum of the constituent quark masses \((m_i + m_j)\). The mass splitting between \(V_{ij}\) and \(P_{ij}\) is given in terms of the quantity \(\frac{\psi(0)}{m_i m_j}^2\), where \(\psi(0)\) denotes the wave function of the the \(q_iq_j\) system evaluated at the origin. Usually, the lack of knowledge about \(\psi(0)\) prevents a quantitative extension of this scheme to mesons composed of different quarks. A new observation that permits such an extension was made last year\(^10\). If the long range confining potential has a linear behaviour, \(V(r) \sim r\), then a simple scaling law based on the Schrödinger equation shows that the square of the wave function, calculated at the origin, is proportional to the reduced mass \(\mu_{ij}\) of the quark-antiquark pair. That is \((m_i + m_j) \langle V_{hf} \rangle\) is independent of the quark flavours, \(i, j\). It is precisely this property that leads to the relationship \(m_V^2 - m_P^2 = \text{constant}^{11}\) for all of the mesons. [If \(i = j\) then there are extra complications due to the possibility of annihilation graphs\(^10\)]. In view of the fact that the reduced mass of a system consisting of one light and one heavy quark is approximately that of the light quark, we see the importance of the light quark scale for all of the meson system, at least for mass splittings. This unification of the mass splittings among the mesons may also be evidence for the applicability of the non-relativistic approach throughout the spectra.

**RADIAL EXCITATIONS AND GLUEBALLS**

The hyperfine interaction just considered, has played an important role in applications of potential models to the light mesons. In a number of calculations\(^12\)-\(^15\), the mixing among the radial wave functions \(1S\) and \(2S\) (and even \(3S\)) comes from such an interaction. These models generally do reasonably well in describing the spectra of the light mesons; here I shall concentrate on the isoscalar-pseudoscalar spectra.

A general feature of these models is the prediction of two radial excited states \((\eta_2, \eta_2')\) of the usual \((\eta, \eta')\) mesons in the region between 1 GeV and about 1.5 GeV. [Ref. 14) included a model which has an additional gluonium component.
I shall refer to the extra state as $G$; the purely radial models will be characterized as $(qq)^*_{\text{models}}$. The details of the flavour content and mixing vary among the different calculations and there is also a slight variation among the predicted masses although these are sufficiently consistent as to be within the errors of the calculations, which I estimate to be of the order of $10$-$15\%$.

In all of these models, except one\(^{14}\), the mixing involves only $1S$ and $2S$ states. Inclusion of a third radial excitation distorts the mass of the heavier $2S$ and moves it up by about $400$ MeV to about $2$ GeV. Although the mixing of even higher radial states may restore the mass to the $1.6$ GeV region, the properties of the heavier $2S(qq)^*$ state must be viewed with some care in all models. The decays $J/\psi + \gamma\eta, \gamma\eta', \gamma\gamma$ have been shown\(^{16}\) to rule out a pure $(\bar{q}q)^*$ interpretation of the $1(1440)$. If the wave functions have not been determined\(^{13}\) then perhaps there is still a slight possibility of getting around this conclusion. It should be emphasized that fits to mass spectra are no longer sufficient in the present context; quite distinct interpretations\(^{17,18}\) can be made of the same mass spectra. Only a simultaneous description of masses and decays will distinguish among them; without wave functions the predictions are not too convincing. (This is a problem which will face the lattice calculations also whenever the mass spectra are finally determined.)

The new features that all of these models have to confront are the determination\(^{19}\) of the $2\gamma$ widths of $\eta$ and $\eta'$, the upper limit\(^{20}\) of $2$ keV for $\Gamma(\gamma + 2\gamma)\eta(\gamma + K\bar{K}\pi)$ and the apparent confirmation of the state $\eta(1275)$. In the recent two-photon experiments\(^{19}\), the new determinations for $\eta$ and $\eta'$ are $\Gamma(\eta + 2\gamma) = 0.56 \pm 0.04$ keV and $\Gamma(\eta' + 2\gamma) = 4.3 \pm 0.1$ keV. It is a stringent test for the potential models to be consistent with these values.

The state $\eta(1275)$ was first reported about seven years ago\(^{21}\). In the past year a confirmation of this claim has been made\(^{22}\). If this claim holds up then I think that all of the purely radial excited models are in trouble. The principal reason for this is the near degeneracy such a state has with the $\pi(1300)$. For this [and other reasons\(^{17}\)] it is unlikely that any of the $(\bar{q}q)^*$ models survive in their present form. Of the list quoted\(^{12}-15\) only the $q\bar{q}+G$ model of Ref. 14) emerges unscathed, albeit with a prediction for $\Gamma(\gamma + \gamma\gamma) \sim 2.5$ keV, uncomfortably close to the present experimental limit.

**CONCLUSIONS**

The successes of potential models in describing the light quark systems have now been extended to include the heavy mesons via the regularity of $m_V^2 - m_F^2$. 
The hyperfine splitting of the mesons has the scale of the light quarks throughout the S-state meson systems. The new measurements of the widths $\eta, \eta' + 2\gamma$ together with the upper limit $\Gamma(\pi + \gamma\gamma)B(\pi + K\pi) < 2$ keV provide a stringent test for potential models. The further confirmation of the existence of the state $\eta(1275)$ would be of great importance for delineating the properties of any future model.

ACKNOWLEDGEMENTS

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19) For a review, see
    H. Kolanoski, Proc. Kyoto Conference (1985) and these proceedings.
A LATTICE MONTE CARLO CALCULATION OF THE GLUON PROPAGATOR

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ABSTRACT

A Monte Carlo calculation of the gluon propagator in the Landau gauge in SU(3) lattice gauge theory is described. Gribov ambiguities are briefly discussed. The results of calculations at $\beta = 5.6$ (200 $4^3 \times 8$ lattices), $\beta = 5.8$ (400 $4^3 \times 10$ lattices), and $\beta = 6.0$ (100 $4^3 \times 8$ lattices) indicate that the gluon propagator resembles a massive particle propagator with asymptotic mass near 600 MeV.

*Presented by Jeffrey E. Mandula
In this talk we will describe a calculation of the gluon propagator in the Landau gauge in SU(3) Wilson lattice gauge theory. Our goal in this calculation is the use of lattice gauge Monte Carlo techniques to estimate a function which, while not directly measurable, is one of the most fundamental quantities in a Yang-Mills theory. In this way, lattice gauge techniques can make contact with continuum approaches to gauge theories. The gluon propagator will be calculated from first principles in QCD with the same level of confidence as other quantities on the lattice. The techniques we use can be employed to calculate other propagators and gauge dependent quantities of interest in field theories. These calculations are important not because they directly verify QCD or show that it is inadequate, but because they offer a window into the way that QCD "works" with theoretical and phenomenological implications.

The principal result of this investigation is that at large distances the gluon propagator in the Landau gauge falls off exponentially like a massive particle propagator. The mass is about 600 MeV. Although the bare gluon propagator has a singularity at \( q^2 = 0 \), in the dressed propagator this singularity has moved up to a finite mass. This is an instance of a purely dynamical Higgs mechanism, without an explicit Higgs field.

The possibility of dynamical mass generation without an explicit scalar field was first pointed out by Schwinger, before the discovery of the conventional Higgs mechanism. He gave an explicit example of this phenomenon, two-dimensional massless quantum electrodynamics. The possible occurrence of a dynamical Higgs mechanism in four-dimensional gauge theories has been studied by Jackiw and Johnson and by Cornwall and Norton. All of the above work is concerned with gauge fields coupled to massless fermions, whereas we work in the pure gauge sector. Our results provide the first hard evidence for the existence of a dynamical Higgs mechanism in a four-dimensional gauge theory.

There is indirect theoretical evidence from other lattice gauge studies that supports the idea that the gluon is massive, although this evidence has not always been interpreted in terms of an effective gluon mass. The fact that as \( N \to \infty \) there survives a linear potential between adjoint sources in SU(N) gauge theory ("the adjoint string doesn't break") suggests that, in the sense of a constituent object, not only is the gluon massive, but that as \( N \to \infty \) its mass becomes infinite.

Bernard has performed lattice Monte Carlo studies of the existence and breaking of an adjoint string in SU(2) gauge theory. As those studies note, the association of a gluon mass with the distance at which the adjoint string breaks is ambiguous, because the measured mass value includes an unknown binding energy to the adjoint source.

The gluon propagator, and especially its infrared behavior, has been the focus of several discussions that aim at explaining confinement or address other fundamental properties of QCD. Arguments have been made, based on truncations of the Schwinger-Dyson equations and Ward identities, that there is a \( 1/q^4 \) infrared singularity in the gluon propagator, and models of confinement based on that singularity have been constructed. Different approximations to the Schwinger-Dyson equations suggest a \( 1/q^2 \ln^2 q^2 \) IR singularity, and, by studying a gauge invariant function somewhat related to the axial gauge propagator, it has been argued that the gluon is effectively massive. Of course, perturbation theory gives a \( 1/q^2 \) behavior. Our calculations directly determine the infrared behavior of the gluon
propagator.

In the rest of this note we will summarize the calculation: the specification of a lattice Landau gauge, our method for calculating in that gauge with the proper functional measure, and the results. We will also comment on phenomenological implications, and make a brief remark about the problem of Gribov copies.

The association between lattice and continuum variables is not unique. We make a maximally local choice for the lattice gauge potential

\[ A_\mu(n) = (U_\mu(n) - U_\mu(n - J)) \text{Traceless/}2ia. \]

We implement the Landau gauge by maximizing the quantity

\[ \text{Re} \sum_\mu (U_\mu(n) + U_\mu^T(n - \mu)) \]

at each site. This implies that the maximally local difference form of the Landau gauge condition, \( \sum_\mu (A_\mu(n) - A_\mu(n - \mu)) = 0 \), holds.

It is very difficult to do Monte Carlo calculations in a fixed gauge by using a Metropolis algorithm on trial updates that preserve the gauge condition, starting from a configuration that satisfies it. Sweep-to-sweep correlations between lattices updated in this way can be much greater than in the usual non-gauge-constrained Monte Carlo method, so that many more sweeps are needed between each pair of statistically independent lattices; this is known to occur in the case of discrete gauge groups. Also, it would be necessary, for each Monte Carlo trial, to calculate the Faddeev-Popov determinant, making the calculation as computationally intensive as the inclusion of dynamical fermions.

To avoid these problems, we update the lattice using the conventional Metropolis Monte Carlo algorithm without any gauge constraint. We then independently gauge transform each lattice into the Landau gauge. This gives us an ensemble of Landau gauge lattices, correctly distributed according to the Wilson action and Faddeev-Popov determinant, without having to explicitly compute a Faddeev-Popov determinant.

The basis for this algorithm lies in the Faddeev-Popov analysis, which we briefly review. The Faddeev-Popov determinant \( \Delta_f(U) \) corresponding to the gauge condition \( f(U) = 0 \) is defined by

\[ \Delta_f(U) \int dg \delta(f(U^g)) = 1 \]

where \( U^g \) is the gauge transform of \( U \) and the integration is over all gauge transformations \( g \). \( \Delta_f(U) \) is gauge invariant by construction. Multiplying the gauge invariant functional integration measure \( DU \) by 1 in the form of Eq. (2), interchanging orders of integration, changing variables \( U \rightarrow U^{g^{-1}} \), and noting that both \( DU \) and \( \Delta_f(U) \) are gauge invariant, gives the functional integral subject to the gauge constraint:

\[ \int DU = \int DU \Delta_f(U)\delta(f(U)) \]

where we have conventionally taken the group volume to be 1.

The expected value of any function of link variables \( \Phi(U) \), gauge invariant or local or not, is given by

\[ \langle \Phi(U) \rangle_f = \frac{1}{Z} \int DU \Delta_f(U)\delta(f(U))\Phi(U) \]
where the normalization $Z$ is the same integral without the $\Phi(U)$ factor. We reverse the Faddeev-Popov analysis by integrating $\int dg$ (on which nothing depends), interchanging the order of integration, and changing variables $U \rightarrow U^g$. We can do the $dg$ integration because of the $\delta$ function, using Eq. (3). Only the value of $\Phi(U^g)$ at that value of $g$ which maps $U$ into the gauge $f = 0$ is needed. The result is

$$<\Phi(U)>_f = \frac{1}{Z} \int DU \ \Phi(U^g(U))$$

(5)

where $g(U)$ is that gauge transformation for which $f(U^g(U)) = 0$.

Equation (5) is the statement of the algorithm. In words it says that the expected value of any quantity $\Phi(U)$ in the gauge $f(U) = 0$ may be calculated from an ensemble of configurations weighted without gauge constraint ($\int DU$), but that each configurations should be gauge transformed into the $f(U) = 0$ gauge before evaluating $\Phi$.

In the foregoing analysis, we have tacitly assumed that the condition $f(U) = 0$ specified the gauge uniquely. However, it is known that many differential conditions, such as $\partial_\mu A_\mu = 0$ in the continuum, are subject to Gribov ambiguities. This may be so with our lattice Landau gauge condition as well, but since it is a global maximization rather than a difference condition, the Gribov ambiguity may be absent.

Our method for transforming a lattice into the Landau gauge consists of sweeping through the lattice, applying at each site the gauge transformations that maximizes the quantity Eq. (1). This procedure is somewhat delicate, since a gauge transformation at one site upsets the gauge condition at the surrounding sites. We know of no proof that the procedure converges. Empirically, however, the method seems to work very well.

We have used this method to compute the gluon propagator in SU(3) pure Yang-Mills theory for $\beta = 5.6$ and 6.0 on a $4^3 \times 8$ lattices and for $\beta = 5.8$ on a $4^3 \times 10$ lattice, all with periodic boundary conditions. Simulations were performed on two machines, a Ridge-32 and Cray-XMP, with different acceptance rates, numbers of hits, and sweeps between each measurement, as well as slightly different update procedures. The data were combined only in the final analysis. We used 200 lattices for $\beta = 5.6$, 400 lattices for $\beta = 5.8$, and 100 lattices for $\beta = 6.0$. Each lattice was separated by 10-30 sweeps, with 10-20 hits per sweep. We fixed the gauge well enough so that $<Tr(\partial_\mu A_\mu)_\text{lattice}>^2$ was always less than 0.01, and often less than 0.0005. To verify that our results are stable to better gauge fixing, we reran the $\beta = 5.6$ case, taking 100 lattices, but spending twice the time gauge fixing, with no significant change in the gluon propagator.

We examined the time displacement correlations of the gauge potential summed over all sites on each fixed $t$ (space-like) hyperplane $A_\mu(t) = \Sigma A_\mu(t, \vec{x})$. In the Landau gauge with periodic boundary conditions, $A_0(t)$ should be a constant, and we measure $Tr(<A_\mu(t)A_\mu(0)>$ as a function of $t$ to monitor that we have adequately fixed the gauge. The dynamical content of the Landau gauge potential is expressed in $\sum Tr(<A_i(t)A_i(0)>$. Both the $<A_i(t)A_i(0)>$ and the $<A_\mu(t)A_\mu(0)>$ correlations are shown in the Figures 1. The errors are statistical errors on the measured propagators from each ensemble, including corrections for residual sweep to sweep correlation in the few cases these were detectable.
To guide the eye, note that the free propagator for mass \( m \), summed over spatial sites, is
\[
\Delta_m(n_t) = C \cosh(aM | n_t - N_t/2 | )
\]  
(6)
where \( \sinh(aM/2) = am/2 \), \( C \) is a normalization constant, and \( N_t \) is the total number of sites in the \( t \) direction. Thus on the periodic lattice an ordinary massless particle propagator becomes \( \Delta_0(n_t) = \text{constant} \), and its momentum-space square becomes the convolution of a constant with itself, which is also constant.

Far from being constant, the actual propagators \( <A_i A_i> \) in Figs. 1 look like massive particle propagators. In fact, using Eq. (6) we can compute an effective mass governing the falloff from each hyperplane to the next. These are given in Table 1. The errors are estimated by binning the data and are statistical only.

Figure 1: Spatial and temporal gluon propagators as a function of \( t \). Errors are statistical only in the individual data points. a) \( \beta = 5.6 \) on \( 4^3 \times 8 \) lattice. b) \( \beta = 5.8 \) on \( 4^3 \times 10 \) lattice. c) \( \beta = 6.0 \) on \( 4^3 \times 8 \) lattice.
Table 1.
Correlation functions versus separation. Errors are statistical only.

<table>
<thead>
<tr>
<th>$\beta = 5.6 \ a^{-1} = 0.80 \text{ GeV}$</th>
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</thead>
<tbody>
<tr>
<td>$4^3 \times 8$ lattice</td>
</tr>
<tr>
<td>$\Delta t$</td>
</tr>
<tr>
<td>1</td>
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<td>2</td>
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<td>3</td>
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<td>4</td>
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<table>
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<tr>
<th>$\beta = 5.8 \ a^{-1} = 1.27 \text{ GeV}$</th>
</tr>
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<tbody>
<tr>
<td>$4^3 \times 10$ lattice</td>
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<tr>
<td>$\Delta t$</td>
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<tr>
<td>1</td>
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<td>2</td>
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</table>

<table>
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<tr>
<th>$\beta = 6.0 \ a^{-1} = 1.69 \text{ GeV}$</th>
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<tbody>
<tr>
<td>$4^3 \times 8$ lattice</td>
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<tr>
<td>$\Delta t$</td>
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<td>2</td>
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A striking aspect of Table 1 is that for each $\beta$ value the effective mass grows with separation. Such behavior is only possible if the spectral function describing the gauge potential propagator is not positive definite. Otherwise, the effective mass must fall monotonically. A rise is entirely possible in the Landau gauge, which has ghosts.

In Figure 2, we plot the $<A_iA_j>$ propagators at all three values of $\beta$, all normalized to one at zero separation. As may be seen from the figure, the agreement with scaling is good. Note that the range of physical distance is different for each value of $\beta$. The cosh form of the propagator results in the last point for each $\beta$ value being above the envelope.
The calculation has a remarkably high intrinsic "signal to noise" ratio. Even at separations of 4 or 5 lattice spacings, the value of the propagator is reasonably well determined with only 100 lattices. This should be compared to calculations of glueballs of comparable mass without using spatially distributed sources. In these calculations, with hundreds or thousands of times as many lattices, the computed correlations are lost in statistical noise beyond 2 spacings.

The calculations reported here have been performed on fairly small lattices and at fairly small $\beta$. In the range of $\beta$ from 5.6 to 6.0, dynamical scaling seems to hold, though there are noticeable deviations from the two loop asymptotic formula. To estimate finite size effects and possible scaling violations, we express the effective masses from Table 1 in physical units, using the values of the inverse lattice spacing determined by Barkai, Moriarty, and Rebbi from the string tension ($\sqrt{\sigma} = 0.42$ GeV) on large lattices at each of our values of $\beta$. Note that in physical units, the $\beta = 6.0$ lattice is smallest, so that finite size effects are greatest, while the $\beta = 5.6$ lattice is coarsest.

It is natural to associate the effective mass determined at asymptotic distances with a constituent gluon mass. From our present results, the best estimate of this quantity is about 600 MeV, with finite size effects, possible scaling violations, and statistical uncertainties of at least $\pm 25\%$. The massiveness of the gluon might lead to an explanation of the apparent
suppression of many-gluon intermediate states in \( J/\psi \) and \( \tau \) decay, and the relative absence of mixing between the lowest quark model states and those with gluonic excitations. In addition, the QCD part of jet production models will be modified.

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IMPLICATIONS OF SCALAR CONFINEMENT FOR CHIRAL SYMMETRY BREAKING IN QCD

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ABSTRACT

I discuss the implications of scalar confinement for gap equation, Coulomb gauge models for chiral symmetry breaking in QCD.
In a prophetic paper, Nambu and Jona-Lasinio proposed that the pion is an almost massless fermion-anti-fermion bound state. A number of authors have attempted to get analytic realizations of this idea within QCD, by deriving "pairing models" or "gap equations" as approximations to the exact QCD equations. Particularly interesting in this regard are several papers showing that chiral symmetry breaking, and a Nambu-Goldstone pion, necessarily occur when the instantaneous Coulomb gauge potential has a confining piece. My purpose in this talk is to outline the derivation of the Coulomb gauge gap equation models, and then to remark that a key assumption in this derivation is contradicted by recent quark spectroscopic evidence on the Lorentz structure of the confining potential.

A simple and general method for obtaining gap equation models proceeds as follows: We start from the Dyson equations for the vector and axial-vector vertices,

\[ \tilde{\Gamma}_\mu(p', p)_{\delta\gamma} = (Z_{\mu \gamma})_{\delta\gamma} + \int \frac{d^4q}{(2\pi)^4} \left[ i \tilde{S}_F(p'+q) \tilde{\Gamma}_\mu(p'+q, p+q)i\tilde{S}_F(p+q) \right]_{\beta\alpha} \times \tilde{K}_{\alpha\beta, \gamma\delta} (p+q, p'+q, q) , \]

\[ \tilde{\Gamma}_5(p', p)_{\delta\gamma} = (Z_{\gamma 5})_{\delta\gamma} + \int \frac{d^4q}{(2\pi)^4} \left[ i \tilde{S}_F(p'+q) \tilde{\Gamma}_5(p'+q, p+q)i\tilde{S}_F(p+q) \right]_{\beta\alpha} \times \tilde{K}_{\alpha\beta, \gamma\delta} (p+q, p'+q, q) , \]

with \( \tilde{S}_F \) the renormalized quark propagator, \( \tilde{K}_{\alpha\beta, \gamma\delta} \) the renormalized quark-anti-quark Bethe-Salpeter kernel, and with \( Z \) the vertex renormalization constant. The vertices \( \tilde{\Gamma}_\mu \) and \( \tilde{\Gamma}_5 \) obey the Ward identities

\[ (p' - p)^\mu \tilde{\Gamma}_\mu(p', p) = \tilde{S}_F^{-1}(p') - \tilde{S}_F^{-1}(p) , \]

\[ (p' - p)^\mu \tilde{\Gamma}_5(p', p) = \gamma_5 \tilde{S}_F^{-1}(p) + \tilde{S}_F^{-1}(p')\gamma_5 . \]

We now make the fundamental approximation of assuming
(i) \( \tilde{k}_{\alpha\beta,\gamma\delta}(p+q, p'+q, q) = \tilde{k}_{\alpha\beta,\gamma\delta}(q) \), \( (3a) \)

with \( \tilde{k} \) a function only of the momentum transfer \( q \), and

(ii) \( \tilde{k} \) has Lorentz-vector couplings on the quark and antiquark lines so there is no explicit breaking of chiral symmetry,

\[
\tilde{k}_{\alpha\beta,\gamma\delta}(q) \ (\gamma_5)_{\gamma',\gamma} = -(\gamma_5)_{\alpha\alpha}, \ \tilde{k}_{\alpha\beta,\gamma\delta}(q), \\
\tilde{k}_{\alpha\beta,\gamma\delta'}(q) \ (\gamma_5)_{\delta',\delta} = -(\gamma_5)_{\beta\beta}, \ \tilde{k}_{\alpha\beta,\gamma\delta}(q). \ (3b)
\]

Substituting (i) and (ii) into the Dyson equations and using the Ward identities twice implies that \( \tilde{S}_F^i \) obeys

\[
\tilde{S}_F^{-1}(p) = (Z_{\mu})_{\mu} p^\mu + \int \frac{d^4q}{(2\pi)^4} \tilde{S}_F^i(p+q)_{\beta\alpha} \tilde{k}_{\alpha\beta,...}(q). \ (4)
\]

Introducing the self-energy \( \Sigma \) by

\[
\tilde{S}_F^{-1}(p) = \gamma_5 p^\mu - \Sigma(p), \ (5a)
\]

\( \Sigma(p) \) then obeys the integral equation

\[
\Sigma(p)_{\delta\gamma} = [\gamma_5 - (Z_{\mu})_{\delta\gamma}]_{\delta\gamma} p^\mu - \int \frac{d^4q}{(2\pi)^4} \tilde{S}_F^i(p+q)_{\beta\alpha} \tilde{k}_{\alpha\beta,\gamma\delta}(q). \ (5b)
\]

Equation (5) is the renormalized gap equation corresponding to the ladder approximation to the Bethe-Salpeter kernel.

If we now assume that \( k \) is given by an instantaneous potential,

\[
\tilde{k}_{\alpha\beta,\gamma\delta}(q) = -4\pi i(\gamma_0)_{\delta\beta} (\gamma_0)_{\alpha\gamma} \frac{4}{3} V(|q|), \ (6)
\]

then Eq. (5) gives the gap equation for the pairing model. In particular, a linear confining potential

\[
V(|q|) = \frac{\sigma}{(|q|^2)^\frac{3}{2}}, \ (7)
\]

gives chiral symmetry breaking solutions to the gap equation.
Thus it appears that one has achieved a concrete realization of the Nambu-Jona-Lasinio idea. Unfortunately, there is a problem. As we have seen, the derivation of the gap equation requires that $k$ have Lorentz vector couplings on the quark and antiquark lines. A Lorentz scalar instantaneous potential analogous to Eq. (6),

$$\tilde{\kappa}_{\alpha\beta, \gamma\delta}(q) = -4\pi i (1)_{\delta\delta} (1)_{\alpha\gamma} \frac{4}{3} V_S(|q|),$$

manifestly breaks chiral symmetry, and hence does not lead to a Nambu-Goldstone pion. However, quark spectroscopy$^5$ and theoretical arguments$^6$ now show quite convincingly that the heavy quark static potential has the form

$$V(r) = V^S(r) + V^V(r),$$

$$V^S(r) = \text{Lorentz-scalar piece} = \sigma r,$$

$$V^V(r) = \text{Lorentz-vector piece} = -\frac{4}{3} \frac{\alpha}{r}.$$  \hspace{1cm} (9)

In other words, while the Coulombic piece has (as expected from perturbation theory) Lorentz vector couplings, the non-perturbative confining potential is predominantly (very likely totally in the $r \to \infty$ limit) Lorentz scalar.

To conclude,$^7$: There is a puzzle here, and an indication that the approximations leading to the gap equation are not valid for the confining part of the potential. One may have to use the full Bethe-Salpeter kernel $\tilde{K}(p+q, p'+q, q)$ - or at least a more sophisticated approximation than the ladder approximation - to understand chiral symmetry breaking in QCD.

Acknowledgment

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7] See also S.L. Adler, "Gap Equation Models for Chiral Symmetry Breaking," to appear in a special issue of Progress in Theoretical Physics (Supplement), dedicated to Professor Yoichiro Nambu on the occasion of his 65th birthday.
CALCULABLE NON-PERTURBATIVE EFFECTS IN SUPERSYMMETRIC GAUGE THEORIES

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ABSTRACT

We discuss the consistency of the instanton calculus for supersymmetric QCD.
The search for theories in which various infinities could cancel among themselves has been initiated a long time ago by Stueckelberg\textsuperscript{1}) with the interesting conclusion that this aim could be achieved if a special relation between the weak and electromagnetic couplings of the electron is satisfied (as far as I know this is the first proposal for unification). More recently, much interest has been devoted to supersymmetric theories\textsuperscript{2}) for a similar reason: bosons and fermions can conspire in order to "soften" the ultra-violet behaviour of the theory. At the perturbative level, those cancellations can be expressed as non-renormalization theorems which imply in particular that vacuum expectation values differing by many orders of magnitude can be maintained at the quantum level (one aspect of the hierarchy problem). Another interesting feature of supersymmetric gauge theories is the fact that some non-perturbative effects (condensates, instanton effects,...) can be computed explicitly as was first noticed by the ITEP group\textsuperscript{3}) in the case of the supersymmetric Yang-Mills theory.

This talk will be mainly devoted to the consistency of the instanton calculus for supersymmetric QCD\textsuperscript{6)},\textsuperscript{7}) (denoted SQCD hereafter).

SQCD is a SU(N) gauge theory with a massless fermion (\(\lambda_i\)) in the adjoint representation, M (flavour) chiral supermultiplets \([\phi_i, \psi_{\alpha i}]\) (Weyl fermion) in the \(\mathbf{N}\) representation and M chiral supermultiplets \([\tilde{\psi}_i, \tilde{\psi}_{\alpha i}]\) in the \(\mathbf{\bar{N}}\) representation (\(i=1,\ldots,M\) and the colour indices will always be implicit. The Lagrangian reads

\[
\mathcal{L} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} - i \bar{\sigma}^{\mu} D_{\mu} \lambda - i \bar{\psi}_j \sigma^{\mu} D_{\mu} \psi_j - i \bar{\tilde{\psi}}_j \tilde{\sigma}^{\mu} D_{\mu} \tilde{\psi}_j - m_{ij} \psi_i \psi^*_j - m_{ij} \bar{\psi}^*_i \bar{\psi}_j - (D^\mu \psi_j)^* (D^\mu \psi_j) - (D^\mu \bar{\psi}_j)^* (D^\mu \bar{\psi}_j) - m_{ij} m_{j'k} (\phi_i \phi_{j'}^* + \phi_{i'} \phi_k^*) + i \sqrt{2} g \phi_j^* \lambda \psi_j - i \sqrt{2} g \bar{\psi}_j \lambda \bar{\psi}_j - i \sqrt{2} g \bar{\psi}_j \lambda \psi_j - i \sqrt{2} g \bar{\psi}_j \lambda \bar{\psi}_j - \frac{i}{2} g^4 \left( (\phi_j^* T^a \phi_j) (\phi_k^* T^a \phi_k^*) - (\bar{\psi}_j^* T^a \bar{\psi}_j) (\bar{\psi}_k^* T^a \bar{\psi}_k^*) \right)
\]

Interesting information about the low-energy physics (vacuum, light excitations) of SQCD have been obtained by Taylor, Veneziano and Yankielowicz\textsuperscript{4}) by using the effective Lagrangian approach. In particular they found the following relations among the condensates.
\[
\langle \lambda \lambda \rangle = \Lambda^3 \prod_{i=1}^{\frac{m}{N}} \left( \frac{m_i}{\Lambda} \right)^{1/N}
\]
\[
\langle \varphi_i \bar{\varphi}_j \rangle = \delta_{ij} \left( \frac{m_i}{\Lambda} \right)^{4} \langle \lambda \lambda \rangle
\]

More generally, gauge and Lorentz invariant Green's functions involving only lowest components of chiral superfields possess remarkable properties due to supersymmetry. Namely, it has been shown\(^5,6\) that Green's functions of the form
\[
\mathcal{G}_{MN}^{pq} = \langle \lambda \lambda (x_1) \cdots \lambda \lambda (x_p) \bar{\varphi}_i \varphi_j (x_{p+1}) \cdots \bar{\varphi}_q \varphi_j (x_{p+q}) \rangle
\]
are: a) independent of the space-time variables \(x_i\); b) independent of \(m^*_i\); c) dependent of \(m_i\) according to
\[
\prod_{\lambda=1}^{\frac{M}{N}} \left( \frac{m_{\lambda}}{\Lambda} \right)^{p+q} \prod_{\ell=1}^{\frac{q}{N}} \left( \frac{m_{i\ell} m_{j\ell}}{\Lambda^2} \right)^{-\frac{1}{2}}
\]

These Green's functions can be computed in the background of an instanton configuration\(^7\). When the masses are non-zero, there are no flat directions and the only classical minimum is
\[
\langle \varphi_i \rangle = \langle \bar{\varphi}_j \rangle = 0
\]

Property a) implies that the Green's functions are the same in the short-distance regime (where the semi-classical approach is meaningful) and in the large distance regime. However, in the massless case, the potential has flat directions, and large vevs are classically allowed (they play the role of infra-red regulators). It is not clear that a consistent Higgs picture also exists in the massive case [for more details about this alternative approach, see Refs. 8]?

The consistency of the method can be checked explicitly in the massive case. For \(N = 2\) and \(M = 1\), we obtain
\[
\langle \lambda \lambda (x) \lambda \lambda (y) \rangle = C m \Lambda^5 \frac{2g^4}{\pi^2} \int d^4x_0 \, d^4s \, \frac{g^4}{s^g} f_x(x) f_y(y) (x-y)^2
\]
\[
= C m \Lambda^5 \frac{4g^4}{\pi^2} ; \quad C = g^4 2^{10} \pi^6 ; \quad f_x = \left( (x-x_0)^2 + s^2 \right)^{-\frac{1}{2}}
\]
\[
\langle \lambda \lambda (x) \varphi \bar{\varphi} (y) \rangle = 2 C g^4 \Lambda^5 \int d^4x_0 \, d^4s \, \sum_{\lambda=1}^{\frac{1}{2}} \left[ \lambda_0 \lambda_0 (x) \lambda_0^2 \right] \frac{1}{D^2 - m^2_1} \lambda_0 (z_1) \frac{1}{D^2 - m^2_2} \lambda_0 (z_2)
\]
\[
\frac{d^4z_1 \, d^4z_2}{D^2 - m^2_1} \lambda_0 (z_1) \frac{1}{D^2 - m^2_1} \lambda_0 (z_2) \frac{1}{D^2 - m^2_2} \lambda_0 (z_2)
\]
The \( \Lambda_0 \) denotes the zero-modes of the gluino [see Refs. 7) for details]. The last expression is in general quite complicated but property c) guarantees that it is \( m \)-independent. Consequently, the calculation can be performed in the limit

\[
\frac{1}{\Lambda^2} \gg (x-y)^2 \gg \frac{1}{|M|^2}
\]

This gives a constant plus various corrections of order \( m^2(x-y)^2 \) whose sum vanishes. The final result is

\[
\langle \lambda \lambda (x) \; \bar{\psi} \bar{\psi} (y) \rangle = \frac{C g^4 \Lambda^5}{20 \pi^4}
\]

in agreement with the supersymmetric Ward identities\(^9\).

Note that in the strictly massless case we would have obtained

\[
\langle \lambda \lambda (x) \; \bar{\psi} \bar{\psi} (y) \rangle = \frac{C g^4 \Lambda^5}{32 \pi^4}
\]

This shows that the limit \( m \to 0 \) is not equivalent to the strictly massless case (where the Higgs description seems more suitable).

Finally let me briefly mention the possibility of a dynamical supersymmetry breakdown in the presence of chiral fermions\(^10\) and various applications of the effective Lagrangian approach for composite models of quarks and leptons\(^11\).

References

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11) See, e.g.:  
We show how the Schwinger Dyson equations of Yang Mills theory yield an effective Lagrangian expressed in terms of electric vector potentials. The classical equations of motion generated by this Lagrangian have a solution corresponding to a confined tube of quantized color electric flux. Semi classical quantization around this solution yields the QCD string.

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In this talk I will show how the Schwinger Dyson equations of Yang Mills theory lead to an effective long distance Lagrangian expressed in terms of electric vector potentials $C_\mu$. This Lagrangian has solutions corresponding to confined tubes of electric flux. Semi classical quantization around this solution then yields the QCD string.

The Yang Mills Lagrangian is given by

$$L_{YM} = 2 \text{Trace} \left( -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} \right)$$

where $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu - i e [A_\mu, A_\nu]$, and $A_\mu = \sum_\alpha A^{\alpha}(x) T_\alpha$. $A^{\alpha}(x)$ are the Yang Mills fields and $T_\alpha$ are the generators of the color gauge group in the fundamental representation. They satisfy the commutation relations $[T_\alpha, T_\beta] = i f_{\alpha\beta\gamma} T_\gamma$ and are normalized so that $2 \text{Trace} T_\alpha T_\beta = \delta_{\alpha\beta}$.

We have used $e$ to denote the Yang Mills coupling constant, i.e.

$$\alpha_s = e^2 / 4\pi$$

$L_{YM}$ was constructed to be invariant under the non Abelian gauge transformations,

$$A_\mu \to \Omega^{-1}(x) A_\mu \Omega(x) + (i / e) \Omega^{-1}(x) \partial_\mu \Omega(x)$$

where $\Omega(x)$ is an element of the color gauge group.

First we introduce electric vector potentials $C_\mu$ and show why they are natural variables for studying long distance Yang Mills theory. To do this we consider the Wilson loop $W_A(\gamma)$ defined as

$$W_A(\gamma) = \mathcal{P} \exp \left[ \int_{\gamma} dx^\mu A_\mu(x) \right]$$

where $\gamma$ is any closed curve and the symbol $\mathcal{P}$ indicates path ordering. Let $|0\rangle$ be the vacuum of Yang Mills theory. Then according to Wilson if the theory confines, then for large loops $\gamma$

$$\langle 0 | W_A(\gamma) | 0 \rangle \to e^{-A(\gamma)}$$

where $A(\gamma)$ is the area enclosed by the loop $\gamma$. Now in Abelian gauge theory we have

$$W_A(\gamma) = e^{(2\pi i / g) \Phi_M(\gamma)}$$

where $g \equiv 2\pi / e$

and $\Phi_M(\gamma)$ is the magnetic flux, which passes through $\gamma$. 't Hooft then used equation (6) to define gauge invariant magnetic flux in Yang Mills theory. We call the constant $g = 2\pi / e$ the magnetic color coupling constant since in SU(N) Yang Mills theory 't Hooft showed that $\Phi_M(\gamma)$ is quantized in integral units (Mod N) of $g / N$. 
t'Hoof\!t then defined a second gauge invariant operator $W_C(\mathcal{L})$ by fixing its commutation relations with $W_A(\mathcal{L})$ such that,

$$W_C(\mathcal{L}) = e^{(2\pi i/e)\Phi_E(\mathcal{L})}$$  \hspace{1cm} (8)$$

where $\Phi_E(\mathcal{L})$ is the electric color flux through $\mathcal{L}$. In Abelian theory $W_C(\mathcal{L})$ is immediately expressible in terms of electric vector potentials $C_\mu$ defined by writing the solution of Maxwell's equations

$$\nabla \cdot \mathcal{D} = 0, \quad \nabla \times \mathcal{H} = \frac{\partial \mathcal{D}}{\partial t} + \mathcal{J},$$

as

$$\mathcal{D} = - \nabla \times \mathcal{C}, \quad \mathcal{H} = -\frac{\partial \mathcal{C}}{\partial t} - \nabla \cdot \mathcal{C}_0  \hspace{1cm} (9)$$

$$\nabla \cdot \mathcal{D} = \mathcal{D} = 0, \quad \nabla \times \mathcal{H} = \frac{\partial \mathcal{D}}{\partial t} + \mathcal{J},$$

as

$$\mathcal{D} = - \nabla \times \mathcal{C}, \quad \mathcal{H} = -\frac{\partial \mathcal{C}}{\partial t} - \nabla \cdot \mathcal{C}_0  \hspace{1cm} (10)$$

In this case $\Phi_E(\mathcal{L}) = \int d^2 \mathbf{s} \cdot \mathbf{D} = \int d^2 x \mathcal{E}_\mu(x)$ and we have

$$W_C(\mathcal{L}) = e^{i g f \int d^2 x \mathcal{C}_\mu(x)}  \hspace{1cm} (11)$$

Mandelstam\(^2\) then used eq. (11) to define $\mathcal{C}_\mu(x)$ in non Abelian gauge theory. The gauge invariance of $W_C(\mathcal{L})$ then implies that $\mathcal{C}_\mu$ transforms as

$$\mathcal{C}_\mu \rightarrow \mathcal{C}_\mu' = \Omega^{-1} \mathcal{C}_\mu \Omega + (i/g) \Omega^{-1} \partial_\mu \Omega \hspace{1cm} (12)$$

Mandelstam called the group of transformations (12) the magnetic color gauge group in contrast to the usual electric color gauge group of transformations\(^3\). Note that the coupling constant $g$ determining the magnetic color gauge group is inversely proportional to the Yang Mills coupling constant $e$, eq.(7). Although Yang Mills theory is invariant under the transformation (12), there is no simple relation between $A_\mu(x)$ and $\mathcal{C}_\mu(x)$ as in Abelian gauge theory. Thus the Yang Mills Lagrangian is invariant under the transformations of the magnetic group, but its explicit dependence upon the fields $\mathcal{C}_\mu$ is in general unknown.

\'t Hooft\(^1\) using his definition of $W_C(\mathcal{L})$ via its commutation relations with $W_A(\mathcal{L})$ showed that in a confining theory where $W_A(\mathcal{L})$ satisfies the area law (5), then $W_C(\mathcal{L})$ necessarily satisfies a perimeter law for large loops,

$$W_C(\mathcal{L}) \rightarrow e^{-\text{perimeter}(\mathcal{L})}  \hspace{1cm} (13)$$

Using eq. (11) expressing $W_C(\mathcal{L})$ in terms of electric vector potentials we conclude that if Yang Mills theory is in the confining phase, then the dual potentials $\mathcal{C}_\mu$ are weakly coupled at long distances. Note that his result is a precise way of saying that the vacuum of a confining theory behaves like a dual superconductor. For, in an ordinary superconductor there are vortices of quantized magnetic flux $\Phi_M = \int \mathbf{B} \cdot d\mathbf{s}$ which produces a linear potential between magnetic...
monopoles. Outside the magnetic vortex as we penetrate a distance \( r \) into the interior of the superconductor the magnetic field dies off exponentially, \( |\hat{B}| \sim e^{-mr} \) (the Meissner effect). Then, since \( \hat{B} = \hat{\nu} \times \hat{A} \), the potential \( \hat{A} \) is weakly coupled at long distances in an ordinary superconductor. On the other hand in a dual superconductor there is a linear potential between quarks and tubes of quantized electric flux \( \Phi_E = \int \hat{D} \cdot d\hat{s} \). Outside the flux tube the field \( \hat{D} \) dies off exponentially and since \( \hat{D} = -\hat{\nu} \times \hat{C} \), we conclude that \( \hat{C} \) is weakly coupled at long distance.

The electric vector potentials \( C_{\mu} \) being weakly coupled at long distances are thus appropriate variables for describing long distance Yang Mills theory. However it is not clear now to use them since the original Yang Mills Lagrangian is expressed in terms of the usual magnetic vector potentials \( A_{\mu} \). We thus begin with dynamics of Yang Mills theory described by the Schwinger Dyson equations for the gluon propagator \( \Delta^{ab}_{\mu}(x,y) \) defined as

\[
\Delta^{ab}_{\mu}(x,y) = \langle 0 | A^a_{\mu}(x) A^b_{\nu}(y) | 0 \rangle
\]

and the multi-gluon vertices \( \Gamma_n \). In momentum space the gluon propagator has the form

\[
\Delta^{A}_{\mu}(q) = \frac{1}{q^2} \varepsilon(q)
\]

(14)

where the dielectric constant \( \varepsilon(q) \) is determined by the vacuum polarization Dyson equation

\[
\varepsilon(q) = 1 + \frac{\Delta^{A}}{q^2}
\]

(15)

The triple gluon vertex \( \Gamma_3 \) is in turn given in terms of \( \Gamma_4 \), the four gluon vertex, etc. There result an infinite set of coupled equations which, when expanded in powers of the Yang Mills coupling constant \( e \), yields the usual perturbation solution. This solution satisfies order by order the requirements of gauge invariance but because of asymptotic freedom ceases to be valid at small \( q^2 \).

We seek a self consistent non perturbative starting point for an alternate expansion of the Dyson equations which also satisfies the requirements of gauge invariance. To do this \(^3\) we replace \( \Gamma_3 \) by \( \Gamma_3(\varepsilon) \) in the Dyson eq. (15) for \( \varepsilon \) where \( \Gamma_3(\varepsilon) \) is the singularity free solution of the Ward identity.
\[
\Gamma_{3,\mu,\nu,\lambda}(q,k,k') = \epsilon(k)(k^2 g_{\nu\lambda} - k'_{\nu} k'_{\lambda}) - \epsilon(k')(k'^2 g_{\nu\lambda} - k_{\nu} k'_{\lambda}),
\]
given by
\[
\tau_{3,\mu,\nu,\lambda}(q,k,k') = \frac{\epsilon(k)k_{\nu}\epsilon(q)k'_{\nu} - (\epsilon(k) - \epsilon(q))}{k^2 - q^2} (k^2 - q^2)
\]
+ cyclic permutations.

Eq. (15) then becomes an integral equation of the form:
\[
\epsilon(q) = 1 + e^2 \int k k' (k^2 - q^2) \epsilon(q) + e^2 \epsilon(\alpha) \int k k' (k^2 - q^2) \epsilon(k) \epsilon(k')
\]
where \(k\) and \(L\) are kinematic factors.

Setting \(q = 1\) on the right hand side of eq. 16 yields the usual perturbative solution with its inconsistent low momentum behavior. One can show that the only possible self-consistent low momentum behavior of \(\epsilon(q)\) compatible with eq. 18 is
\[
\epsilon(q) \rightarrow -q^2/M^2, \quad \text{as } q^2 \rightarrow 0
\]
where \(M^2\) is an undetermined scale parameter. We have solved eq. (18) numerically and have obtained a self consistent solution for all \(q^2\) having the low \(q^2\) behavior given by eq. (19) (see Fig. 1).

![Figure 1](image)

The solution of the integral equation for \(\epsilon(q)\). The solid curve is the input \(\epsilon(q)\). The circles are the output \(\epsilon(q)\) calculated from the right hand side of the integral equation (18).

Denoting by \(\Delta^0_A(q)\) the limiting low momentum behavior of the gluon propagator obtained from the solution of eq. (18), we obtain from eq.'s (14) and (19).
\[
\Delta^0_A(q) = M^2/(q^2)^2
\]
Using \( \Delta_0^A(q) \) as a starting point we can develop a new expansion of the Dyson equations which order by order satisfy all the requirements of gauge invariance (all the Ward identities).

Because of the factor \( M^2 \) appearing in the numerator of \( \Delta_0^A \) all the vertices \( \Gamma \) of this new expansion will then have the form

\[
\Gamma \approx \sum_n (e^2M^2)^n I_n
\]  

(21)

where the \( I_n \) is an \( n \) dimensional integral over \( \Delta_0^A \). Since the effective coupling constant in eq. (21) has the dimensions of \( M^2 \), the high momentum contributions to \( I_n \) can be neglected for calculating \( \Gamma \) at low \( q^2 \) and we are left with a pure low momentum problem with a dimensionful coupling constant \( e^2M^2 \), whose value reflects the effect of high momentum on low momentum. On the other hand, for the same reason, the low momentum contributions to eq. (21) are extremely divergent and thus eq. (21) does not provide a viable expression for calculating \( \Gamma \).

We will see that this is a reflection of the fact that the \( A_\mu \) are not appropriate variables for calculating the long distance behavior of Yang Mills theory. We will now show how to reexpress the physics described by eq. (20) and (21) in terms of electric vector potentials \( C_\mu \).

We will end up with a low energy theory described by an effective Lagrangian expressed in terms of electric vector potentials \( C_\mu \) which yields a confining theory and in particular the QCD string.

Let us first write the zero order theory, that is an Abelian theory having a propagator \( \Delta_0^A \), in terms of electric vector potentials \( C_\mu \). An Abelian theory of a relativistic dielectric medium with dielectric constant \( \varepsilon(q) \) is described by an \( A_\mu \) propagator of the form, eq. (14), or by a \( C_\mu \) propagator \( \Delta_0^C, C_{x,y} = \langle 0|C_\mu(x)C_\nu(y)|0\rangle \), which in momentum space has the form

\[
\Delta_0^C(q) = 1/q^2\mu(q), \quad \text{with } \mu(q^2) = 1/\varepsilon(q)
\]  

(22)

Thus for a medium with dielectric constant \( \varepsilon = -q^2/M^2 \), \( \mu = -M^2/q^2 \), and hence \( \Delta_0^C(q) \) eq. (22) has no singularity at \( q^2 = 0 \). (Just as \( \Delta_0^A(q) \) has no singularity at \( q^2 = 0 \) in an ordinary superconductor). To give a momentum dependence to \( \Delta_0^C(q) \) we must include the next term in the low momentum expansion of \( \mu(q) \). Thus we take

\[
\mu(q) = -M^2/q^2 + 1/f^2
\]  

(23)

where \( f^2 \) is a dimensionless parameter which cannot be determined from our solution of the Dyson equations since it is sensitive to our choice of \( \Gamma_3 \). Inserting eq. (23) into (22) yields
the $C_\mu$ propagator $\Delta^0 C$ given by:

$$\Delta^0_C(q) = \frac{f^2}{q^2 - M_f^2}, \quad (24)$$

where $M_f = fM$. Thus in terms of the $C_\mu$ variables the starting point is an Abelian theory with a massive $C_\mu$ propagator. The singular $\Delta^0_A(q)$ and the smooth $\Delta^0_C(q)$ describe the same physics.

Before describing the interactions in terms of the $C_\mu$ variables, let us note some features of this zero order Abelian theory. We write the constitutive equations $\mathcal{D} = \varepsilon \mathcal{E}$ and $\mathcal{B} = \mu \mathcal{H}$ in the coordinate representation with $\varepsilon = \mu^{-1} = 3^2/\mathcal{M}^2$ (including the $1/r^2$ contribution to $\varepsilon$ will not change this discussion essentially). Using eqs. (10) for $\mathcal{D}$ and $\mathcal{H}$ we obtain

$$-\mathcal{\nabla} \times \mathcal{E} = \frac{1}{\mathcal{M}^2} \partial^2 \mathcal{E}, \quad -\partial \mathcal{E} / \partial t - \mathcal{\nabla} \cdot \mathcal{B} = \frac{1}{\mathcal{M}^2} \partial^2 \mathcal{B} \quad (25)$$

We thus see that $\mathcal{E}$ and $\mathcal{B}$ are not determined in terms of the potentials $C_\mu$. We can always add to any solution $\mathcal{E}$ and $\mathcal{B}$ of eqs. (25), solutions of the homogeneous equations $\partial^2 \mathcal{E} = 0$, $-\partial^2 \mathcal{B} = 0$. Thus the fields $\mathcal{E}$ and $\mathcal{B}$ are independent dynamical variables. The Abelian zero order theory is thus described by 10 variables $C_\mu, \mathcal{E}$, and $\mathcal{B}$. The variables $\mathcal{E}$ and $\mathcal{B}$ will play the role of Higgs fields when we include interactions. We see here from purely kinematic reasons that such fields are necessarily present even in the non interacting theory. Finally we note the solutions of the homogeneous eq's for $\mathcal{E}$ and $\mathcal{B}$ represent massless excitations which accompany the excitation of mass $M_f$. If Yang Mills theory confines, the interactions must remove these residual massless excitations present in the zero order theory.

We now express the interactions as represented by eq. (21) in terms of the $C_\mu$ variables. This cannot be done directly, because in non Abelian theory there is no direct correspondence between the variables $C_\mu$ and $A_\mu$. Furthermore the expression for the interactions in terms of the $A_\mu$ variables is given by a divergent sum eq. (21). However we know that each term in this series is invariant under non abelian gauge transformations and hence the interactions in the $C_\mu$ description must be invariant under non abelian gauge transformations (12) of the magnetic gauge group. We also know that only long distances i.e., low momenta are involved. We will see that these facts will determine the essential features of the interactions.

We thus seek the Lagrangian, $L_{\text{eff}}(C_\mu, \mathcal{E}, \mathcal{B})$, which represents the long distance interacting theory in terms of the $C_\mu$ variables. Let us denote $L_{\text{eff}}^{(0)}(C_\mu, \mathcal{E}, \mathcal{B})$ the Lagrangian
generating the equations of the non interacting theory such eqs (25). We then must have

\[ \text{Leff}(\mathbf{C}, \mathbf{E}, \mathbf{B}) = L^{(0)}_{\text{eff}}(\mathbf{C}, \mathbf{E}, \mathbf{B}) \quad \text{when } \mathbf{C}, \mathbf{E}, \mathbf{B} \text{ are Abelian.} \]  

(26)

This means all the interactions of the effective long distance theory arise from the non Abelian nature of the fields just as in the original Yang Mills theory.

(II) \text{Leff must be invariant under the transformations (12) of the magnetic color gauge group and hence its dependence upon } \mathbf{C} \text{ must enter in one of the following combinations}

\[ g_{\mu\nu} = \delta_{\mu\nu} C_{\nu} - i g [C_{\mu}, C_{\nu}], \quad \text{or} \quad D_{\mu} f = \partial_{\mu} f - i g [C_{\mu}, f], \]  

(27)

Since \( g = 2\pi/e \), we see immediately that the expansion parameter of the theory is \( (2\pi/e)^2 \) or \( 1/\alpha_s \); i.e. the perturbation expansion will be valid when \( \alpha_s \) is large, i.e. at low momentum.

(III) \text{Since we integrate only over small momentum } q^2 < M_f^2 \text{ we need include only a minimal number of derivations in Leff. Higher derivative terms will give contributions of the order } \( (p/M_f)^2 \) \text{ to processes at momentum } p \text{ and will not be important for momentum } p \ll M_f, \text{ for which the theory is applicable.}

The derivatives appearing in \text{Leff are thus determined by } L^{(0)}_{\text{eff}}. \text{ Then the dependence of } \text{Leff on } \mathbf{C} \text{ is uniquely determined by (II), that is, by combining knowledge of } L^{(0)}_{\text{eff}} \text{ and gauge invariance. The only feature of } \text{Leff which is not completely determined by the above considerations is its dependence upon the variables } \mathbf{E} \text{ and } \mathbf{B}. \text{ \text{Leff} thus contains a gauge invariant function } \mathbf{W} \text{ such that}

\begin{align*}
\text{a) } & \mathbf{W}(\mathbf{E}, \mathbf{B}) \text{ contains no derivatives,} \\
\text{b) } & \mathbf{W}(\mathbf{E}, \mathbf{B}) = 0, \text{ if } \mathbf{E} \text{ and } \mathbf{B} \text{ is abelian.} \end{align*}  

(28)

To understand the meaning of \( \mathbf{W} \) consider the Yang Mills vacuum in the absence of quantum fluctuations. Then

\[ \mathbf{C} = 0 \text{ and hence,} \]

\[ \text{Leff} = \mathbf{W}(\mathbf{E}, \mathbf{B}). \]  

(29)

The vacuum is then described by the classical equations of motion

\[ 0 = \delta \text{Leff}/\delta \mathbf{B} = \delta \mathbf{W}/\delta \mathbf{B} \]

\[ 0 = \delta \text{Leff}/\delta \mathbf{E} = \delta \mathbf{W}/\delta \mathbf{E} \]

One can show that the lowest energy stable solution will always have \( \mathbf{E} = 0 \) so let us concentrate on the dependence of \( \mathbf{W} \) upon \( \mathbf{B} \). Suppose \( \mathbf{B} \) has the form given in Figure (2) i.e. it has a non trivial minimum for some non vanishing value of \( B_0 \), determined by the equation

\[ \delta \mathbf{W}/\delta \mathbf{B} = 0. \]  

(30)
One can then show that the magnetic condensate, $\hat{B}_0$

(1) removes the massless excitations present in the zero order Abelian theory,

(2) produces a magnetic pressure which confines electric flux.

$W(\hat{E}, \hat{B})$ represents the part of long distance dynamics which is not determined by combining our solution of the Dyson equations with invariance under the transformations of the magnetic gauge group. The important feature of $W$ is that it has the shape shown in fig. 2, yielding $\hat{B}_0 \neq 0$.

We now write the complete form of the effective Lagrangian $L_{\text{eff}}$ describing long distance Yang Mills theory. It is convenient to combine the fields $\hat{E}$ and $\hat{B}$ into an antisymmetric tensor $\tilde{F}_{\mu \nu}$ defined as

$$\tilde{F}_{\mu k} = -B_k, \quad \tilde{F}_{ij} = -\varepsilon_{ijk}E_k.$$  \hspace{1cm} (31)

Then

$$L_{\text{eff}} = 2\text{trace} \left( \frac{1}{M^2} (D_\mu \tilde{F}_{\alpha \beta}) (D_\mu \tilde{F}_{\alpha \beta}) - \frac{1}{f^2} (\tilde{F}_{\mu \nu})^2 - W(\tilde{F}) \right).$$  \hspace{1cm} (32)

It is clear that $L_{\text{eff}}$ is invariant under the non Abelian gauge transformation (12) of the potentials $C_{\alpha \beta}$ accompanied by the transformation

$$\tilde{F}_{\alpha \beta} \rightarrow \Omega^{-1} \tilde{F}_{\alpha \beta} \Omega.$$

(33)

of the fields $\tilde{F}_{\alpha \beta}$ to ensure that it vanishes when $\tilde{F}_{\alpha \beta}$ is Abelian. $W(F)$ must be a Lorentz invariant gauge invariant function of commutators like $[\tilde{F}_{\alpha \beta}, \tilde{F}_{\mu \nu}], [\tilde{F}_{\alpha \beta}, [\tilde{F}_{\mu \nu}, \tilde{F}_{\gamma \delta}]]$ to ensure that it vanishes when $\tilde{F}_{\alpha \beta}$ is Abelian. It must of course have the shape shown in Fig. 2.

To understand further the meaning and origin of $L_{\text{eff}}$, we can expand $L_{\text{eff}}$ in powers of $C_{\mu}$ after eliminating $\tilde{F}_{\mu \nu}$ via its equation of motion. The coefficient of $(C_{\mu})^n$ in this expansion gives the zero order multi-gluon $C_{\mu}$ vertices $\Gamma_n^{C_{\mu}}$. Because $L_{\text{eff}}$ is invariant under the...
transformations of the magnetic gauge group these vertices will satisfy the same Ward identities as the original $A_\mu$ vertices, $\Gamma_n$. In particular the Ward identity for $\Gamma_3^C$ has the form (16) with $\epsilon$ replaced by $\mu$. Replacing $\epsilon$ by $\mu = -M^2/q^2 + 1/f^2$ in the solution (17) yields an expression for zero order $\Gamma_3^C$ which agrees with the result obtained from the expansion of $L_{\text{eff}}$. Similarly all the zero order $C_\mu$ vertices $\Gamma_\mu^C$ derived for $L_{\text{eff}}$ can be obtained by solving the infinite set of Ward identities beginning with $\Gamma_2^C = \mu = -M^2/q^2 + 1/f^2$.

For the color group $SU(2)$ we have found solutions of the classical equations of motion

$$\frac{\delta L_{\text{eff}}}{\delta C_\mu} = \frac{\delta L_{\text{eff}}}{\delta \mu} = 0$$

representing a confined tube of electric flux. For $SU(N)$ it is easy to show that the electric flux is quantized in units of $e/N$ where $e$ is the Yang Mills coupling constant. Our $SU(2)$ flux tube contains flux $e/2$ and has the following color structure.

$$E_3 = E_3 T_3, \quad B = B T_1 + B T_2, \quad C = C_3 T_3, \quad D = D_3 T_3$$

We choose the $z$ axis to be the axis of the flux tube and introduce cylindrical coordinates $(r, \phi, z)$. Then

$$E_3 = E(r) \delta_2, \quad B = B_1 (r) \delta_r, \quad B_2 = B_2 (r) \delta_\phi, \quad C = C(r) \delta_\phi, \quad D = D_3 (r) \delta_z$$

The electric field thus lies along the 3 direction in color space and the $z$ axis in ordinary space; the magnetic field on the other hand lies along the 1 and 2 directions in color space and along the $r$ and $\phi$ direction in ordinary space. A cross section of this flux tube is shown in Fig. 3:

![Cross section of the flux tube.](image)

The dots represent the electric field lines which point along the $z$ axis.
The electric flux is contained in a tube of width \(1/M_f\). The field \(D_3(r)\) dies off exponentially for distances \(r > M_f\). The magnetic fields \(B_1\) and \(B_2\) approach at large distances a common value \(|\vec{B}_0|\) determined by (30). Inside the flux tube they fall off as shown in Fig. 4. Thus we see that the magnetic field \(\vec{B}\) plays the role of a Higgs Field.

![Figure 4](image)

The magnetic field \(B_1(r)\) and the electric field \(D_3(r)\) as functions of the distance \(r\) from the center of the flux tube.

We can compare this solution to a Nielsen-Olesen magnetic vortex in which quantized Abelian magnetic flux is confined by an electrically charged scalar Higgs field \(\phi(x)\). Correspondingly, the above static cylindrically symmetric solution of our equations represents a unit \(e/2\) of quantized color electric flux confined by a non Abelian magnetic field \(\vec{B}\), which assumes a non vanishing vacuum expectation value \(\vec{B}_0\) at large distances. The QCD flux tube can thus be regarded as a dual Nielsen-Olesen vortex, obtained from Yang Mills theory without scalar fields.

We next want to carry out semi classical quantization about this classical solution. We note that the coupling constant of the theory is \(g^2/M_f^4\) since by redefining variables, we can write \(L_{\text{eff}}\) in a form where the coupling constant appears as a factor \(1/(g^2/M_f^4)\). Semi classical quantization then leads to an expansion in \(g^2/M_f^4\). In leading order we have \(g^2/M_f^4 \to 0\) and the flux tube becomes the QCD string. Its quantization leads as usual to linear Regge trajectories. The details of this quantization has not yet been worked out, and it is not clear which version of the quantum string will result. The corresponding problem of the semi classical quantization of the Nielsen Olesen vortex\(^5\) was carried out by Gervais and Sakita\(^6\) and their methods could be used to obtain a similar expansion about our QCD flux tube solutions.
We are now working on the simpler problem of the quantization of $L_{\text{eff}}$ about the vacuum solution, $\hat{B} = \hat{B}_0$. The resulting quantum fluctuations will restore the rotational and color symmetry of the vacuum and only massive physical excitations will survive. Perturbation theory about the physical vacuum will then be an expansion in powers of $(g/M_f^2)^2 = (1/\alpha_s(q^2))(1/M_f^4)$ with a propagator $\Delta_C^{0}(q^2) = 1/(q^2 - M_f^2)$, and integrations which go over only long distances, i.e. $r > 1/M_f$. The effect of short distances are contained in the parameter $M_f$ and in the structure of $L_{\text{eff}}$. Since the expansion parameter is $1/\alpha_s(q^2)$, such an expansion could serve as a possible basis for a perturbative long distance QCD phenomenology.

As a final remark we note that the full $\zeta$ propagator, $\Delta_C(q^2)$ will have the same long distance behaviour as $\Delta_C^{0}(q^2)$ except for renormalization effects, since the perturbative corrections have no long distance singularities. Thus the full $\zeta$ propagator will have the long distance form,

\begin{equation}
\Delta_C = \frac{1}{q^2 - M_f^2}.
\end{equation}

However since in non Abelian theory there is no simple relation between $\Delta_A$ and $\Delta_C$ we cannot infer any information from eq. 37 about the long distance behaviour of $\Delta_A(q^2)$. In particular we would have no reason to think that $\Delta_A(q^2) \rightarrow \Delta_A^{0}(q^2) = M_f^2/q^4$ at low $q^2$.

We have considered above pure Yang-Mills theory without quarks. To calculate the potential between heavy quarks at finite distances we must introduce the quark gluon interactions explicitly. However the calculation of the resulting quark-antiquark potential becomes tractable only when non Abelian contributions are neglected everywhere except at large distance where they are essential for giving the vacuum energy density $W(\hat{B}_0)$ a non vanishing value. When this is done, comparison with phenomenological $\bar{q}q$ potentials then fixes $M_f = 1.26$ GeV, and $W(\hat{B}_0) = -(178 \text{ MeV})^4$.

References.
2) S. Mandelstam, Phys. Rev. D19, 2391 (1979)
Abstract:

An introduction to the Polyakov string theory is made. Then, we show how to derive this string theory from $\Phi^3$ field theory, using exact propagators. We obtain a version of the Polyakov string which has a zero-critical dimension in agreement with the previously known unphysicalness of the $\Phi^3$ field theory.
We would like to present some results concerning the relations between field theory and string theory. Strings are now again fashionable. They promise to be a Theory Of Everything, from which at low energies, i.e. accessible to present experiments, one would recover our field interpretation of data. Then, the interest is focused on the derivation of fields from strings:

\[
\text{STRINGS} \Rightarrow \text{FIELDS}
\]

Here, on the contrary, we consider strings as a mean to calculate field theory, i.e. we are interested in the reverse derivation:

\[
\text{FIELDS} \Rightarrow \text{STRINGS}
\]

an approach which aims at being useful to calculate the S-matrix for present-day experimental processes. We intend to briefly discuss the following:

1 - An Introduction to the POLYAKOV string
2 - How to rigorously derive the Polyakov string from \(\Phi^3\) field theory
3 - The (zero)-critical dimension of \(\Phi^3\) field theory.

1 - The Polyakov string theory

Old string models, such as the Veneziano bosonic string or the Neveu-Schwartz-Ramond fermionic string are physical in only 26 or 10 dimensions, which renders them inappropriate in our "low energy" 4-dimensional world. However, in 1981 Polyakov\(^1\) found generalizations for them which are physical for \(d \leq 26\) or \(d \leq 10\), and are therefore much more attractive.

The Polyakov string is described by a functional integral

\[
\mathcal{Z}_p = \sum_{\text{topologies}} \int [d\,g_{ab}(\xi)] \int [dx^\mu(\xi)] \exp (-L)
\]

where \(L\) is an action defined on a 2-dimensional disk \(D\) on which a \((\xi^1, \xi^2)\) coordinate system is drawn. The sum over topologies means that holes can be cut in \(D\) and holes joined together through handles. The \(x^\mu(\xi)\)'s \((\mu = 1, ..., d)\) describe the embedding of \(D\) in a \(d\)-dimensional space. The \(g_{ab}(\xi)\) variables define a metric on \(D\) and are independent of the \(x^\mu\)'s. The action is the most general one invariant under general coordinate transformations and renormalizable:

\[
L(x^\mu, g_{ab}) = A \left(1/2\right) \int_D d^2\xi \sqrt{g} g^{ab} \partial_a x^\mu \partial_b x^\nu + B \int_D d^2\xi \sqrt{g} \left(C (1/4\pi)\right) g^{2\xi} \sqrt{g} R
\]
for the bosonic string. (The "old" Veneziano string is recovered by taking $g_{ab} = 6_{ab}$). In the following we only consider the first term in eq. (2). What is interesting about $Z_P$ is that one can integrate over most of its variables and obtain:

$$Z_P = N \sum \int [d \sigma(\xi)] \exp (-\mathcal{L})$$

where $d \sigma(\xi)$ expresses a dilatation of the metric, i.e. the conformal degree of freedom and is an effective lagrangian

$$\mathcal{L} = [(26 - d)/2\pi] \int d^2 \xi (\partial \sigma)^2 + \psi_0^2 \int d^2 \xi \exp(2\sigma)$$

(with $\sqrt{g} = \exp(2\sigma)$)

where the part of the first term with the factor 26 is coming from the integration over the two non-conformal degrees of freedom of the metric $g_{ab}(\xi)$. For $d > 26$ the theory explodes, giving 26 as the critical dimension of the full Polyakov string theory. We shall see below that the $\phi^3$ field theory gives a restricted version of the Polyakov string where the factor 26 has disappeared in the first term.

2 - Derivation of the Polyakov string from $\phi^3$ field theory

Let us consider the propagator expression for a scalar field in Feynman's representation

$$\int_0^{\infty} d\xi \exp [-\alpha \Delta x^\mu \Delta x^\nu]$$

where $x^\mu$ are momentum-space variables which sit on the faces of a given planar Feynman graph drawn on $D$. One can derive a discretized version of the first term of the lagrangian displayed in (2) by making the identification:

$$(1/2) \partial^2 \xi, \sqrt{g} \leftrightarrow \alpha$$

$$g_{ab}, \partial x^a \partial x^b \leftrightarrow \Delta x^\mu \Delta x^\nu$$

where an infinitesimal area $\partial^2 \xi$, is associated with each edge. $\alpha$ can then be interpreted as a $d$-dimensional metric used to build a 2-dimensional one through $\sqrt{g}$. The integration
measure for $g_{ab}$ is now split into

$$[d g_{ab}] = [d \sqrt{g}] [d \hat{g}_{ab}]$$  (7)

where $\sqrt{g}$ (or $\alpha$) plays the role of the conformal degree of freedom. $[d \hat{g}_{ab}(\xi)]$ takes care of the non-conformal degrees of freedom and the integration over $\hat{g}_{ab}(\xi)$ is obtained through a summation over all possible graphs, each graph contributing to a fixed value $\hat{g}_{ab}(\xi)$ on each surface element $\partial^2 \xi$ of $D$. We will see in the next section how to each edge of the Feynman graph (or more precisely of its dual which has its vertices in the center of the faces of the original Feynman graph and with edges across the edges of the Feynman graph) a particular value of $\hat{g}_{ab}(\xi)$ can be computed.

If all values $\hat{g}_{ab}$ are given the same weight, i.e. if the same number of graphs contributes to each value of $\hat{g}_{ab}$, then we have recovered the entire integration measure for $Z_P$:

$$[d \sqrt{g}(\xi)] [d \hat{g}_{ab}(\xi)] [d \chi(\xi)]$$  (8)

and the complete Polyakov bosonic string will have been shown to be equivalent to $\Phi^3$ field theory.

However, we will see that is not the case for $\Phi^3$, because the integration over the non-conformal degrees of freedom will be frozen.

3 - The critical dimension of $\Phi^3$

We now argue that there is preferred value for $\hat{g}_{ab}$ when summing over all planar graphs. This preferred value will introduce a delta-function $\delta(\hat{g}_{ab} - \bar{g}_{ab})$ in the measure which destroys the integration over the non-conformal degrees of freedom of the metric, so that the term with a factor $26$ in the effective lagrangian (4) vanishes. Instead, we will have

$$L_{\Phi^3} = \left[-d/24\pi\right] \int d^2 \xi (\partial \Phi)^2 + \mu^2 \int d^2 \xi \exp(2\Phi)$$  (9)

for the effective lagrangian, giving a zero-critical dimension for the $\Phi^3$ field theory, in
agreement with the known unphysicalness of $\Phi^3$ for positive dimensions$^2$).

This property can be proved by using an algorithm$^3$ that we devised for the systematic construction of all planar graphs, and which shows that most of the propagators can be put parallel to the axis $\xi^1 = 0$ or $\xi^2 = 0$ (see figure 1 below).

All graphs can be constructed by piling up rows of rectangular cells. Those cells contain areas (shown hatched in figure 1) on which the construction of sub-diagrams is performed in the same way as for the entire graph, but at a smaller "scale". $g_{ab}$ is defined through

$$\sum_{\mathbf{P} \in \mathbf{d}^2\xi} \Delta \xi^a \Delta \xi^b \cdot (1/2) g_{ab} C(\xi)$$

\[ (10) \]

($\mathbf{P} =$ propagator or edge)

where $C(\xi)$ is independent of the graph for every $d^2\xi$. It can be easily proved that the number of graphs having

$$\hat{g}_{11}/\hat{g}_{22} \to 0, \quad \hat{g}_{12} = 0$$

is infinite compared to those having other values of $\hat{g}_{ab}$. In fact, $\hat{g}_{11}/\hat{g}_{22} \to 0$ because to each graph having a finite number of edges parallel to $\xi^2 = 0$ in figure 1 in some area $d^2\xi$, one can associate an infinite number of graphs having an infinite number of edges parallel to $\xi^2 = 0$ on $d^2\xi$. Also, $\hat{g}_{12} = 0$ because in the same way, almost all edges are parallel either to the $\xi^1 = 0$ or to $\xi^2 = 0$ axis. This fixes completely $\hat{g}_{ab}$, which has only two degrees of freedom, by providing a delta function in the integration measure.

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3) R. Hong Tuan, LPTHE Orsay preprints 85/4 and 85/43.
ABSTRACT

Difficulties concerning renormalization of Yang-Mills theories in light cone gauge are pointed out. They arise in loop calculations and may affect even gauge invariant quantities.
The light cone gauge, namely a gauge in which the condition
\[ n^\mu A^\mu = 0 \]  \hspace{1cm} (1)
with \( n^2=0 \) is imposed on the potentials \( A^\mu \), describing Yang-Mills fields, has
been often considered in recent years in applications to phenomenology as well
as in more formal treatments. Its main virtue is the possibility of an intuiti-
ve partonic interpretation; after redundant degrees of freedom have been elimi-
nated, the number of independent fields equals the number of the "free" quanta
in the theory.

This circumstance makes the formulation particularly suitable in connec-
tion with supersymmetry 1).

The resulting theory is however quite singular; in particular some pa-
thologies occur which are briefly discussed in the sequel.

A quantization based on a null plane algebra of operators 2) leads to
the following free vector propagator
\[ D_{\mu\nu}(k) = \frac{1}{k^2+i\varepsilon} \left[ -g_{\mu\nu} + \frac{n^\mu k^\nu + n^\nu k^\mu}{nk} \right] \]  \hspace{1cm} (2)
where the "spurious" singularity \( \frac{1}{nk} \) is interpreted as a Cauchy principal va-

tue.

This prescription, when inserted in loop calculations, leads to inte-
grals whose convergence cannot be controlled by power counting. Indeed the Wick
rotation entails the contribution of extra spurious poles.

If, for instance, the vector self-energy is computed at one loop in
SUSY \( N=4 \), we get a double pole at \( \omega=2, \) \( 2\omega \) being the number of dimensions in the
dimensional regularizat
scheme, coming from the integral
\[ \int \frac{d^2\omega}{(2\pi)^2} \left\{ \frac{(2\pi)^{k+2\omega}}{k^2+i\varepsilon} \frac{1}{(p-k)^2+i\varepsilon} \frac{p}{nk} \right\} = \]  \hspace{1cm} (3)
which should instead converge by power counting.

As noticed in ref. 3), this singularity would lead to an inconsistency
for the \( \beta \)-function, when computed directly from the self-energy.

Besides, the failure of the power counting criterion prevents the usual
renormalization procedure. This result casts some doubts on previous calculations involving loops, even for gauge invariant quantities, as they are not based on a consistent subtraction scheme.

A new recipe for handling the "spurious" singularity has been recently proposed in ref. 1) and 4):

\[ \frac{1}{[nk]} = \frac{n_k^*}{nkn^*_k+i\varepsilon} \]  

(4)

where \( n^*_\mu \) is the parity transformed of \( n^\mu \). In ref. 5) it has been proven that this prescription necessarily follows from an equal time algebra of operators. It entails the presence of a "ghost" propagating along a generating line of the light-cone. This can be intuitively understood, as the prescription (4) is not real.

If the \( \beta \) function in SUSY N=4 is computed from the self-energy at one loop using the prescription (4), it is found to vanish in agreement with the result obtained in other gauges. Moreover the Wick rotation is possible and the power counting criterion for convergence is recovered.

Still some difficulties arise in pursuing a consistent renormalization program 6). They are rooted in the appearance of the vector \( n^*_\mu \) and of singularities with a non local character in the proper vertices, namely of poles at \( \omega=2 \) with residues which are not polynomials in the external momenta 4). We stress that this pathology is peculiar of the light-cone gauge; indeed one can prove that other algebraic gauges do not share it.

The main consequence is the proliferation of new tensorial structures with divergent coefficients, which are allowed by the Lee-Ward identities as well as by dimensional analysis.

Non local divergences would entail non local counterterms; plenty of them can be envisaged starting from the basic non local quantity \( (n^\mu)^{-1}nA \), which is homogeneous in \( n^\mu \) and dimensionless.

Nevertheless one can prove that non local singularities in primitively divergent diagrams always decouple in the Green's functions, being proportional to the gauge vector \( n^\mu \). Hence if one contents himself with making finite only the Green's functions, this can be done at one loop level by means of a linear local transformation on the fields, which respects the unitarity property.
of the theory 6), 7).

For instance in the Yang-Mills theory interacting with Dirac fermions, we can define renormalized quantities as

\[ A^{(o)}_{\mu}(x) = R \psi_{\nu} A^{(o)}_{\nu}(x) \]
\[ \psi^{(o)}(x) = R \psi(x) \]
\[ g_{0} = Z_{3}^{-\frac{1}{2}} g, \]
\[ m_{0} = m + M, \]

with

\[ R_{\mu} = Z_{3}^{\frac{1}{2}} \left[ g_{\mu} - \frac{1}{2} g_{3} \left[ 1 - \frac{1}{2} \left( 1 - Z_{2}^{-1} \right) \frac{\gamma_{3}^{*}}{2n^{*}} \right] \right] \gamma_{3} \]
\[ R = (Z_{2}Z_{3})^{-\frac{1}{2}} \left[ 1 - \frac{1}{2} \left( 1 - Z_{2}^{-1} \right) \frac{\gamma_{3}^{*}}{2n^{*}} \right]. \]

At one loop the choice

\[ Z_{3} = 1 + g^{2} \left( 1 - \frac{Z_{2}}{3} n_{j} \right) \frac{1}{16\pi^{2} \epsilon} \]
\[ \gamma_{3} = 1 + \frac{3g^{2}}{8\pi \epsilon}, \]
\[ Z_{2} = 1 + \frac{g^{2}}{12\pi \epsilon}, \]
\[ \gamma_{2} = 1 + \frac{g^{2}}{6\pi \epsilon}, \]
\[ M = - mg^{2}/4\pi \epsilon, \]
\[ \epsilon = 2 - \omega, \]

succeeds in making finite all the Green's functions.

The generalization at any order of the loop expansion along this line is still an open problem. To give an example of the difficulties which are involved, the fermion-fermion-boson vertex contains the kernel

\[ \frac{1}{nn^{*}} \left[ n^{* \nu} - \frac{n^{* \nu} q_{j} n}{n_{j}} \right] \]

with a divergent coefficient, which cannot obviously be controlled by the Lee identity

\[ q_{\mu}^{*} \Gamma_{\mu}^{\nu}(q+p,p) = g\Gamma^{\nu}(p+q) - g\Gamma^{\nu}(p). \]

Now at one loop the linear local transformation (5-12) which is completely determined by the requirement of making finite the boson and fermion propagators succeeds in making finite also the fermion-fermion-boson Green func-
tion thanks to a fine tuning in the weights of the counterterms. This has been verified by an explicit one loop calculation, but we do not know, unfortunately, any reason why it should occur in general. In those conditions it is clear that an inductive proof of renormalizability is still out of reach.

REFERENCES

STOCHASTIC QUANTIZATION OF FREE STRING FIELD THEORIES

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ABSTRACT

After an elementary review on stochastic quantization the main features of stochastic quantization of free string field theories are outlined.
In this talk I would like to address myself to non-experts both in the field of stochastic quantization and of string field theories. I would like to attract your interest for these quite interesting and fascinating subjects and hopefully will be able as well to create some elementary understanding of the basic concepts involved. Concerning the string field part of my talk I will mainly report on results obtained together with I. Bengtsson.

Starting off with a brief sketch of stochastic quantization I would like to recall the important analogy between Euclidean field theory and statistical mechanics. Put into simple words the Euclidean path integral density \( \exp( - \text{Euclidean action}) \) has a close relation to the Boltzmann distribution of a statistical system in equilibrium, implying that Euclidean Green functions may be interpreted as correlation functions of a statistical system in equilibrium.

Let us observe that in a computer simulation of Euclidean Green functions (as e.g. in Monte Carlo calculations) the equilibrium configurations, which are needed for the averaging procedure, can in fact be generated only after some elapse of computer (!) time. One may, however, take as well a different attitude and calculate the correlation functions while the system is still in nonequilibrium. The desired equilibrium values can then be extracted by a careful study of the time evolution of the correlation functions up to large computer times.

Abstracting these ideas, Parisi and Wu formulated the following concept of stochastic quantization:

1) Supplement the fields \( \phi(x) \) with an additional coordinate, the "fictitious" time

\[
\phi(x) \rightarrow \phi(x,t), \quad x = (x_0, x_1, x_2, x_3)
\]  

2) Demand that the fictitious time evolution of \( \phi \) is described by a stochastic differential equation that allows for relaxation to equilibrium, as e.g. the Langevin equation

\[
\frac{\delta S}{\delta \phi} + \frac{\partial \phi}{\partial t} = - \frac{\partial S}{\delta \phi} + \eta
\]
Here $\eta$ is a Gaussian white noise and $S$ the Euclidean action. It should be remarked that eq. (2) without the time derivative and the noise term is just the classical field equation.

3) Solve for $\phi_\eta = \phi_\eta (x,t)$ and define correlation functions $(F(\phi_\eta))_\eta$ by performing Gaussian averages over $\eta$. $F$ denotes an arbitrary functional of $\phi$.

4) Euclidean Green functions are obtained by taking the limit $t \rightarrow \infty$

$$
(F(\phi)) = \lim_{t \rightarrow \infty} (F(\phi_\eta))_\eta
$$

Several interesting features emerge from the above stated quantization scheme (for a forthcoming review see $^3$): Firstly it can be shown in various ways $^2$, $^4$ that the new method is equivalent to the conventional path integral quantization. Furthermore I would like to point out that for gauge theories no gauge fixing is required $^2$; the additional fictitious time co-ordinate allows for a new type of regularization scheme $^5$; there is an intrinsic connection to supersymmetry $^6$ and quantization is possible also directly in Minkowski space $^7$. Last but not least a new numerical scheme is available $^8$.

Let us now discuss string fields. It was one of the most fascinating innovations in field theory to enlarge the concept of a field that is defined on space-time points to being defined on strings. The consequences of such a generalization are far reaching and involved; this deserves a detailed discussion, which is not intended in this talk. Let me just point out that the string field, denoted by $|\psi\rangle$, contains ordinary fields of arbitrary high spin and mass. It is actually a vector in a specific Fock space spanned by certain harmonic oscillators. The gauge invariant classical action for free bosonic strings as given in $^9$ reads (see also $^{10}$)

$$
S = \int d^2z \langle\psi(x)| H - q T_3^{-1} q |\psi(x)\rangle
$$

where $H$, $q$ and $T_3$ are operators defined in terms of the above mentioned oscillators $^{11}$. Varying the action (4) with respect to $|\psi\rangle$, we obtain the field equation

$$
0 = (H - q T_3^{-1} q) |\psi(x)\rangle = H \rho |\psi(x)\rangle
$$
and it can be shown that $P$ is a projector, i.e. $P^2 = P$. According to the procedure of stochastic quantization outlined above we associate to (5) the following Langevin equation:

$$\frac{\partial}{\partial t} |\Psi(x,t)\rangle = - HP |\Psi(x,t)\rangle + |\eta(x,t)\rangle$$

(6)

The same projection operator $P$ as in the classical field equation appears now in the drift term of the stochastic differential equation, where also a noise string field $|\eta\rangle$ has been introduced. Eq. (6) constitutes the most natural extension of the stochastic quantization concept from ordinary field theory to string field theory; given the fact that $|\Psi\rangle$ and $|\eta\rangle$ carry no indices, (6) appears to be rather unique.

The solution of (6) can easily be found in analogy to the Maxwell field \cite{2} and is given in Fourier space by

$$|\Psi(k,t)\rangle = \int d\tau \left[ e^{-H(t-\tau)P} + (1-P) \right] |\eta(k,t)\rangle$$

(7)

Calculating a gauge invariant quantity as e.g. $\langle\langle \Psi | HP | \Psi \rangle\rangle$ we obtain from (7)

$$\langle\langle \Psi(k) | HP | \Psi(k') \rangle\rangle = \lim_{t \to \infty} \text{tr}[P(1-e^{-2Ht})] \delta(k-k') = \text{tr}P \delta(k^*k')$$

(8)

where the trace is defined with respect to a complete set of states in the Fock space. We want to recall that the stochastic scheme is perfectly consistent without having added gauge fixing terms or ghost fields to the action (4).

Another virtue of the stochastic quantization method is that it helps to overcome the so called indefiniteness problem. By this we mean that the Euclidean string field action, in common with ordinary field theories containing fields with spin $\geq 2$, is not positive definite, so that straightforward path integral quantization is not possible (for a discussion concerning gravity see \cite{12}). In the stochastic approach the indefiniteness problem manifests itself when one tries to give a probabilistic interpretation of the noise \cite{1}. A rather natural prescription of how to deal with this can be found and it is finally possible to extract a well defined path integral prescription for free string fields. It seems an exciting challenge to apply the stochastic method also for the interacting case.
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A COMMENT ON THE ORIGIN OF SPIN-DEPENDENT FORCES
IN HEAVY QUARKONIUM SYSTEMS

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Abstract

It is pointed out that in models of quark confinement the usual analysis
of spin dependent forces in terms of scalar or vector potentials may be in-
appropriate. The special rôle that the self energy plays in such models
and its consequences for heavy quarkonia are discussed.
Q̅Q bound states are usually discussed within the framework of potential models and their spin dependent forces are then related to the Lorentz character of the exchange forces, vector versus scalar. Self energies are usually not considered though, as I will discuss, in the situation of color confinement they cannot be separated from interaction energies. This feature has been discussed, or is at least transparent, within several confinement models. The fact that tree and one-loop contributions are connected by color confinement is on the one hand interesting from a theoretical point of view but has on the other hand also phenomenological consequences. The latter are most transparent for the case of heavy quarkonia, especially for the LS force.

The reason why we are intuitively induced to neglect self energies can be seen when considering the hydrogen atom. The electric field is obtained as a superposition of those of the proton and the electron, \( \vec{E}_1 \) and \( \vec{E}_2 \). The total electrical energy

\[
E = \frac{1}{\epsilon_0} \int \frac{3}{\pi} (E_1^2 + E_2^2 + 2E_1E_2)
\]

separates into proton and electron self energies and an interaction term. The proton self energy is included into the renormalized proton mass. The electron self energy is included in the electron energy

\[
E_{\text{kin}} = \frac{m + \delta m}{\sqrt{1 - v^2}}
\]

except for the off shell correction, the Lamb shift, which is small since the electron is almost on shell. The interaction term can be treated as a potential in the Dirac equation. The fact that the electron field moves along with the electron should lead to a boost of the electric field energy which is however essentially contained in the boost of the electron mass, eq. (2).

In a model of confinement - and I will take a flux tube picture as an example - there is no simple superposition principle as used in eq. (1). The facts that the fields obey boundary conditions and that the energy is obtained in the static approximation by minimizing with respect to the geometry of the boundary makes clearly a separation between interaction- and self energies impossible. It is therefore likewise impossible to incorporate the boosting of the moving field configuration into the quark masses.
Neglecting this boost may be a reasonable approximation for heavy quarkonia but leads to a wrong connection

\[ \alpha' = \frac{1}{4\sigma} \]

(3)

between Regge shape and string tension for light quarks.

A framework where the analysis of spin dependent forces for quarkonia can be done almost along conventional lines, i.e. using the Breit-Fermi approximation, but applied to both self and interaction energies, have been proposed in refs. 5), 11). One considers a version of the MIT bag model where the spherical approximation is used (realistic flux tubes are almost spherical), and where the radius \( R \) of the bag is determined in a self consistent way as for light quark bags. For a spherical cavity the confined gluon propagator with bag boundary conditions can be calculated explicitly, and the only place where the condition of vanishing color comes in is the \( s \) wave part of the Coulomb propagator. Explicitly

\[
G_c(x,x') = \frac{1}{4\pi|x-x'|} + \sum_{L=0}^{\infty} \frac{\ell+1}{\ell} \frac{(rr')^\ell}{4\pi R^{2\ell+1}} P_\ell(x \bar{x}')
\]

(4)

Obviously the boundary term for \( \ell = 0 \) is ill-defined, but if the propagator is connected on both ends to all color sources the numerator vanishes as well for colorless states. Still the \( s \) wave parts of self and interaction energy remain ambiguous and it is only the sum of (tree and 1-loop) graphs

\[
\begin{align*}
\text{Q} & \quad + \quad 2 \quad \text{Q} \\
\bar{Q} & \quad + \quad \text{Q} \quad \bar{Q}
\end{align*}
\]

which is well defined. In other gauges this "infrared" singularity appears in different ways and cancels again for colorless states. A similar discussion for a linear confinement potential is given in ref. 4).

In the case just presented one can define what one means by self and interaction energy up to for the \( s \) wave contribution. It is this "up to" sense that we will use these terms in the following. The intimate connection between self and interaction energies is however not just a matter of principle, a matter of an ambiguous contribution, it is of course
to be expected that both parts are of the same order of magnitude and this is what is found for both light\(^8,9\) and heavy\(^5\) quark systems. For light quark systems both the boundary contributions and the off shell corrections of the free self energy are important. For heavy quarkonia the latter can be neglected as has been done in ref. 5. For static quarks in a spherical cavity the self energy yields a repulsive mirror charge potential\(^10\) behaving as \(\sim 1/4\pi(R-r)\) near the boundary and thereby to confinement. The \(1/c^2\) corrections lead via a Breit-Fermi-type analysis to modified spin dependent terms for the interaction part and to new spin dependent and orbital corrections arising from the self energy. The latter, given in full detail in ref. 5, contain in particular a negative LS term. The essentially scalar nature of the confining self energy part can be understood qualitatively by nothing the place where the \(\gamma^\mu\) appear in the graphs

\[
\begin{array}{c}
\gamma^\mu \\
\gamma^\mu
\end{array}
\]

Though the main aim of my contribution was to recall some features of a general nature that one can study explicitly in a special model, the MIT bag model, I would also like to mention that this analysis of spin dependent forces meets reasonably well with experimental results. Writing the \(\hat{L}_S\) contribution, which is most important in the context, as

\[
E_{\hat{L}_S} = a \cdot \langle \hat{L}_S \rangle
\]

we obtain\(^5\) values of \(a = 30\) MeV for the \(c\overline{c}\) 1P states and of \(a = 14\) MeV and 12 MeV for the \(b\overline{b}\) 1P and 2P states respectively.\(^12\) This is in fair agreement with the corresponding experimental values of\(^13\) 34.9 ± .2 MeV for the \(c\overline{c}\) 1P states and of\(^14\) 12 ± .4 and 10.8 ± 3 MeV for the \(b\overline{b}\) 1P and 2P states.
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12) The absolute position of the P wave multiplets is somewhat off, due to the fact that we have used a simple Coulomb potential at short distances.
14) J. Lee-Franzini, Talk presented at this meeting.
COMMENT ON THE MECHANISM OF DECAY OF HEAVY QUARKONIA INTO LIGHT HADRONS

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Abstract: The mechanism of production of quark-antiquark pairs in the formation of exclusive final states appears to be sequential. This seems to imply nonperturbative and long-distance characteristics for this process.
A crucial aspect of heavy quarkonium decay into light hadrons concerns the mechanism of formation of the final state. It is generally supposed that this proceeds through the formation of a multi-gluon intermediate state, followed by the conversion of these gluons into quark-antiquark pairs. The distance scale where the quark pair creation takes place is important in deciding whether this is a perturbative process (at short distances) or it is nonperturbative, like string breaking for example. In a string breaking model for the formation of the final state the quark antiquark pairs are produced sequentially. In such a model it is possible to explain the observed suppression of the decay of psiprime into the rho-pi final state\(^1\). In a recent preprint\(^2\) it was noted that decays into three meson final states offer the possibility to check whether the sequential mechanism is responsible for the formation of the final state. For example, the final state phi-pi-pi cannot be produced sequentially whereas omega-K-Kbar can. At the conference some evidence was presented which favors the sequential mechanism\(^3\). In the discussion which followed my description of this mechanism, Petronzio and Mueller\(^4\) questioned whether the concept of sequential fragmentation was gauge invariant in QCD. The whole notion of sequential fragmentation is appropriate for a nonperturbative description of the formation of the final state, as in string breaking. Therefore the evidence for sequential fragmentation in quarkonium decay means that the quark-antiquark pairs of the final state are not produced at short distances where a perturbative description would be valid.

2.) G. Karl and San-Fu Tuan, Hawaii preprint UH-511-571-85, January 1986
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ABSTRACT: We present numerical data for 2 and 8 dynamical quark flavours on $2 \times 6^3$, $4 \times 6^3$ and $6^4$ lattices. We discuss the dependence of the order of the transition on the number of flavours and on the number of sites in the temperature direction. An estimate of the critical temperature in physical units is given.

* Presented by B. Petersson.
I. Introduction

In the last few years there has been a considerable interest in the behaviour of nuclear matter at finite temperature and density. On the one hand nuclear matter at high temperature and density will be produced in heavy ion experiments, planned for the near future. On the other hand, QCD should give precise predictions on the thermodynamic behaviour of strongly interacting matter.

Within QCD it is natural to assume that at high temperatures and densities nuclear matter turns into a new phase, a quark-gluon plasma. The critical behaviour is, however, a non-perturbative phenomenon. For this the only method known is lattice QCD. In the SU(3) pure gluon theory quite extensive calculations also on large lattices have been performed in the last years. However, one would expect that the production of quark-antiquark pairs would severely influence the critical behaviour. Thus, although the calculations in full lattice QCD are very cumbersome because of the fermion determinant, studies of the finite temperature behaviour in this theory were started some time ago by several groups.

In this contribution we will report on some continuation of this work, in particular investigating the dependence on the number of quark flavours and of the lattice size.

II. The Formalism

For the derivation of the finite temperature field theory and its lattice regularization we refer to the literature. In our calculations we have used the partition function

$$Z(g^2, N_T, N_s, n_f, m_a) = \int \Pi dU e^{-S_G(U)} [\det (\mathcal{D} + m_a)]^{n_f/4}$$

and expectation values have been calculated with the measure in the above integral. Here $a$ is the lattice spacing, $N_T$ and $N_s$ are the number of sites in the temperature and space directions respectively, and $m$ is the bare quark mass.

For the gluon part we use the Wilson action

$$S_G(U) = \beta \sum_{\mu \nu} \left( 1 - \frac{1}{3} \text{Re} \text{tr} [U_\mu(x) U_\nu(x + \hat{\mu}) U'_\mu(x + \hat{\nu}) U'_\nu(x)] \right)$$

where $U_\mu(x)$ are SU(3) matrices on the links $(x, \hat{\mu})$ of the lattice and $\beta = 6/g^2$.

For the quark part we use the Kogut-Susskind regularization. Thus we take

$$(\mathcal{D} + m_a)_{x,x'} = \sum_{\mu = 1}^{4} \frac{\tau}{2} \Gamma_\mu(x) \left[ U_\mu(x') \delta_{x',x+\hat{\mu}} - U'_\mu(x') \delta_{x,x'+\hat{\mu}} \right] + m_a \delta_{x',x}$$

where

$$\Gamma_\mu(x) = \frac{1}{2} \sum_{\nu \neq \mu} \sum_{x'} \left[ \delta_{x,x'} \delta_{x',x+\hat{\nu}} - \delta_{x,x''} \delta_{x'',x+\hat{\nu}} \right]$$
\[ \Gamma_{\mu}(x) = (-1)^x_1 + x_2 + \ldots + x_{\mu-1}. \]  

The eigenvalues of $D + m$ come in complex conjugate pairs, because of a "chiral" symmetry of $D$, from which follows

\[ \Gamma_{\pm}(x) D(x,x') \Gamma_{\pm}(x') = -D(x,x') \]  

where

\[ \Gamma_{\pm}(x) = (-1)^{x_1 + x_2 + x_3 + x_4}. \]

Furthermore,

\[ D^+ = -D. \]

For the parameter $n_f = 4$, the theory corresponds to four degenerate quark flavours, at least in the classical continuum limit. The interpretation is discussed in detail in ref. 6. Following Hamber et al., \(7\) we define the theory for $n_f \neq 4$ by formula (1). The measure is still positive definite:

\[ Z = \int \Pi U(x) e^{-S_G(U) + \frac{n_f}{2} \ln[-D^2 + m^2a^2]} . \]

The theory has the following symmetries.

(i) For $n_f \to 0$ or $ma \to \infty$ we refine the pure gluon theory. In this limit we can define an order parameter

\[ L(\bar{x}) = \text{tr} \prod_{t=1}^{N_T} U_\bar{x}(\bar{x},t) . \]

For $<L> = 0$ we have confinement (of static quarks). The existence of a phase with $<L> \neq 0$ corresponds to deconfinement. If the phase transition is first order it may still exist in the interior of the $(n_f, ma)$ plane.

(ii) For $ma \to 0$, any $n_f$, the chiral invariance discussed above, can be described by an order parameter

\[ <\bar{x}x> = \lim_{m \to 0} \text{Tr} (D + ma)^{-1} \]

In the strong coupling region at least this symmetry is broken at low temperatures \(8\) and restored at high temperatures through second order phase transition \(9, 10\).

The physical temperature and volume are given by

\[ T = (N_T a)^{-1} \quad \text{and} \quad V = (N_0 a)^3 \]

where the lattice distance can be determined from other measurements of e.g. masses, if it is universal (scaling). In the asymptotic scaling region
In our calculations we have used the pseudofermion method. Other groups have also used different algorithms, like the microcanonical method\(^ {4}\) or the Langevin algorithm\(^ {11}\). It is encouraging that where calculations using different methods can be compared, the results agree quite well\(^ {12}\).

For our \(n_f = 2\) data the numerical procedure is described in detail in refs. 2 and 13. For \(n_f = 8\) we have used a heat bath method for the inversion of \((\mathcal{D} + \mathcal{M})\), needed in the update of the \(U\)'s. We discard the first 4 pseudofermion sweeps and average the following 20. The 6U was chosen so that the acceptance became \((73 \pm 8)\%\). More details will be discussed in ref. 14.

IV. The Results

We have up to now performed calculations for \(N = 6\) and \(N = 2\) and 4 for \(n_f = 2\) and 8. We have also some preliminary data on a 6\(^4\) lattice for \(n_f = 8\). In this report we will concentrate on the measurements of \(\langle L \rangle\). We also compare with the earlier \(n_f = 0\) (quenched) data of Kogut et al.\(^ {15}\).

The results for \(N = 2\) are shown in Fig. 1. At \(n_f = 0\) there is a first order phase transition. Clear two-state signals have been observed. Moving in \(n_f\) for fixed \(m^2 = 0.1\) a sharp but continuous transition was observed for \(n_f = 2\). We now show that this transition becomes very broad for \(n_f = 8\).

A similar behaviour is seen for \(N = 2\) if one moves at fixed \(n_f\) for decreasing \(m^2\). This calculation was made for \(n_f = 2\) only\(^ {16}\). The transition gap goes to zero for some finite mass. It does, however, not disappear as fast as was claimed in calculations using the hopping parameter expansion\(^ {17}\).

For \(N = 4\) we again show results for \(m^2 = 0.1\) and \(n_f = 0, 2\) and 8 in Fig. 2. Note that the scales are different. Still, the behaviour is now quite different. Although the first order phase transition seems to disappear between \(n_f = 0\) and \(n_f = 2\) it comes back at \(n_f = 8\). In fact, at \(n_f = 0\) \(\Delta L \approx 0.4\), at \(n_f = 2\) \(\Delta L < 0.15\) and at \(n_f = 8\) we see a clear two-state signal over 2000
iterations at $\beta = 4.81$ with $\Delta L = 0.5$. The latter was already observed by Kogut et al.\cite{18} and by Gavai\cite{12}. The value of the critical coupling constant is slightly different in the three measurements. Kogut et al. get $\beta_C = 4.67 \pm 0.1$, Gavai $\beta_C = 4.78$ and we get $\beta_C = 4.81$. The spatial volume in the first two works is $8^3$ instead of our $6^3$. One should investigate if the difference is due to finite size corrections and/or systematic errors in the algorithms.

In Fig. 3 we show the "world data" on the finite temperature transition in lattice QCD. The behaviour of the $N_t = 2$ data is quite as expected. The phase transition disappears somewhere in the $(m_a, n_f)$ plane, because the fermion term breaks the corresponding symmetry. The question is only if the transition disappears already for physical values of $n_f$ and $m_a$, and what is the influence of a possible chiral transition on $\langle L \rangle$. Unfortunately, $N_t = 2$ is probably in the strong coupling region. Here a second order phase transition is predicted\cite{9,10} with $\langle L \rangle$ being exponentially small in the phase where chiral symmetry is broken\cite{10}.

The $n_f$-behaviour of the $N_t = 4$ data seems very puzzling. It has been suggested that the $n_f = 8$ first order transition would be a zero temperature transition, related to the breakdown of asymptotic scaling for large $n_f$\cite{18}. To check this suggestion we have started calculations with $n_f = 8$ on a $6^4$ lattice. Preliminary results show $\beta_C \gtrsim 4.9$, well above the critical $\beta_C = 4.81$ on the $4 \times 6^3$ lattice. A $T = 0$ transition should be at the same place for $N_t = 4$ and 6, apart from finite size effects. From the discussion above these are probably smaller than $\Delta \beta_C \lesssim 0.03$. Thus another explanation for the first order transition at $n_f = 8$ should probably be considered.

Finally a calculation of the $\rho$ mass at $n_f = 2$, $\beta = 5.45$ on a $6^4$ lattice was performed by Billoire and Marinari\cite{19}. Their results give $a^{-1} = 800 \pm 200$ MeV at this $\beta$-value. The $N_t = 4$ data would then correspond to $T_C = 200 \pm 50$ MeV. A similar estimate from calculations with $n_f = 3$ has been given by C. Rebbi\cite{20}. However, with such a low ultraviolet cut-off, these numbers should probably not be taken as a quantitative prediction for the continuum.

V. Concluding Remarks

In the pure SU(3) theory the results on rather small lattices was already a good guide for the results on large lattices: the transition is first order and
\[ T_c \approx \sqrt{\sigma/2}, \] where \( \sigma \) is the string tension. These results are confirmed on the large lattices. Therefore one might hope that the same happens in full lattice QCD. This is the justification for working with rather small lattices.

However, as we have seen, the \( N_f = 2 \) and \( N_f = 4 \) data are possibly contradictory as to the order of the phase transition at least for large \( n_f \). The value of \( T_c \) obtained at \( n_f = 2 \) and \( 3 \) might of course still be at least roughly correct.

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Figure 1: The dependence of $\langle L \rangle$ on $\beta$ at $N_t = 2$ for $n_f = 0$ (o), $n_f = 2$ (o) and $n_f = 8$ (•).

Figure 2: The same as Fig. 1, but for $N_t = 4$. 
Figure 3: $\beta_c$ as a function of $n_f$, transitions of first order (\(\triangle\)) of at least second order (\(\times\)), not decided (\(\square\)). The data for $n_f = 3$ are from ref. 3 and 21, for $n_f = 4$ from ref. 12, 18, 22, for $n_f = 0, 2, 8$ as described in the text. The lower points are for $N_T = 2$, the upper ones for $N_T = 4$. The upper point for $n_f = 4$ is for $N_T = 6$. 
EFFECT OF DYNAMICAL QUARKS ON THE SU(2) GLUEBALL MASS

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ABSTRACT

We show in SU(2) Gauge theory with two degenerate flavours and 8$^4$ lattice that under the alternative assumptions of flavour independence of $\Lambda_{\overline{MS}}$ or that $(<\bar{q} q>_{N_f=0} / <\bar{q} q>_{N_f=2})_{m=0}^{1/3}$ behaves like the corresponding spacings ratio, the two to zero flavour glueball mass ratio is $2.03^{+0.34}_{-0.21}$. 
1. Introduction

Several physical properties of QCD, such as the spontaneous breaking of chiral symmetry and the existence of a critical temperature below which quarks and gluons are free particles and chiral symmetry is recovered, have been shown in the late few years by numerical calculations in the lattice.

In this line, many calculations of the hadronic mass spectrum have been succesfully made.

Another important quantity which can be computed numerically in the lattice is the glueball mass spectrum. From an experimental point of view it should be very interesting to get precise predictions about such mass spectrum because of the difficulty to see their experimental signal. Many numerical computations on such mass spectrum have been made for the pure gauge SU(2) and SU(3) theories, giving for the 0+ glueball a mass around 1 GeV. Theoretical prejudices, and the accumulated experience in the quenched aproximation, suggest that the effects of dynamical fermions would not be relevant for these results.

2. The numerical experiment

In this communication we present some preliminary results on a computation of the plaquette-plaquette correlation function in a SU(2) lattice gauge theory with dynamical fermions and we will try and get some physical conclusion from these results. The motivation to do such a heavy calculation was some previous results on a 4×4 lattice which indicated important changes in the plaquette-plaquette correlation function when light quarks are switched on the system for a SU(2) gauge theory.

Our computation has been made on an 8×4 lattice and SU(2) gauge group (we used the icoshedron group) and effects of dynamical fermions are taken into account by means of the pseudo-fermions method.

All technical details of our program can be found in refs. 3, 5. We have computed the plaquette-plaquette correlation function for two values of the quark mass $m_a = 0.2$ and $m_a = 0.3$ (a is the lattice spacing) and $\beta$ fixed to the value 2.2. The value of $\beta$ was chosen at the beginning of the asymptotic scaling region for the quenched theory. The reason for doing so is that dynamical fermions lower the system temperature. The fact that the Poliakov loop is very small (0.022 for $m_a = 0.3$, 0.023 for $m_a = 0.2$) tells

(*) Recently, B. Berg et al. have claimed that the glueball state with lowest mass is the $2^{++}$, not the $0^{++}$ as previously thought.
us that boundary effects are not relevant in this region of the parameters.

Additionally we used antiperiodic boundary conditions in the time direction in order to decrease the boundary effects.

The mean values are taken over a total number of gauge sweeps equal to 2000 and 3000 for \( ma = 0.3 \), \( ma = 0.2 \) respectively, each containing 40 Monte Carlo (MC) pseudo-fermion iterations, with 3 hits per site. The system was previously thermalized at \( \beta = 2.2 \) and \( ma = 0.3 \) with 500 gauge sweeps, with the corresponding pseudo-fermion iterations.

The pure action is the standard Wilson action and we have used Kogut-Susskind fermions\(^6\) with two degenerate flavours.

Let us report the numerical results. We have computed the time slice plaquette-plaquette correlation function\(^3\) at the separations \( t = 0, 1, 2, 3, 4 \) from the correlation of the plaquette operator \( \text{Tr} (U_1 U_2 U_3 U_4) \) (\( U_1, U_2, U_3, U_4 \) are the four link variables around a plaquette), and the mean energy per plaquette. Unfortunately we cannot use the plaquette-plaquette correlation function at time distance \( t=3 \) because statistics is not sufficient. In fact in some cases the correlation at \( t=3 \) is negative even after 1000 MC iterations. In Fig. 1 we plot the plaquette-plaquette correlation function against the number of MC gauge sweeps for \( \beta=2.2 \) and \( ma=0.2 \). It can be seen that the correlations for time separations \( t=0, 1, 2 \) averaged over groups of 700 and 1000 iterations fluctuate around the mean value.

3. The glueball mass from \( \Lambda_{\text{MS}} \)

Using the numerical results for the correlation function at \( t=0 \) and \( t=1, t=2 \) (\( ma=0.2 \)) and fitting this correlation function to an exponential (we neglect the possible mixing between the plaquette operator and \( \bar{q}q \) states\(^9\)) we get for \( m_g a \) the values

\[
(\text{am}_g)_{0,1} = 2.12 \pm 0.01, \quad (\text{am}_g)_{0,2} = 1.79 \pm 0.21,
\]

where \( m_g \) stands for the glueball mass. The estimation of the statistical errors in these values of the mass in unphysical units has been made computing the same quantity over groups of 1000 MC iterations and calculating the quadratic deviation of the mass values obtained in this way from the value obtained from the total sample. To give a numerical prediction for the glueball mass in physical units we need to estimate the value of the lattice spacing \( a \) at \( \beta=2.2 \) for the theory with two flavours. The best we can do in this sense is to assume we are in the asymptotic scaling region\(^8\), and to estimate the value of the lattice spacing from the renormalization group equation

\[
a = \Lambda_{\text{lat}}^{-1} \cdot e^{-1/2 \theta(\beta_0^2)} \cdot (\beta_0^2)^{-\beta_1/2 \theta_0} \quad (3.1)
\]
\[ \beta_0 = (4\pi)^{-2} \cdot \left( \frac{N_1}{3} N - \frac{2}{3} N_F \right) \]  
(3.1)

\[ \beta_1 = (4\pi)^{-4} \cdot \left( \frac{34}{3} N^2 - \frac{10}{3} N \cdot N_F - (N^2 - 1) \right) \left( \frac{N_F}{N} \right) \]

\( N \) = number of colours

\( N_F \) = number of flavours

From eq. (3.1) one can see that we need to know also the value of \( \Lambda_{\overline{\text{MS}}} \) for the two flavour theory in order to estimate the value of \( a \). Assuming, from the results on \( \Lambda_{\overline{\text{MS}}} \) obtained by means of the string tension calculations (zero flavour theory) and from the experimental measurements of \( \Lambda_{\overline{\text{MS}}} \) for the four flavour theory, that \( \Lambda_{\overline{\text{MS}}} \) is practically independent of \( N_F \) and using the corresponding conversion formulae to \( \Lambda_{\overline{\text{MS}}}^{(10)} \) we get

\[ \frac{\Lambda_{N_F=2}^{\text{latt}}}{\Lambda_{N_F=0}^{\text{latt}}} = 0.552 \]  
(3.2)

which implies

\[ \left[ \begin{array}{c} a|_{N_F=0}^{\text{latt}} \\ a|_{N_F=2}^{\text{latt}} \end{array} \right] \beta=2.2 = 3.6, \quad \left[ \begin{array}{c} a|_{N_F=0}^{\text{latt}} \\ a|_{N_F=2}^{\text{latt}} \end{array} \right] \beta=2.2 = 2.0 \]  
(3.3)

On the other hand, the values of \( m_{ga} \) for \( \beta=2.2 \) in the pure gauge theory reported in ref. 11, where the same lattice size was used, combined with equation (3.3) and our numerical results for \( m_{ga} \), give the predictions for the mass ratio \( r = m_{N_F=2}/m_{N_F=0} \)

\[ r(0,1)=2.36^{+0.05}_{-0.06} \quad r(0,2)=2.03^{+0.34}_{-0.21}, \]  
(3.4)

where the arguments stand for the time separations used to evaluate the masses(*).

4. The glueball mass from \( \langle \bar{\psi} \psi \rangle \)

However, we feel that the assumption of independence of \( \Lambda_{\overline{\text{MS}}} \) on the number of flavours is mainly based on the great uncertainty in the experimental determination of \( \Lambda_{\overline{\text{MS}}} \) in the four flavour theory\(^{12}\), and, therefore, too strong to be taken as the basis of our prediction. Because of that, we prefer to make another argument independent of this assumption which will tell

(* We are in the process of calculating the pure gauge glueball mass from the plaquette-plaquette correlation function\(^8\)).
us how well this assumption works. This argument consists on noticing (see ref. 13) that \( \langle \bar{\psi} \psi \rangle \) should scale as a times an anomalous dimension factor independent of \( N_F \) up to the order it has been computed. Then, the following expression must hold

\[
\left( \frac{\langle \bar{\psi} \psi \rangle_{N_F=0}}{\langle \bar{\psi} \psi \rangle_{N_F=2}} \right)^{1/3} = \frac{a_{N_F=0}}{a_{N_F=2}}
\]

(4.1)

Taking the value of \( \langle \bar{\psi} \psi \rangle \) in the quenched theory at \( \beta=2.2 \) from ref. 14 and doing a linear extrapolation of our results for \( \langle \bar{\psi} \psi \rangle \) at \( \beta=2.2 \) and \( m=0.1 \) and \( m=0.2 \) to \( m=0 \) we get

\[
\frac{a_{N_F=0}}{a_{N_F=2}} = 1.92 \pm 0.16 - 0.12
\]

(4.2)

in fair agreement with (3.3).

5. Conclusions

In conclusion, we have run a computer simulation of a SU(2) lattice gauge theory fermions. We have measured the plaquette plaquette correlation function and Poliakov loop for \( \beta=2.2 \) and \( m=0.3, m=0.2, \) and \( \langle \bar{\psi} \psi \rangle \) for \( m=0.2, m=0.1 \). From the numerical results plus the alternative assumption that we are in the asymptotic scaling region and \( \Lambda_{\overline{MS}} \) is independent of the number of flavours or that \( \langle \bar{\psi} \psi \rangle_{N_F=0}/\langle \bar{\psi} \psi \rangle_{N_F=2} \) behaves like the ratio of the corresponding spacings, we get for the mass of the glueball a factor 2 with respect to the pure gauge theory.

In view of these results, we think that further investigations on these points are needed.

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Fig. 1.- Expectation values of the plaquette-plaquette correlation function at $t=0$, $1$, $2$ against the number of MC gauge iterations for the quark mass $m_a=0.2$ and $g=2.2$. The solid lines are the mean values over 3000 MC iterations. The $\Delta$ are the averages over consecutive groups of 700 MC iterations and the $\cdot$ are the same but over groups of 1000 MC iterations.

The $X$ are successive averages over groups of 350 MC iterations.

The scale for the solid line is on the left-hand side axis and it is logarithmic. The scales for all the other symbols (on the right-hand side) are linear and centred around the mean values.
Abstract
The possibility of a bound-state effect in the quark-quark interaction is discussed. If such dynamical diquarks exist, they should play a role in the formation and decay of a QCD plasma, as probed in relativistic heavy-ion collisions. In an intermediate temperature range, just above the critical temperature, a diquark plasma component is expected to be important. Quantitative results within a particular model, the Stockholm diquark model, are presented. Possible experimental signatures of such a diquark plasma are suggested.
1. Introduction

A Stockholm group (S. Fredriksson and myself, previously also M. Jändel and T.I. Larsson) has for the past few years investigated the hypothesis that there exist a bound state in the quark-quark interaction, a diquark. So far the model has been confronted with data from deep inelastic lepton-nucleon scattering\(^1,2\), \(e^+e^-\) annihilation\(^3\), and hadron-hadron scattering\(^4,5\). It is in line with world data, and for some phenomena diquarks seem to provide the best explanation on the market.

The Stockholm diquark model prescribes that two quarks of unequal flavour and colour can form a scalar diquark. Excited states do not seem to be needed to account for the data, so for simplicity and economy we assume that only the ground state is bound, i.e. a diquark has colour \(3^*\) and \(\mathcal{J}^P = 0^+\).

Since diquarks are not colour singlets, they are subject to confinement, but they could occur

- inside hadrons,
- in the fragmentation process, or
- in a QCD plasma,

on an equal footing with single quarks.

Diquarks are expected to occur inside spin-\(\frac{1}{2}\) baryons (as well as in multiquark states, such as dibaryons). We assume that \(p = u(ud)\) and \(n = d(ud)\).

Since the diquark is an extended object, the interaction amplitude is suppressed by a form factor. This \(Q^2\) dependence can explain the scaling violation in deep inelastic structure functions\(^1,2\), as well as the \(p_T\) and \(\phi\) dependences of high-\(p_T\) proton production in hadron-hadron collisions\(^4,5\).

Here I will focus on the third possible arena for diquarks in physics: the QCD plasma. If diquarks are bound, it is energetically favourable for two quarks in the plasma to form a diquark. However, for very high temperatures, i.e. \(T \gg \) diquark binding energy, we expect that most diquarks will dissociate into two quarks. It is in an intermediate temperature range, just above the critical temperature, that diquarks could conceivably be an important component in the plasma. Thus, it is possible that the experimentally observable decay of the plasma, which reflects the plasma properties at "low"
temperatures, is influenced by the presence of diquarks.

In the following, I present a model of the plasma, in order to arrive at a quantitative estimate of the fraction of quarks in the plasma that are bound in diquarks.

2. Model of the plasma

Regard a QCD plasma with $N_i$ particles of type $i$ in a volume $V$ as an ideal relativistic gas of fermions and bosons. Consider the process

$$u + d \rightarrow (ud)$$

as a chemical reaction, and use thermodynamics to calculate the relative abundances, assuming thermal and chemical equilibrium. In this approach, the gluonic degrees of freedom are regarded as an effective "heat bath", which is responsible for the thermalisation.

We have contributions to the grand canonical partition function $\Xi$ from particles of type $i$:

$$T \ln \Xi_i = \int \sigma_i(E) \left( \exp \frac{E - \mu_i}{T} \pm 1 \right)^{-1} dE ,$$

where $T$ is the temperature and $\mu_i$ is the chemical potential. The plus sign is applicable for fermions and the minus sign for bosons.

The integrated one-particle state density function is

$$\frac{E}{m_i} \int_{\varepsilon_i}^{\infty} \rho_i(E) dE = \sigma_i(E) = n_i \int_{\varepsilon_i}^{\infty} \frac{V d^3P}{(2\pi)^{3/2}} = \frac{n_i}{6\pi^2} V \frac{E^2 - m_i^2}{2} \frac{3}{2}$$

for $E > m_i$. $\eta_i$ is the degeneracy factor, $\varepsilon_i$ the energy, and $m_i$ the mass of one particle of type $i$.

For the particle number densities $n_i = N_i/V$ and the pressure $P$ of the plasma one gets

$$n_i = \frac{1}{V} \frac{\partial}{\partial \mu_i} (T \ln \Xi_i) = \frac{1}{V} \int_{m_i}^{\infty} \rho_i(E) \left( \exp \frac{E - \mu_i}{T} \pm 1 \right)^{-1} dE ,$$

and

$$P = \frac{1}{V} T \ln \Xi = \frac{1}{V} \sum_{i} \int_{m_i}^{\infty} \sigma_i(E) \left( \exp \frac{E - \mu_i}{T} \pm 1 \right)^{-1} dE .$$
The condition for chemical equilibrium is

\[ \mu_u + \mu_d = \mu_D , \] (6)

where we have introduced the notation \( D \) for the \((u \ d)\) diquark, and the condition of total quark-number conservation

\[ 2n_D + n_u + n_d = n , \] (7)

where the quark-number density \( n \) is the total net number of quarks (including those bound in diquarks) per unit volume.

For the effective quark and diquark masses we use \( m_q = 400 \text{ MeV} \) and \( m_D = 500 \text{ MeV} \), corresponding to a diquark binding energy of 300 MeV. This choice can be motivated by comparing the nucleon \((qD)\) mass \( \approx 900 \text{ MeV} \) to the delta \((qqq)\) mass \( \approx 1200 \text{ MeV} \). Also, from classical considerations one can show \(^6\) that when two objects of mass 400 MeV bind to form a state of mass 500 MeV, the radius of the system becomes 0.25 fm, which is the value favoured by our earlier analyses of \( \mu p \) and \( pp \) scattering\(^1,5\).

We take into account only the particle types \( i = D \) and \( i = q \) (= \( u \) or \( d \)), i.e. we neglect strange and heavier quarks, and also pair production of even the lightest quarks. This can be motivated, since production of real quark-antiquark pairs should be negligible when \( T \ll 2m_q \). We are interested in the low-temperature structure of the plasma, and we assume the effective lightest quark mass to be 400 MeV.

For simplicity we consider the isoscalar case

\[ n_u = n_d . \] (8)

In this case one also gets

\[ \mu_u = \mu_d = \mu . \] (9)

This is applicable to the forward baryon-rich fragmentation region in the CERN SPS heavy-ion programme, with \( ^{16}O, ^{32}S, \) and, possibly, \( ^{40}Ca \) beams\(^7\).

In order to solve the system (4) - (9) for \( \mu \), and get the diquark fraction, the pressure, and the volume of the plasma solely as
functions of temperature (and of the initial conditions at plasma formation), we need an additional relation; an equation of state for the cooling plasma. Here I will present results for two cases, namely constant volume or constant pressure. Probably neither of these scenarios is very realistic, and the results should be considered as preliminary. It has been argued, for instance, that \( p \propto T^4 + \text{const} \) is a good candidate as the equation of state.

3. Results and conclusions

The resulting diquark fractions

\[
x_D = \frac{\text{n}_D}{(\text{n}_D + 2\text{n}_q)}
\]

as functions of temperature are shown in Figures 1 and 2 for some special cases of constant volume and constant pressure, respectively.

![Figure 1](image1.png)

**Fig. 1** Diquark fractions in the plasma as functions of the temperature for constant volume. Curve A corresponds to the density \( n = 4n_0 \) and curve B to \( n = 7n_0 \). \( n_0 \) is the density of ordinary nuclear matter.

![Figure 2](image2.png)

**Fig. 2** Diquark fractions as functions of temperature for constant pressure. Curves A, B, C, and D correspond to the pressures

- A : \( P = P(T=.15,n=4n_0) \)
- B : \( P = P(T=.20,n=7n_0) \)
- C : \( P = P(T=.30,n=4n_0) \)
- D : \( P = P(T=.30,n=7n_0) \)

(\( T \) in GeV)

The line \( x_D = 0.5 \) has been indicated, since it is of special interest whether \( x_D \) is greater than 0.5 or not at the hadronisation temperature. \( x_D = 0.5 \) corresponds to equal number of diquarks and single quarks, which is the situation we have in spin-\( \frac{1}{2} \) baryons.
If $x_D > 0.5$ there are "too many" diquarks for the plasma to be able to decay completely into colour-singlet hadrons by baryon and meson formation alone. The excess diquarks do not readily split up into quarks, as the system is too cool. Instead, we expect the formation of multiquark states, such as three-diquark systems or dibaryons, as an interesting experimental signature of this type of baryon-number rich QCD plasma decay.

Since the hadronisation temperature is widely believed to be of the order of the pion mass, it seems that all cases depicted in the Figures would lead to this excess-diquark effect. However, one must be careful in noting that those curves were calculated using constant volume or constant pressure, whereas a more realistic equation of state would presumably imply that the pressure and density decrease as the plasma cools. This would tend to disfavour the diquark component. More detailed work is in progress.

To conclude:

- If diquarks are bound, they should occur as a component in the QCD plasma.
- If the binding energy $\geq T$, they should be important.
- If the diquark fraction $> 0.5$ at the hadronisation temperature, we expect multiquark states such as dibaryons as an experimental signature.

Inspiring discussions with Sverker Fredriksson, Magnus Jändel, and Olle Edholm are gratefully acknowledged, as is a travel grant from the Wallenberg foundation.

References

INITIAL TEMPERATURE IN ULTRARELATIVISTIC NUCLEAR COLLISIONS

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Abstract

The thermalization problem in high-energy nuclear collision is studied in the parton model with nuclear transparency and distributed contraction as the essential physical input. The initial temperature and thermalization time are determined in the approximation of short collision time.

One of the program organizers of this meeting jokingly gave this talk a subtitle: "How to bake a quark-gluon pie." It is not a bad joke because, whether to bake a pie or to make a plasma, it is important to know what temperature the oven should be set at. It is not possible at this point to determine from first principles the initial temperature \( T_i \) in a high-energy nuclear collision, but an educated estimate is not impossible, and it is especially important to make that estimate when the funding for the construction of a heavy-ion collider is being considered. I present here a study on the subject based on the parton approach to the thermalization problems.\(^1\)

A question that can be asked is whether \( T_i \) would increase indefinitely if we let \( s \to \infty \). Our answer is no, modulo nonscaling effects such as the growth of \( \sigma_T \) and \( dN/dy \) with \( s \), on which there are no firm theoretical predictions.
Basically, the problem is to convert the organized, kinetic energy of the incident partons into random, thermal energy through interaction and compression. The difficulty in achieving arbitrarily high $T_i$ is due to the basic properties of partons in high-energy collisions: (a) nuclear transparency and (b) distributed contraction. The former refers to negligible nuclear stopping power on fast quarks, which underlies the phenomenological fact that in $a + b \rightarrow c + X$ the inclusive distribution of $c$ in the fragmentation region of $a$ at high energy is independent of $b$, if $b$ is a hadron, or only mildly dependent on $b$ if $b$ is a nucleus, $A$. Distributed contraction refers to the quantum mechanical property that the uncertainty in longitudinal position $\Delta z$ of the partons is contracted in accordance to their rapidities; thus $\Delta z \approx \mathcal{L} / \cosh y$ where $\mathcal{L}$ is the hadronic length scale of about 1 fm. This implies that only the fast partons are contracted to a thin disk; the wee partons are not. Combining these two properties we see that the thermal energy density cannot increase indefinitely, even as $s \rightarrow \infty$, because by (a) the fast partons from the two nuclei, though contracted, do not interact effectively, while by (b) the slow partons that do interact strongly are not densely packed.

An important part of the procedure is to recognize the global change of the rapidity distribution in a space-time cell as a function of the proper time $\tau$. Let the location of the cell be specified by $(\tau, \eta)$, where $\eta$ is the spatial rapidity, $\eta = 1/2 \ln (t + z)/(t - z)$. Fixing $\eta$ at any value and increasing $\tau$, we know in the no-interaction case that the parton rapidity distribution $P(\tau, \eta, y)$ for $\tau \approx 0$ is very broad, from $-Y$ to $+Y$, and decreases in width until it approaches $\delta(y - \eta)$ as $t \rightarrow \infty$, at which point the partons have hadronized and end up in a detector located at $\eta$. Property (b) mentioned above leads to the independence of the width $2\Delta$ of $P(\tau, \eta, y)$ on $\eta$ for fixed $\tau$. The dependence of $\Delta$ on $\tau$ is calculable.

In the realistic situation parton interactions change $P(\tau, \eta, y)$ to the true distribution $F(\tau, \eta, y)$. The determination of how $P(\tau, \eta, y)$ becomes $F(\tau, \eta, y)$ is hard, and is the heart of the thermalization problem. Only partial attacks on the problem have so far been attempted. Fortunately, for our purpose of calculating $T_i$ and the thermalization time $\tau_i$, the full solution is not needed, provided that high accuracy is not demanded. In the following discussion we need only regard $F(\tau, \eta, y)$ as the thermal distribution, characterized by the temperature $T(\tau)$ at every cell.

Now, in the free-streaming case we have $P(\tau, \eta, y)$ from which the energy density $\varepsilon'(\tau)$ can be calculated from kinematics and the parton density of the colliding nuclei, assuming that each nucleon has partons in the sea (i.e. central rapidity region) as prescribed in the parton model. In the thermalized system we have $F(\tau, \eta, y)$ from which the energy density $\varepsilon(\tau)$ can also be
calculated. The key point in our approach is to identify \( e'(\tau) \) with \( e(\tau) \) at \( \tau = \tau_i \) as a first order requirement in the approximation of short collision time. Our reasoning is that the change \( P \to F \) is due to local interaction in rapidity. There are exchanges of partons among neighboring cells and creation of partons within the cells. So long as partons in the fragmentation region do not contribute to \( e(\tau) \) in the central region (due to nuclear transparency), energy conservation and boost invariance along a fixed \( \tau \) hyperbola within the central region require that

\[
e'(\tau_i) = e(\tau_i) \tag{1}
\]

Independent of the normalizations of the distribution functions we can also separately obtain

\[
e'/\rho' = f(\tau), \quad e/\rho = 3T - \tau^{-1/3} \tag{2}
\]

where \( \rho' \) and \( \rho \) are the parton densities in the free-streaming and thermal cases. \( f(\tau) \) is completely calculable and behaves as \( m'_T l/4\tau \) at small \( \tau \), where \( m'_T \) is the transverse mass of the partons. We shall set \( m'_T l \approx 3 \). At large \( \tau \), \( f(\tau) \) approaches \( m'_T \). In Fig. 1 we show both parts of (2). The first intercept at the smaller \( \tau \) value corresponds to the solution we seek because in that case we have

\[
\tau_i = 1/4 
\]

which exhibits the correct inverse relationship between \( \tau_i \) and \( \tau_i \). The second intercept at higher \( \tau \) has no physical significance. The realistic solution for the physical system presumably follows a smooth curve shown in Fig.1; its distance from the asymptotes depends on the detailed properties of soft interaction. If the collision time \( \tau_c \) is not too large, the intercept of the asymptotic lines of (2) provides a good approximation for the true values of \( \tau_i \) and \( \tau_i \). As this point there is no good estimate of \( \tau_c \). The conservative view that it is on the order of 1fm/c is not persuasive, since the pion mass or the perturbative-QCD scale parameter \( \Lambda \) is not necessarily relevant for a highly turbulent and high-density system whose natural length scale may well be far less than the usual hadronic scale. The fact that \( \tau_i \) is around 0.15 fm/c (see below) attests to the possibility that \( \tau_c \) can be equally small at early times when the energy and parton densities are high.

Since \( e'(\tau) \sim \tau^{-2} \) and \( e \sim \tau^4 \), we obtain from (1) and (3) a value for \( \tau_i^2 \) that depends only on the properties of the colliding nuclei

\[
\tau_i^2 \sim \int dz \rho'(z) \sim A^{1/3} \tag{4}
\]

where the integration is over the average longitudinal nuclear distance, \( L = 4R_A/3 \). Thus \( \tau_i^2 \) is a measure of the integrated parton density per unit
transverse area. In principle, in an $A + A'$ collision ($A$ right-moving, $A'$ left-moving) $T_i$ for $y > 0$ depends on $A^{1/6}$ while for $y < 0$ the dependence is $A'^{1/6}$, but in practice it may be hard to verify the difference experimentally, since, apart from other complications, thermal conductivity in the plasma tends to equalize the temperature in the system. Putting in typical numbers for nuclear and quark densities, we get for $AA$ collision

$$T_i = 180 A^{1/6} \text{ MeV}$$

$$\tau_i = 0.27 A^{-1/6} \text{ fm/c}$$

The $A$ dependence is very weak. For $16 \leq A \leq 238$, we have $300 \leq T_i \leq 500$ MeV and $0.1 \leq \tau_i \leq 0.2$ fm/c. The value for $\tau_i$, being smaller than $\mathcal{L}$, does not imply inconsistency with the uncertainty principle because the high temperature in a cell arises from partons with $|y - \eta| \sim \sinh^{-1} \mathcal{L}/2\tau_i$, so the corresponding uncertainty in longitudinal position $\Delta z$ is roughly $\tau_i$, not $\mathcal{L}$. The average value of $T_i$ estimated is well above the phase transition temperature $T_c (= 200$ MeV) determined by nonperturbative QCD calculations on the lattice. The results are for $s$ large enough so that the central region with no net baryon number can be developed. The projected RHIC energy range would be adequate; any higher energy would not lead to any significant gain in the formation of a quark-gluon plasma.

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Fig. 1. Sketch of log-log plot of $\frac{\varepsilon'}{\rho'}$ and $\frac{\varepsilon}{\rho}$ versus $\tau$. The broken curve represents a possible real solution to the problem.
Abstract

We analyse the leading baryon spectrum in nuclear collisions in the framework of the Dual Parton Model. The agreement with available data on hadron–nucleus and nucleus–nucleus collisions is very good. The nuclear stopping power is studied as a function of the event selection criteria.

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The question of how much energy is lost by a proton when it collides with a nuclear target is of crucial relevance in order to know whether or not a quark-gluon plasma can be created in nuclear collisions. This energy loss determines the spectrum of the leading baryons in the final state of the collision and allows us to estimate the nuclear stopping power.

In this paper we shall analyse this problem in the framework of the Dual Parton Model (DPM)\(^1\). So, let us recall briefly the fundamentals of the DPM. In the DPM particle production takes place in chains formed between constituents of the colliding hadrons. In a single nucleon-nucleon collision only valence constituents are involved: each nucleon splits into a slow quark and a fast diquark and two quark-diquark chains are formed. However if a given nucleon undergoes a multiple collision, its sea constituents must also be excited. In general if a nucleon suffers \(m\) inelastic nucleon-nucleon collisions, \(m-1\) quark-antiquark pairs from the sea have to be excited. In the DPM the energy taken by valence and sea constituents is fixed, without free parameter, by the structure functions (which are obtained from very general Regge arguments). The leading baryon spectrum is obtained by performing the convolution between structure and fragmentation functions. The latter are not specified in the model and must be taken from a fit to the proton-proton data. Once the fragmentation function is fixed, the spectrum in nuclear collisions is determined without additional parameters.

In nuclear collisions it is necessary to perform the average over all multiple-scattering configurations. This can be done by using the Glauber-Gribaux\(^2\) model. In the case of nucleus-nucleus collisions, a detailed and unambiguous calculation of these nuclear averages has been carried out in ref. 3.

Before continuing an important remark is in order. As it is evident from the physical picture described above, the incoming nucleons lose their identity as a result of the collision with the nuclear target. Therefore the leading baryons are produced outside the nucleus and thus they do not propagate through nuclear matter. This is in fact the standard picture of space-time development of the hadronic interaction, which is at the origin of the nuclear transparency.

Let us turn now to the results. In ref. 4 the fragmentation function for the leading baryon has been obtained from a fit to the pp and pp spectrum. It has been shown in this same reference that the DPM prediction for the leading baryon spectrum in proton-nucleus and antiproton-nucleus are in perfect agreement with the experimental data. The baryon spectrum as a function of the number of collisions is presented in figure 1. We observe in this figure the softening of the baryon spectrum with the increasing number of collisions. Let us point out that in the DPM there is a fundamental difference between the first and subsequent collisions. Such a difference was found in ref. 5 from a phenomenological analysis of the experimental data and, as it was mentioned there, it is a straightforward consequence of the DPM. Indeed, in a single inelastic collision energy is removed from the diquark by a valence quark, whereas in all extra collisions the energy is taken by sea constituents whose momentum distribution function is very different from that of valence quarks.
In figure 2 we compare the prediction of the model with the experimental data of ref. 6 on
the leading proton spectrum in α-α collisions at the ISR(√s = 31 GeV/Nucleon). We see a very
good agreement except at the edges of the rapidity space where spectator protons and
intra-nuclear cascading (which have not been taken into account in our calculation) give
important contributions.

Let us turn next to heavy ion collisions. In figure 3 we have plotted the leading baryon
rapidity densities for total and central 16O + 207 Pb collisions at laboratory momentum of 200
GeV/c/Nucleon. In this case by central collision we mean a collision in which all the nucleons
of the oxygen participate. We see that in going from total to central collisions a significant
increase of the baryon density in the central rapidity region (and of the mean baryon rapidity
shift) is obtained. By comparing this baryon density with the corresponding one for all charged
particles (which has been computed in ref. 7) we see that the baryon contamination is about
10% at y* = 0 and becomes 100% at y* = -1.5. In order to study more in detail this effect, let
us consider events where all the baryonic fragments of the lightest nucleus have a longitudinal
momentum fraction x lower than certain fixed value x_c (x_c < 1). Obviously by decreasing x_c we
are selecting events in which baryons are more and more "stopped" (i.e. events more and more
inelastic). However, by doing so, the cross section corresponding to these events gets smaller
and smaller. Therefore the interesting point is to study how much the baryons can be stopped
without having a negligible cross section for the corresponding events. For definiteness let us
consider the cases of p-Pb and O-Pb collisions. Let us define the ratio R as follows: in the case
of p-Pb(O-Pb) R is the cross section of the selected events divided by the inelastic cross
section (the cross section of central collisions, respectively). In both cases R measures the
relative rate of the selected events. In figure 4 we have plotted the average rapidity shift versus
the ratio R for these two cases. We see from this figure that there is a great difference between
p-Pb and O-Pb collisions. In the first case by decreasing R one obtains a regular increase of the
rapidity shift. On the contrary for O-Pb collisions in order to obtain a significative increase of
rapidity shift one has to consider events with very low values of R. This means that there are
substantial differences in nuclear stopping power between hadron-nucleus and nucleus-nucleus
collisions when highly inelastic events are considered.

References
   For a review see A. Capella in "Partons in soft hadronic reactions", edited by R. Van de Walle
2) R. Glauber, in "Lectures in Theoretical Physics", edited by W.E. Brittin and L.O. Dunham


Figure 1

$P_{lab} = 200$ GeV/c

$\frac{dN^{q\bar{q}'}}{dy}$

$y - y_{max}$

Figure 2

$\alpha \cdot \alpha \sqrt{s} = 31$ GeV/nucleon

DPM

$\frac{dN}{dy}$

$y$
Figure 3

Figure 4
Abstract

Factorial rapidity moments of high energy collisions are shown to select short-range fluctuations of non-statistical origin. The observation of a power-law behaviour of these moments as a function of the rapidity resolution would give a clear signal of an expanding quark-gluon plasma.
The present work was motivated by the observation of apparently large short-range fluctuations in the two cosmic ray events with the largest multiplicity\[^{[1]}\]. However, the techniques proposed by A. Biatas and myself in a recent publication\[^{[2]}\] allow one to analyze short range rapidity fluctuations observed in various high energy collisions with high multiplicity. For instance, they may be applied to the very high energy collisions at SPS and Fermilab colliders, and in the future experiments on heavy-ion collisions at Cern. Here, I will only sketch the main results of our paper, reference \[^{[2]}\], with a special emphasis on the physical sources and implications of short-range fluctuations in rapidity.

One unavoidable source of short-range fluctuations is of pure statistical nature. In an event-by-event analysis, the limited multiplicity by itself leads to statistical fluctuations in rapidity. Usually, one gets rid of these fluctuations by summation over many events, but in a way which also eliminates non-statistical fluctuations. Our first goal was to find a method able to keep information on these fluctuations.

As we shall see now, non-statistical fluctuations, if they exist, are related to the space-time development of the collision. For hard collisions, i.e., an interaction time much shorter than $1$ in fermi units, short-range fluctuations in rapidity are to be associated to jets. If a long-lived system is formed, for instance a quark-gluon plasma phase, rapidity fluctuations can develop during the expansion, as we will explain further on. Rapidity is a very suitable variable for this kind of analysis.

Since a long time\[^{[3]}\], it has been recognized that the quantum spreading in rapidity of longitudinal wave-packets can be neglected. Considering the uncertainty relation between longitudinal position ($Z$) and momentum ($P_L$), it is possible to write, for a given rapidity fluctuation $\delta Y$:

\[
\delta Y \cdot \delta Z \cong \frac{\delta P_L \cdot \delta Z}{E_{C.M.}} \cong \frac{\hbar}{E_{C.M.}} \rightarrow 0 ,
\]

where the center-of-mass energy $E_{C.M.}$ is considered to be large.

As a consequence, a given physical perturbation $\delta Z$ will leave a sizeable track in the rapidity distribution, provided one can neglect further rescattering. More precisely, $Y$ and $Z$ behaving as classical conjugate variables, one gets $\delta y \approx \frac{\delta Z}{t}$, where $t$ is the time at which the perturbation took place. Consequently, the determination of the characteristic scales of short-range rapidity fluctuations is the second requirement on our method.

As we have shown in ref.\[^{[2]}\], the measure of factorial rapidity moments $<F_1(M)>$, as a function of the number of rapidity bins, meets our two requirements, namely it eliminates the statistical noise, and is sensitive to the scales of fluctuations. Let us divide the total rapidity interval $\Delta Y$ under study into $M$ equal intervals $I_j$ ($j = 1, M$) and let us
denote \( K_j \), the number of particles (or charged particles) observed in the interval \( I_j \). One can easily compute the event-by-event factorial moment \( F_i(M) \):

\[
F_i(M) = \frac{1}{M} \sum_{j=1}^{M} \frac{K_j(K_j-1)...(K_j-i+1)}{N(N-1)...(N-i+1)},
\]

where \( N \) is the multiplicity in \( \Delta Y \) the total interval. The rapidity resolution \( \Delta Y \), can be varied at will in the limits imposed by the experimental resolution by summing over adjacent bins.

As shown in ref. [2], averaging the factorial moments over the event sample eliminates the statistical noise, but keeps the information on eventual non-statistical fluctuations. Introducing the probability \( p_j \) to find a particle in the interval \( I_j \), and averaging over the sample of high multiplicity events under consideration, one obtains:

\[
\langle F_i(M) \rangle = \int P(p_1, p_2, ..., p_M) dp_1 dp_2 ... dp_M \times \frac{1}{M} \sum_{j=1}^{M} (Mp_j)^i,
\]

where \( P(p_1, p_2, ..., p_M) \) describes the dynamical distribution of probabilities \( p_j \). Note that the average over factorial moments of the observed rapidity distribution allows one to reach the normal moments \( \langle \frac{1}{M} \sum_{j=1}^{M} (Mp_j)^i \rangle \) of the dynamical distribution. This distinction is related to the fact that, at finite multiplicity \( N \), one has to take into account the difference between frequencies \( \frac{K_j}{N} \) and their asymptotic limits \( p_j \) at large \( N \). One has also to keep in mind that \( F_i(M) \) are not the factorial moments of the multiplicity distribution, which do not correspond to event-by-event fluctuations.

Let us consider now the properties of the normal moments, and of their average, in presence of fluctuations which are now restricted to non-statistical ones. In ref. [2] we have checked on numerical examples that the moments reflect the existence of scales in rapidity fluctuations. Varying the rapidity resolution, one notes important modifications in the behaviour of \( \langle F_i(M) \rangle \) when the resolution \( \frac{\Delta Y}{M} \) is of the order of the scale \( \delta Y \) of the input distribution fluctuations. Constant when the resolution is much smaller than \( \delta Y \), the moments show large instabilities when one varies the resolution over the critical value \( \delta Y \).

As a physical application of this method, let us discuss the
possibility of producing short-range fluctuations in the formation of a quark-gluon plasma in high energy collisions. It is known that the quark-gluon plasma will expand during a (proper) time of order of 10-15 fermi units\textsuperscript{[4]}, this time being determined by the ratio of the quark-gluon and pion gas degrees of freedom. The basic physical assumption is that the interaction size $\delta Z$ remains of the order of 1 fermi during the plasma expansion. Comparing with the expansion time, one expects that the plasma phase cannot stretch during the expansion without separating a few times in parts. Moreover, there is some probability that, at each step of separation some pieces of the plasma will hadronize.

In a very simplified description, one can figure out a bubble of plasma of size $\Delta_0 Z$ formed at the initial time $\tau_0$; after a time $\tau_1 \approx 2\tau_0$, this bubble has a size $\Delta_1 Z \approx 2\Delta_0 Z$, and separates into 2 parts, of maximum interaction size $\Delta_2 Z$, with some probability of hadronization of one at least of these pieces. After that, each piece evolves independently from the other, and at a time $\tau_0 \approx 2\tau_1 \approx 4\tau_0$ will separate again. This process may go on till complete hadronization a few steps of the cascade later. Though very simplistic, this model introduces one to the more general type of fluctuations which can be characterized as intermittency. Intermittency is qualitatively defined as a set of fluctuations of different sizes superimposed one on each other. It is possible to show\textsuperscript{[2]} that, independently of the number and nature of the steps, the intermittency pattern of fluctuations can be characterized by a specific behaviour of the moments, namely a power-law behaviour:

$$\langle F_1(\mathbf{M}) \rangle = M^{\varphi_1},$$

where the exponents $\varphi_1$ are numbers characteristic of the strength of the intermittency pattern. In this case, intermittency is described by the probability $p_j$, (see equation 3) written as a product of independent random functions representing the cascading steps. One gets the following formula:

$$p_j = \frac{1}{\nu} \prod_{\nu} W,$$

where $\nu$ is the number of cascading steps, the moments $\{W^1\}$ of the random distributions being related to $\varphi_1$. Note again that formula (4) is independent on $\nu$ and on the detailed structure of the random variable $W$.

One important condition for an intermittency pattern to show up in the quark-gluon plasma evolution is related to the diffusion properties of the medium. Indeed, if the diffusion is fast, the "old" perturbations will be washed out and no superposition of fluctuations of different scales will remain possible. On the contrary, if the perturbations get reasonably
"frozen" some intermittency pattern, even weak, will remain as a signal of the stages of the quark gluon phase expansion. To our knowledge, the features of diffusion are not well known at this moment, and should be studied.

In any case, if intermittency patterns are observed in high multiplicity events of high energy collisions, this will give an interesting signal of the quark-gluon plasma formation. It has been checked, by comparison, that monte-carlo simulations of the background mechanisms of ordinary soft Nucleus-Nucleus interactions[5] lead to no variation of the factorial moments, together with the predicted elimination of the statistical noise. This is an encouragement to propose our method of analysis for the high multiplicity events which will be produced in the near future at Cern for heavy-ion collisions or in the large multiplicity fluctuations at pp colliders.

References


ABSTRACT

A brief general review of string and superstring theories of phenomenological interest is presented. It is intended mainly for the experimentalists.
The fact that some lectures on strings (and not only the hadronic ones) have managed to leak into a Moriond session devoted to QCD is another manifestation of the enormous interest in the subject among theoretical physicists. There are a few reasons to this:

1. Superstrings are at present the only consistent theories unifying gravity with gauge interactions [1].

2. As such they contain a quantum theory of gravity which is one-loop finite (perhaps finite to all orders of perturbation theory) and free of gravitational as well as gauge and mixed anomalies [2].

3. The consistent superstring theories show a great degree of uniqueness. Recently [3] more string theories have been constructed along lines similar to the heterotic string construction [4], but so far the most promising theory for phenomenology is the $E_8 \times E_8$ heterotic superstring. These theories have no freely adjustable parameters, except for an overall scale, to be identified with the Planck scale.

4. String theories have exhibited a remarkable number of miracles, namely one seems to be getting for free many properties which are not put into the theory in its basic construction. The mathematical origin for these miracles (some of them will be described below) is still unclear.

5. Unfortunately the uniqueness of the theory applies also to the embedding space-time. The heterotic string theory is ten-dimensional, and a Kaluza-Klein interpretation must be evoked, where six of the ten dimensions are compactified to a very small volume, the compactification radius being of the order of $10^{-33}$ cm.

A bosonic string is a generalization of the concept of a point-like particle, where the physical object is one dimensional. The first quantized theory of a point particle is derived from its generating functional defined on its line of propagation in space-time (or rather, all possible lines):

$$Z = \int_{\text{world-lines}} \mathcal{D}x^\mu e^{-\int d\tau}$$

(1)
Fig. 1: A point particle world-line

Similarly, the relativistic bosonic string propagates along its "world-sheet", depicted in Fig. 2 both for an open and a closed string

Fig. 2: The string world-sheet

The generating functional is [5]

\[ Z = \int D\sigma^{\alpha} \int Dx^{\mu} \ e^{-\frac{1}{4\pi\alpha'} \int_{S} \sqrt{g} g^{\alpha\beta} \partial_{\sigma} x^{\mu} \partial_{\sigma} x^{\nu} \eta_{\mu\nu} d^{2}\sigma + \ldots} \]

\( g^{\alpha\beta} \) is the world-sheet metric, \( \eta_{\mu\nu} \) is the embedding space metric.

The above expression is classically invariant under reparametrizations in both spaces and the additional terms involve the other embedding space background fields and a cosmological constant.

The superstring is (roughly) obtained by supersymmetrizing the above expression, thereby introducing a background gravitino and the supersymmetric partners of the space-time coordinates.

Interactions between strings are included when the sum over surfaces
in $\mathbb{Z}$ is allowed to extend over different topologies. For the closed string the genus of the surface (its number of handles) is equal to the order of perturbation theory ("number of loops") to which it contributes.

When this is done the first and perhaps the most important miracle of string theory follows: the two dimensional quantum gravity on the world-sheet (this is what eq. (2) really is) becomes an interacting theory of an infinite tower of states containing quantum gravity in the embedding space. The other states, lying on Regge trajectories, present a hierarchy of gauge symmetries.

For the superstring another miracle emerges at this point: the resulting 10-d field theory (with the Gliozzi, Scherk, Olive projection [6]) is also supersymmetric and the tachyon is eliminated.

So far, we have discussed the first quantized formulation of the string. In addition, the spectrum and the low-energy interactions are generally determined in the light-cone gauge.

A satisfactory covariant second quantized formulation exists only for the free string [7] and the field theory of the interacting string remains probably the outstanding open theoretical problem [8].

The phenomenologically interesting theories are:

- **type I**
- **heterotic**

As illustrated, heterotic string theory has as physical objects only oriented closed strings. From this it follows that the only possible interaction can be depicted as:

The type I superstring has more possible vertices, but both theories contain only one parameter, a scale.
\[ \alpha'(\text{slope}) \rightarrow M_s = \sqrt{1/\alpha'} \sim M_{\text{Planck}} \]

Gravity (coupling constant \( \kappa \)) and the gauge interactions (g) are unified, with the following relations:

\[
\begin{align*}
\text{type I} & : \kappa \sim g^2/\alpha' \\
\text{heterotic} & : g^2 \sim \kappa^2/\alpha'
\end{align*}
\]

Note that these are relations among 10-d coupling constants, that have the following dimensionalities: \([\kappa, g, \alpha'] = [t^4, t^3, t^2]\).

For energies much lower than \( M_s \), the effective field theory of the massless string states is 10-d \( N=1 \) supergravity [9,2], with a gauge group:

\[
\begin{align*}
\text{type I} & : SO(32) \\
\text{heterotic} & : SO(32), E_8 \times E_8
\end{align*}
\]

In ref. [4] more string theories are described, with other gauge groups but no space-time supersymmetry.

The 10-d supergravity contains a sugra multiplet: \( g^{\mu\nu} \) (graviton; 35 degrees of freedom), \( \psi^\alpha \) (gravitino; 56), \( B^{\mu\nu} \) (antisymmetric tensor field; 28), \( \lambda^8 \) (spinor; 8), \( \phi \) (dilaton; 1), and in addition the super Yang-Mills multiplet with gluons and gluinos.

The massive states of the string theory are important in rendering the supergravity theory finite and anomaly-free.

Unfortunately, the string theory is not yet at a stage where phenomenology can be done reliably from first principles only. Some speculations are necessary before predictions are made:

1) Compactification \( M^{10} \rightarrow M^d \times K \), where \( M^d \) is \( d \)-dimensional Minkowsky space and \( K \) is the internal space with \( 1/R_{\text{com}} \sim M_{\text{com}} \ll M_p \).

2) A stable hierarchy \( M_\nu \ll M_p \) probably requires supersymmetry in the effective 4-d theory obtained after compactification [10]. This, together with some mild assumptions, requires \( K \) to be of the Culabi-Yan type (Kähler, \( SU(3) \) holonomy) and the gauge group broken:

\[ SO(32) \rightarrow SO(26) \text{ with fermions in real representations (not interesting !)} \]

\[ E_8 \times E'_8 \rightarrow E_6 \times E'_8 \text{ with matter in 27 of } E_6 \text{ (possibly} \]

\[ \text{with additional matter in 54 and 105 of } E_6 \text{.} \]

\[ \text{and some further matter in 845 of } E_6 \text{.} \]
also singlets) and \( E_8' \) interacting with \( E_8 \) only gravitationally. This is the model we discuss in the following.

3) A further breaking of the gauge group to the standard model can be done via the Hosotani mechanism [11], but this requires at least an additional \( U(1) \) at low energies:

\[
E_8 \rightarrow SU(3) \times SU(2) \times U(1) \times U_E(1)
\]

If the additional group is indeed only \( U_E(1) \), the quantum numbers of all known particles with respect to this \( U_E(1) \) are fixed. The possibility that \( E_8 \) breaks down directly to \( SO(10) \) or \( SU(5) \) (no additional \( U(1)'s \)) is discussed in [12].

Experimentally, one can place some lower limits on such an additional \( Z_E \) boson. Its coupling constant is (theoretically, demanding grand unification with the other couplings) [13]:

\[
\alpha_E = g_E^2 / 4\pi = 0.0164 \quad \text{at} \quad E \sim M_w
\]

Ignoring mixing with the ordinary \( Z_0 \) and at equal mass, \( Z_E \) would be produced by \( pp \) at CERN with \( \sigma_E / \sigma_0 = 0.22 \pm 0.02 \). This ratio decreases for higher \( m_E \). The different quantum numbers under \( U_E(1) \) imply:

\[
B(Z_E \rightarrow e^+ e^-) / B(Z_0 \rightarrow e^+ e^-) = 1/3
\]

Unless the \( Z_E \) is just under the \( Z_0 \) peak (ruled out by neutral current experiments) the absence of experimental evidence for it at CERN implies \( m_E > 110 \text{GeV} \). Similar limits are obtained from neutral current data (neutrino interactions, parity violation in \( \ell \ell \) and atoms, and forward-backward asymmetry in \( e^+ e^- \rightarrow \ell^+ \ell^- \)).

4) Supersymmetry breaking can be achieved by a gaugino condensate in the hidden sector \( (E_8') \), whose effect propagates to the observable sector by gravitation [14].
The \( E_8 \) grand-unified families contain the following fields:

<table>
<thead>
<tr>
<th># of fields</th>
<th>SU(5) qu.nos.</th>
</tr>
</thead>
<tbody>
<tr>
<td>conventional ((q,q^c,l,l^c))</td>
<td>15</td>
</tr>
<tr>
<td>2 Higgs ((H,\tilde{H}))</td>
<td>10</td>
</tr>
<tr>
<td>2 Neutrals ((N,\tilde{N}))</td>
<td>2</td>
</tr>
</tbody>
</table>

Each particle coming with its supersymmetric partner. This implies some exotic particles still to be discovered.

The low-energy phenomenology requires the knowledge of \( K \), the manifold on which 6 dimensions compactify. So far Calabi-Yau manifolds have not been classified and this prevents a direct search for the one most suited for describing nature (if any). Given a choice of \( K \), one can in principle derive:

1. The number of \( E_8 \) generations [10]. Examples with 3, 4 or more families are known.

2. Yukawa couplings [15]. Many Yukawa couplings vanish due to discrete symmetries, for example.

3. The form of the \( d = 4 \) supergravity. There are indications that it is of the "no scale" type [16]. This is quite satisfactory because in those theories there is a natural generation of hierarchy. The usual "no scale" machinery gives typically (for a three generation model; masses are in GeV) [13]:

\[
\begin{align*}
\mu_q & \gtrsim 360 \\
\mu_g & \gtrsim 180 \\
\mu_l & \gtrsim 100 \\
\mu_\gamma & \gtrsim 27
\end{align*}
\]

One potential difficulty with superstring models (like any other supersymmetric GUT) is proton decay. The proton can decay either via \( X \)-boson exchange or via Yukawa couplings to squarks or color-triplet-Higgs. \( X \) boson exchange usually gives longer lifetimes in this kind of models, because the proliferation of particles slows the running of coupling constants and unification is at scales higher than the typical non-susy GUT (\( 10^{15} \text{GeV} \)). Yukawa couplings are much more dangerous and those potentially leading to proton decay must vanish for the model to make any sense. Luckily, as described above, some Yukawa couplings do vanish with Calabi-Yau compactifications.

Let me briefly mention some other phenomenological aspects of
superstring theory:

1. These models contain two axions. This may solve the strong CP problem, but those axions must be consistent with all the present limits: accelerator, reactor, astrophysical and cosmological.

2. The low-energy gravity is different from Einstein's. There are $R^2$ corrections which may be important for the existence (or non-existence) of singularities, and for cosmological solutions.

3. Big bang cosmology will be modified considerably due to the proliferation of states (string excitations) at temperatures of the order of $M_p$.

It is quite remarkable that after so much of concentrated theoretical effort, string theories which are so much constrained have not yet collapsed as candidates for a realistic theory of nature. Much more theoretical progress is needed before hard predictions are made, but the likely experimental signals at low energy could be: additional gauge boson(s), exotic particles, axions, and supersymmetric partners.

Acknowledgements

It has been a pleasure working with J. Ellis, K. Engvist, C. Gomez and D.V. Nanopoulos on some aspects of superstring phenomenology.

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8. For a first step in this direction see


15. A. Strominger, Phys. Rev. Lett. (1985) 2547. See also E. Witten, Ref. 11.

WHY IS A STRING THEORY ULTRAVIOLET FINITE?*

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Abstract

String theories have recently been regarded as candidates for the fundamental theory of nature. In this short review, an elementary exposition is presented on the key ingredients involved for having an ultraviolet finite quantum theory.

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I. INTRODUCTION:

String theories, developed in the late sixties and early seventies, have enjoyed a great recent comeback, after having remained dormant for more than a decade.\(^1\) They have been widely regarded as candidates for the ultimate unified theory of Nature, chiefly because they are capable of leading to finite theories of quantum gravity. Although they were originally formulated in an S-matrix framework, they may yet be transformed into some kind of "generalized" quantum field theories. Whatever the outcome, they certainly differ greatly from the conventional quantum field theories (QFT) which we have grown accustomed to in recent years. For physicists not working on the subject, particularly the experimentalists, the field seems mysterious. In this short review, I will present an elementary exposition on the key ingredients involved for having an ultraviolet finite quantum theory.

The ultraviolet behavior of a Feynman integral with loops can always be characterized by \(\int d^D p \cdot p^{-q}\), where \(D\) is the number of space-time dimensions and \(q\) is an integer depending on the number of propagators involved. Since one cannot avoid having diagrams with \(D > q\), QFT is meaningless without a certain amount of massaging. Making sense out of nonsensical mathematical expressions has been the preoccupation of field theorists for the past half century. In each case, one always manages to find a "taming operator" such that when it is applied to a Feynman integral, a finite number is obtained. Indeed, the history of QFT can be loosely described as that of taming the ultraviolet divergences.

Finding the correct taming operator requires the understanding of quantum physics. First of all, a renormalization procedure must be adopted. To proceed further, one finds that symmetry plays a crucial role. Starting with the gauge invariance of QED, next the non-abelian gauge theories, either for QCD or the standard electroweak theory, and more recently supersymmetry and supergravity, it has always been the symmetry which is chiefly responsible for the divergence cancellation. Identifying the relevant symmetry is therefore the key step involved in the construction of a taming operator.

The question remains whether this scenario can go on indefinitely within the context of the conventional QFT. The problem becomes particularly acute in regards to gravity—it appears that the usual procedure cannot render the theory finite; loop divergence cannot be removed. It could happen that a more clever introduction of new symmetries, hitherto unrecognized, would again lead to a finite quantum gravity within the conventional field theory framework, it is equally plausible that a new avenue of approach should be adopted. Indeed, string theory is the most likely candidate.
II. LOOP DIVERGENCE IN A PARTICLE FIELD THEORY:

To illustrate the problem, let us consider the simplest situation of a bosonic propagator, \( G(0, x) = \int dp (p^2 + m^2)^{-1} \exp(ipx) \), where \( dp = (2\pi)^{-d} d^d p \).

Graphically, we normally represent this propagator by a line segment, from the point 0 to the point x. This symbolic line segment certainly does not correspond to the actual trajectory of the particle; being a probability amplitude, it can be reexpressed as a sum over all allowed paths from 0 to x. This "world-line" representation becomes particularly useful for our present purpose if we consider the "loop" diagram, \( G(0, 0) = \int dp (p^2 + m^2)^{-1} \), which is always ultraviolet divergent.

Using the identity \( a^{-1} = \int ds \exp(-as) \) for the integrand and next performing the momentum integration, the propagator can be written as

\[
G(0, 0) = (1/2) \int D(s) W(s) P(s)
\]

where \( D(s) \) is the ordinary integration measure, \( ds \), with the integration domain \( 0 \leq s < \infty \), \( W(s) = (2s)^{-D/2} \) and \( P(s) = \exp(-m^2 s/2) \). This can be interpreted as a sum over closed paths according to their path lengths, s. For a fixed length s, the probability amplitude is \( P(s) \) and the density of paths is \( W(s) \). The UV divergence of the original Feynman loop integral has now been translated into the divergence of the world-line sum, (2), at \( s=0 \). That is, the UV divergence in this sum comes from the small loop contribution, as is expected intuitively. More interestingly, the zero-point (vacuum) energy of a point particle can also be expressed in terms of the propagator: \( V(m^2) = - \text{Tr} \ln G \), which is again a divergent one-loop integral. In the world-line representation, \( V(m^2) \) is again given by (1) if one substitutes the measure by \( D(s) = -ds/s \). Again, the UV divergence can be associated with the divergent contribution from small loops in the world-line representation.

Although these examples are of no immediate physical consequence, other divergences must be regulated in order to provide a meaningful quantum field theory. Using the vacuum energy as a typical example of a potentially divergent loop integral, we shall see next that in a string theory, due to the presence of a new symmetry, ultraviolet divergences are removed automatically.
III. MODULAR INVARIANCE AND ULTRAVIOLET-FINITENESS:

 Whereas a particle corresponds to a point, a string can be parametrized by a line segment, (a circle in the case of a closed string). As a closed string evolves, it traces out a cylindrical surface in the space-time continuum. Therefore, the vacuum energy for a closed string, in contrast to a closed path in the case of a particle, is characterized by a two-dimensional torus. Instead of a world-line sum, a "world-sheet" sum over all distinct tori can be used to represent its vacuum energy. Is the ultraviolet divergence present due to "short" tori?

 One must first learn how to label a torus. Given a parallelogram, a torus can be constructed by identifying the opposite sides. Since a parallelogram can be specified by two independent basis vectors \( \vec{w}_1, \vec{w}_2 \) in a plane, any function \( f \) over the torus can in turn be thought of as a periodic function on a two-dimensional Euclidean space such that \( f(\vec{r}) = f(\vec{r} + n_1 \vec{w}_1 + n_2 \vec{w}_2) \) with \( n_1, n_2 \) arbitrary integers. Furthermore, since a two-vector can equally well be represented by a complex number, we will treat \( \vec{r}, \vec{w}_1 \) and \( \vec{w}_2 \) as complex numbers in what follows. It is also clear that only the length ratio and the relative phase between \( \vec{w}_1 \) and \( \vec{w}_2 \) are relevant, which can be characterized by the ratio \( \tau = \vec{w}_2 / \vec{w}_1 \), with the convention \( \text{Im}(\tau) > 0 \). We therefore tentatively arrive at a conclusion that distinct tori can be specified by points in the upper-half complex plane.

 The vacuum energy for a free closed string can next be obtained, e.g., by summing \( V(s^2) \) over its particle spectrum. In particular, we can identify the parameter \( s \), up to a factor, with \( \text{Im}(\tau) \). Writing \( \tau = \tau_1 + i\tau_2 \), one arrives at an expression similar to (1)

\[
V(\text{string}) = -\frac{1}{2} \int D(\text{torus}) W(\tau_2) P(\tau)
\]

where \( P(\tau) \) can be expressed in terms of known functions. The integration measure for the torus can be chosen as \( d\tau_1 d\tau_2 / \tau_2 \). What remains to be specified is the domain of integration.

 The association of tori with points of the upper-half complex plane is not one-to-one. For instance, it is not difficult to convince oneself that changing the pair of vectors \( (\vec{w}_1, \vec{w}_2) \) to \( (\vec{w}_1, \vec{w}_2 + \vec{w}_1) \) yields the same torus. Therefore, there is a redundancy, described by the transformation \( \tau \mapsto \tau + 1 \). Provided that the integrand in (2) also remains invariant under \( \tau \), we can immediately restrict the integration domain to the strip \( \tau_2 > 0 \) and \(-1/2 \leq \tau_1 \leq 1/2 \). In this parametrization for the world-sheet sum, \( \tau_2 \) is analogous to the parameter \( s \) of (1), so that it characterizes the "length" of a torus.
A similar analysis shows that we also have a symmetry \( S: \tau \rightarrow -\frac{1}{\tau} \), which can be understood as changing the pair \((\tilde{u}_1, \tilde{u}_2)\) into \((\tilde{u}_1 - \tilde{u}_2, \tilde{u}_2)\). This allows us to further restrict \(|\tau| > 1\). Together with our earlier restriction, we arrive at the fundamental domain of integration for \( (2): |\tau| > 1, -1/2 < \tau < 1/2, 0 < \tau_2 < \infty \). In particular, the "zero-length" torus region has now been excluded and the vacuum energy is now ultraviolet finite!

The totality of transformations generated by \( S \) and \( T \) forms the "modular group". It is a symmetry group of any function defined over distinct tori. At a classical level, a string theory obviously possesses this property. However, this symmetry may not be preserved at a quantum level, a phenomenon we have grown accustomed to, generically referred to as an "anomaly". One of the most far-reaching findings of the string studies is the fact that this symmetry can be preserved at certain critical space-time dimension, e.g., for a bosonic string, \( D_{\text{critical}} = 26 \). Therefore, at our current level of understanding, for a finite string theory, one must force oneself to work at a space-time dimension different from four. Of course, to make contact with our more familiar four-dimensional world, additional step such as spatial compactification must be performed. This, however, will not alter the picture for UV finiteness which is our main concern here.

**IV REMARKS:**

More than a half century has passed since its invention, the divergence difficulty of QFT remains the central problem facing our attempt to understand the interactions of fundamental particles. Unlike previous efforts where UV finiteness was invariably achieved by discovering new symmetry in the context of conventional QFT, one finds a much more economical route by adopting a string picture. In this approach, a new kind of symmetry emerges naturally, which provides an automatic UV-cutoff in a self-consistent fashion.

It should be pointed out that I have painted too simplistic a picture for the overall finiteness of a string theory. First of all, even at the level of the vacuum energy for free strings, (2) in general remains ill-defined because of the presence of a tachyon, which leads to a divergence at the upper-limit of \( \tau_2 \)-integration. This can of course be turned around and used in further restricting the acceptable string models.

For a realistic unified theory, we must also bring in chiral fermions, which then leads to superstrings. This in turn forces one to be concerned with the chiral anomaly. This then takes us right back to the exciting
story of superstrings of last few years, which has just been reviewed by the previous speaker.

Shortly before the recent outburst in superstring activities, S. Weinberg, in his contribution, entitled "The Ultimate Structure of Matter", to a volume in honoring G. F. Chew's sixtieth birthday, addressed the old question in physics: What are the fundamental entities of which we regard our universe as being composed—particle or fields? After briefly clarifying the history of quantum field theory since its inception in years 1929 and 1930 by Heisenberg and Pauli, Weinberg carefully reviewed several swings back and forth between quantum field theory and the S-matrix approach, before finally arriving at the recent renaissance of QFT in the seventies. However, in spite of the impressive successes of modern QFT, Weinberg pointed out that the question what the ultimate form the fundamental theory would take remained unanswered. Weinberg concluded by stating "If our quantum field theories of which we're so proud are just the debris of some really fundamental theory which describes all of physics including gravity, it may be that the really fundamental theory will have nothing to do with fields; it may not look like a quantum field theory at all. I think we have to leave it as an open possibility that maybe, it will be something like an S-matrix theory." Could these prophetic words be fulfilled with superstrings?

References:

1) J. H. Schwartz, Phys. Reports 89, 223 (1982); "Superstrings" by J. H. Schwarz, World Scientific (1985), and references therein. See also the contribution by Schwartz in Ref. 3.


SUPERGRAVITY MODELS FROM SUPERSTRINGS

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Abstract

A brief review (mainly for the experimentalists) on how the four-dimensional supergravity models are obtained from the superstrings through a truncation mechanism is presented.
Superstring theories are a further step in attempting to unify all present-day known interactions. The most exciting feature of the superstring theories is the inclusion of gravitational interactions in a natural way. The construction of string theories is much more restrictive than that of pointlike particle field theories. Quantization of these theories is only consistent in certain space-time dimensions. The presence of tachyon states in bosonic string theories seems to require supersymmetry as the only possible way of curing the problem. The resulting theories predict the dimensionality of space-time to be ten instead of 4. Furthermore, the examination of anomalies and finiteness predicts the choice of the gauge group to be $O(32)$ or $E_8 \times E_8$. However, the theory still has a lot of problems. One of the problems is to study the low-energy phenomenology of superstring theories in four dimensions after compactification. Since to make low energy predictions from the string theory itself is rather difficult, one would like to have an approximation to the theory first to ease up the calculations. One of the approximations could be the field theory in the zero-slope limit, i.e. the field theory of the ground states of strings. Such an approximation has been considered by Green and Schwarz, which corresponds to a chiral $N=1$ supergravity multiplet coupled to an $N=1$ super Yang-Mills multiplet plus appropriate anomaly cancelling terms. However, a direct compactification of this approximate theory from $d = 10$ to $d = 4$ is still too difficult. Witten therefore introduced a way to mimic the compactification mechanism in order to obtain a four-dimensional theory which is termed as "truncation to 4 dimensions". In the following, we will review this truncation process (very briefly). The general structure of obtaining this 4-dimensional theory is shown in figure 1.

Superstrings (with gauge group $E_8 \times E_8$ or $O(32)$)

(zero-slope limit)

Chapline-Manton Model (with appropriate anomaly cancelling terms)

(compactification : Witten's method, truncation from 10-D to 4-D).

$4$-D $N=1$ supergravity + $N=1$ Super Yang-Mills

($M_4 \times K_6$, $M_4$ = 4-D Minkowski space-time,
$K_6$ = compact space of dimension 6)

FIG. 1
Let us remind ourselves the field content of 10-Dimensional (Chapline-Manton model) $N = 1$ sugra + $N = 1$ super Yang-Mills model. In the supergravity sector, there are 128 on-shell degrees of freedom. Its bosonic fields are: a metric $g_{MN}$, an antisymmetric tensor $B_{MN}$ and a real scalar $\phi$, where $M, N = 1, ..., 10$. Its fermion fields are: a gravitino $\Psi_M$ and a spinor $\lambda$. In the super Yang-Mills sector, there is a gaugino spinor $\chi^a$ corresponding to each gauge field $A^a_M$ of the gauge group $G$ ($G$ is, for example, $E_6 \times E_8$).

**Truncation to 4-dimensions**

We will begin our discussion by considering first the supergravity sector.

(1) As a first step, we assume all fields to be independent of the last six coordinates, i.e. $X^M$, $M = 5, ..., 10$. Result: an $N = 4$ supergravity theory in four-dimensions. The compact space $K_6$ is a six-torus.

(2) To further reduce to $N = 1$ supergravity, one notice the following: the rotation group of $X^M$ is $O(6)$ or $SU(4)$. The four supersymmetry $Q^a$ transform as a 4 of $SU(4)$. One then picks an $SU(3) \times$ subgroup of this $SU(4)$. The $Q^a$'s transform as $1 + 3$ under this $SU(3) \times$. One then choose the singlet as the only supersymmetry left. Result: an $N = 1$ supergravity theory in 4-dimensions.

The inclusion of the Super Yang-Mills multiplet is straightforward and one can again use the arguments above. A more interesting case is to require an invariance under a diagonal group of one of the $E_8$'s. The resulting field content after this truncation mechanism is shown in figure 2.

<table>
<thead>
<tr>
<th>Before &quot;Compactification&quot;</th>
<th>After &quot;Compactification&quot;</th>
</tr>
</thead>
<tbody>
<tr>
<td>Supergravity sector : $g_{MN}$</td>
<td>Supergravity Sector : a(2, 3/2) multiplet (graviton, gravitino)</td>
</tr>
<tr>
<td>$B_{MN}$</td>
<td>$\phi$</td>
</tr>
<tr>
<td>$\Psi_M$</td>
<td>Two (1/2, 0) multiplets (scalars $\phi$ and $\sigma$, pseudoscalars $\theta$ and $\eta$, and their fermion partners)</td>
</tr>
<tr>
<td>$\lambda$</td>
<td></td>
</tr>
<tr>
<td>$A^a_M$</td>
<td>One (27) family of $E_6$ and its scalar partners.</td>
</tr>
</tbody>
</table>

**FIG. 2**
Summary

This truncation mechanism suggested by Witten certainly preserves some properties of the 10-Dim theory after compactification but oversimplifies the whole picture. For example, it preserves

(1) an \( N = 1 \) supersymmetry
(2) an \( E_6 \times \bar{E}_6 \) gauge symmetry
(3) has a flat potential (to describe the smallness of the masses of the light fields from Planck mass)
(4) and keeps the classical scaling property. (The dilatons \( \phi \) and \( \sigma \)).

However, it has

(1) only 1 (27) family of \( E_6 \) and
(2) does not give rise to the extra light singlets which associate with the deformation of the shape of the manifold.

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BAGGED SUPERFIELDS AND LIGHT PSEUDO - GOLDSTONE FERMIONS

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Abstract

By means of a SUSY generalization of Bogoliubov - MIT confinement we restrict supermultiplets of particles to small regions of space.

SUSY PCAC is used to introduce pseudo Goldstone particles living outside the bag. Supersymmetry is automatically broken so that the pseudo Goldstone fermions are much lighter than the bosons.
It is not completely straightforward to describe quarks and leptons as composite particles.

If we form the product of mass \( m \) and extension \( R \) of bound states that we are used to, we find a product \( mR \) which is bigger than one. For humans it is \( \sim 10^{44} \), for atoms \( \sim 10^6 \), for nuclei \( \sim 10^2 \) and for nucleons \( mR \sim 5 \). (For pions which is exceptional \( mR \) is slightly smaller than one). These results are natural consequences of Heisenberg's uncertainty relation

\[
\Delta x \cdot \Delta p \geq \frac{\hbar}{2}
\]

Even a massless particle will have an energy \( E = p \) which is of order of \( R^{-1} \) if it is confined to a region of extension \( R \).

One of the most beautiful results in physics is the fantastic agreement between the observed magnetic moment of the electron and its calculated value from QED. The assumption of a point electron give the correct results with 11 digits. If the electron has an extension \( R \) then \( R \lesssim 5 \times 10^{-17} \text{cm} \).

The energy due to confinement of a preon inside such a distance is something as 200 GeV and not \( 0.5 \times 10^{-3} \) GeV which is the energy of an electron at rest. If the electron with \( mR \lesssim 10^{-5} \) or the neutrino where probably \( mR \lesssim 10^{-10} \) are composite systems they are definitely of a kind that is different from ordinary bound states. Among the hadrons only the pion could serve as a model. The small mass of the pion is most easily explained by looking at it as a quasi Goldstone particle generated by a spontaneous breaking of global SU(2) x SU(2) chiral symmetry. Naive \( q \bar{q} \) models for the pion often has trouble to get the pion light enough, the Goldstone picture represents it as a collective excitation with an indefinite number of \( q \bar{q} \) pairs.

In a supersymmetric extension of the Nambu-Goldstone mechanism the Goldstone bosons will also have associated fermionic partners\(^1-3\). This is the most elegant way I know of to obtain very light or even massless composite fermions.

What is not so pleasant however is that the nice supersymmetry that was introduced to get fermions as light as bosons has to be broken if we want to make a model of reality. And it has to be broken such that the bosons become much heavier than the fermions.

I shall describe just this last kind of symmetry breaking in a model very similar to the MIT bag model. The bagged particles\(^4\)
will be called urons and will be represented by superfields confined to a small sphere of radius $R$ where $R$ is extremely small. Uron bags have an energy $\sim R^{-1}$ and have themselves nothing to do with quarks and leptons. The minimum number of constituents in a uron bag is two, we can imagine some kind of urcolour to be the dynamical mechanism that leads to the confinement. The energy of the bagged urons of spin $1/2$ will be determined by the Bogoliubov-MIT boundary condition

$$\left. -i \hat{r} \varphi \right|_{r = R} = \varphi$$

We make a supersymmetry transformation of this to find the boundary condition for scalar urons:

$$\mu_r \varphi = \frac{d\varphi}{dr} \left. \right|_{r = R}$$

$\mu_r$ is the uron mass.

As we want to have Dirac particles we need a minimum of two chiral superfields: $\varphi_L$ and $\varphi_R$. Each of these chiral superfields also carry an index characterizing the transformation property under a global flavour symmetry group $G$ of the uron Lagrangian inside the bag. To each generator $t_a$ of this group there corresponds super-currents

$$V_a = \varphi_L^+ t_a \varphi_L - \varphi_R^+ t_a \varphi_R$$

and

$$A_a = \varphi_L^+ t_a \varphi_L + \varphi_R^+ t_a \varphi_R$$

These supercurrents, transforming under the adjoint representation of the group $G$ (=SU?) fullfil conservation equations (inside the bag)

$$D^2 V_a = D^2 A_a = 0$$

$$D^2 A_a = 8 \mu_r \varphi_L^+ t_a \varphi_L$$

$$D^2 A_a = 8 \mu_r \varphi_R^+ t_a \varphi_R$$
By expanding the chiral superfields in number and \( \theta \) components one find that \( V_a \) and \( A_a \) contain the ordinary vector and axial currents as well as currents with spinorial nature.

For vanishing uron masses the currents \( A_a \) are strictly conserved inside the bag, but they have (as in the MIT bag where quarks are confined) a delta function source term on the surface. This is because the boundary condition breaks chiral symmetry.

In the MIT bag one uses PCAC and ensures the continuity of the isovector axial currents by introducing an external Goldstone pion field coupled to the bag surface\(^5\). By a supersymmetric extension of PCAC - SUSY PCAC, we ensure the continuity of the axial supercurrents \( A_a \) and their divergences\(^8\) through the uron bag surface by introducing chiral superfields on the outside of the bag.

If urons are massless so are the Goldstone fermions and bosons associated with the long distance properties of the axial currents. If we give the urons a mass \( m_{ur} \) we get relations between \( m_{ur} \) and the masses of the pseudo Goldstone fermions \( \tilde{\nu}_f \) and bosons \( \mu \).

As the boundary conditions break supersymmetry \( \mu_f \) and \( \mu \) are different\(^4\):

\[
\begin{align*}
\mu_{ur} &= \frac{1}{2} \mu_f \quad \text{(7)} \\
\mu_{ur} &= \frac{\mu}{2} \frac{UR(1+UR)}{2+2\mu R+\mu^2 R^2} \quad \text{(8)}
\end{align*}
\]

For a very small radius \( R \) of the uron bag \( (\mu R \ll 1) \) one has

\[
\mu = \sqrt{\frac{4\mu_{ur}}{R}}, \quad \tilde{\nu}_f = \frac{R}{2} \mu^2 \quad \text{(9)}
\]

The pseudo Goldstone fermions are therefore much lighter than the bosons: this was the desired result. We note from eq 8) that if \( R \to \infty \) then \( \mu \to \tilde{\nu}_f \) and it can be shown that supersymmetry is recovered. As seen from eq.9 it is not possible to make a fermion massless without getting its SUSY boson partner massless. If one wants to build a realistic theory for quarks and leptons from this model then neutrinos necessarily have a mass. Another general property regardless of the group \( G \) is that quarks and leptons must transform under the adjoint representation of this internal flavour symmetry group of the urons.

The energy of uron bags and the energy scale \( \Lambda \) relevant for the...
confinement of quarks is proportional to $R^{-1}$. Therefore it is only at extremely high energy that the fundamental dynamical variables which are the quark fields will come into play. Today their only manifestation would be through the effective interaction of the pseudo Goldstone fermions arising from the simultaneous breaking of supersymmetry and chiral symmetry.

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STRUCTURE FUNCTIONS AND PARITY VIOLATION IN SUPERSYMMETRIC QCD

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Abstract: Two topics in supersymmetric QCD are shortly reviewed: the structure functions in unpolarized and polarized lepton-nucleon scattering and the influence of possible parity violation on the deep inelastic scattering.

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1. In this talk, I discuss selected topics in deep inelastic scattering assuming that supersymmetry\(^1\) shows up in this process. In our phenomenological approach we do not follow any particular prescription of how to introduce and break supersymmetry, partly because of lack of satisfactory model, we assume instead that the resulting theory of strong interactions contains, in addition to quarks and gluons, new elementary fields – the spin \(1/2\) gluinos \(\lambda\) and spin \(0\) squarks \(s\) – which masses are parameters somewhere in the region not excluded by the present supersymmetry searches. Since the lagrangian of supersymmetric quantum chromodynamics (SQCD) has been presented already in the preceding talks (and can be easily found in the literature\(^2\)) I save space and time referring to them. Let me start with commenting the results in electron-unpolarized proton scattering.

2. The cross-section for the scattering of electron of momentum \(k\) and helicity \(\pm\) off the proton of momentum \(P\), with the exchange of momentum \(q\), reads

\[
d\sigma^{\text{ep}}/dxdy = 4\pi\alpha^2 s/Q^4 \left[ xy^2 F_1(x,Q^2) + (1-y)F_2(x,Q^2) \pm xy(1-y/2)F_3(x,Q^2) \right]
\]

where \(s = (k+P)^2\), \(Q^2 = -q^2\), \(x = Q^2/(2P.q)\), \(y=q.P/k.P\) and \(\alpha = e^2/4\pi\) is the strong coupling constant.

The functions \(F_i, i = 1,2,3\) contain in general the structure functions \(F_i\) resulting from the exchange of the photon \((\gamma\gamma\) term), \(Z^0\) \((ZZ\) term) and from the interference \((\gamma Z\) term)

\[
F_i = F_i^{(\gamma\gamma)} + \alpha F_i^{(\gamma Z)} + \beta F_i^{(ZZ)}
\]
with $\alpha, \beta$ depending on $Q^2$ (through the $Z^0$ propagator) and the lepton electro-weak couplings. (Their explicit form may be found in ref.3.). To see how supersymmetry enters the structure functions one has to express them in terms of parton densities:

$$F_1 = \frac{1}{2} \sum_{i=1}^{n_f} A_i (q_i + \bar{q}_i)$$

$$F_2 = x \sum_{i=1}^{n_f} A_i (q_i + \bar{q}_i + s_i + \bar{s}_i)$$

$$F_3 = \sum_{i=1}^{n_f} B_i (q_i - \bar{q}_i)$$

with $A_i, B_i$ depending on the $Z^0$ propagator and the couplings of leptons and partons ($B_i = 0$ when only photon is exchanged). The appearance of supersymmetry is twofold. First, the superpartner fields enter the structure functions explicitly, leading e.g. to the violation of the Callan–Gross relation $F_2 = 2 \times F_1$, second, it modifies the quark densities in the course of $Q^2$ evolution. The later, when framed in terms of the evolution equations 4) means that the quark and gluon equations get modified by the new splitting functions (quark–squark–gluino, gluon–gluino–gluino and gluon–squark–squark) and two additional equations (for gluinos and squarks) couple. Out of many detailed results 4) the most spectacular seems to be the decrease with $Q^2$ of the number of valence quarks form 3, below the squark production threshold, to 2 at infinite $Q^2$. Instead, the sum of valence quarks and squarks is conserved. When looking at moments of the structure functions (3), which are directly measurable, similar behaviour is observed e.g.
$$\int_0^1 F_3(x, Q^2) dx = \begin{cases} \text{constant in QCD} \\ \text{decreasing in SQCD} \end{cases}$$ 

as illustrated in Fig. 1.

![Graph](image)

Fig. 1 – The sum rule (4) as function of $Q^2$ in QCD (broken line) and SQCD (solid line). The squark mass is chosen at 25 GeV.

3. Defining the polarized proton $\Delta P = 1/2(P_+ - P_-)$ as the difference of helicity $+$ and $-$ states one can write the lepton-polarized proton cross-section

$$d\sigma^{\pm\Delta P}/dxdy = 4\pi\alpha^2 s/Q^4 \left[ \pm 2xy(1-y/2)G_1(x, Q^2) + xy^2 G_3(x, Q^2) + (1-y)(G_4(x, Q^2) + G_5(x, Q^2)) \right]$$

The functions $G_1, G_3 - G_5$ contain the spin-dependent structure functions $G_4, G_5$ from $(\gamma\gamma), (ZZ)$ and $(\gamma Z)$ contributions. In terms of parton spin densities $\Delta q = q_+ - q_-$ and $\Delta s = s_R - s_L$ (the squark $s_R(s_L)$ interacts with quark $q_+(q_-)$) one writes $^{3}$
The QCD relation $6_4 + 6_5 = 2 \times 6_3$ is in general broken in SQCD. The $Q^2$-evolution of the spin densities is again slightly more complicated than in standard QCD due to the appearance of new fields. The most interesting physically is the nonconservation of the spin carried by the valence quarks. Constant in QCD it decreases, in most cases to zero, in SQCD. The spin of the proton is thus carried by the sea components. Analogically, the sum rules on the spin-dependent structure functions, known in QCD, are all broken in SQCD. For instance, the celebrated Bjorken sum rule

$$
\int_0^1 \left[ G_1^{(\gamma \gamma)P} - G_1^{(\gamma \gamma)N} \right] dx = \begin{cases} 
\text{constant in QCD} \\
\text{decreasing in SQCD}
\end{cases}
$$

Analogically behave the $(\gamma Z)$ and $(ZZ)$ parts of $G_1$, as seen in Fig. 2, as well as the sum rules on $G_3$ and $G_4 + G_5$.

Fig. 2 - The same as in Fig. 1 but for the Bjorken sum rule (6).
To summarize, the effects of supersymmetry are present in deep inelastic scattering above the threshold for superpartner production, in general however are rather weak and appear very slowly with $Q^2$. In the search for supersymmetry in the inclusive $eP$ scattering one is forced to look for more refined effects. Below I present an example of this kind.

4. Supersymmetric QCD allows for parity violation. This may happen when the squarks $s_L$ and $s_R$ have nondiagonal mass matrix and the mass eigenstates $s_1$ and $s_2$ have different masses. In terms of the $s_L$, $s_R$ mixing angle $\Theta$, parity violation occurs for $\Theta \neq 45^\circ$. This can be seen by looking at the quark–gluino–squark interaction lagrangian in which the coupling of e.g. the squark $s_1$ is different to right- and left-handed quark

$$L_{\bar{q} \lambda s} = \sqrt{2} g [\cos \Theta \bar{\ell}_a T^a q_L s_1^* - \sin \Theta \bar{\ell}_a T^a q_R s_1^* + ...], \tag{7}$$

$g \cos \Theta$ as compared to $g \sin \Theta$ ($T^a$ are the SU(3) group generators). This means precisely parity violation.

The effect can be observed in deep inelastic scattering since, in the first order in $g$, a nonzero structure function $F_3(x2)$ appears, due to the diagrams of Fig.3. To measure it one has to select $F_3(x,Q^2)$ by scattering the combination ($e^+_+e^-$) or ($e^-_+e^-_+$) on protons and compare the result with the standard model prediction in which only $F_3^{(1)}$ and $F_3^{(2)}$ are nonzero. If the squarks mix nonideally, as suggested by supergravity models, the effect of SQCD may be visible.
Fig. 3. The diagrams which produce nonzero $F_{3}^{(\tau 2)}$ in SQCD for nonideal squark mixing.

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ABSTRACT

The low-energy values of the gauge coupling constants are very insensitive to those at a large energy scale if the corresponding theory is asymptotically divergent. This in general requires more than the standard three generations of quarks and leptons. We explore this scenario within the extension of the standard model provided by $N = 1$ supersymmetry with $n$ quark and lepton generations. We introduce as a parameter the mass threshold $M_{ss}$ of supersymmetric partners of ordinary particles. We find that, if all gauge theories enter the non-perturbative regime at an energy of the order of $10^{17}$ GeV, the low-energy values of gauge coupling constants are well reproduced with five quark and lepton generations and $M_{ss} < 5$ TeV. Furthermore, all fermion masses should be below 250-300 GeV.
A large fraction of theorists is becoming accustomed to the idea that superstrings at an energy scale of the order of the Planck mass will provide the ultimate theory of the world. However, there is a scenario according to which the values of the gauge coupling constants of the SU(3) × SU(2) × U(1) group at low energy (i.e., at about the Fermi scale) are largely insensitive to those at a very large scale and, therefore to some extent independent of our understanding of the physics up there\(^1\). This can happen if all known gauge theories behave at energies larger than the Fermi scale as asymptotically divergent theories and not as asymptotically free. This is simple to understand: let us consider the relation provided by the renormalization group between the gauge couplings at a low-energy scale (\(\mu\)) and those at a larger one (\(\Lambda\))

\[
\frac{4}{\alpha(\mu)} = \frac{4}{\alpha(\Lambda)} + C \frac{\ln(\Lambda/\mu)}{\mu}
\]

(1)

The above expression is only correct at one-loop level but still suitable for our illustration purposes. The sign of the constant \(C\) in Eq. (1) determines whether the theory is asymptotically free (minus) or not (plus). In the first case, the value of the gauge coupling at low energy (\(\alpha(\mu) \gg \alpha(\Lambda)\)) is obtained as the difference of two large numbers: \(1/\alpha(\Lambda)\) and \(\ln(\Lambda/\mu)\). Indeed \(\alpha(\Lambda) \ll \alpha(\mu)\) from which follows:

\[
\frac{4}{\alpha(\mu)} \ll \frac{4}{\alpha(\Lambda)} \quad \text{or} \quad \ln(\Lambda/\mu) \ll \frac{4}{\alpha(\mu)}
\]

(2)

Any uncertainty in the value of \(\alpha(\Lambda)\) will therefore be amplified when it is propagated to \(\alpha(\mu)\). The opposite is true for divergent theories where \(\alpha(\mu) \ll \alpha(\Lambda)\) and therefore \(1/\alpha(\mu) \gg 1/\alpha(\Lambda)\), \(1/\alpha(\mu) \sim O(c\Lambda/\mu)\). In this case the value of \(\alpha(\mu)\) is essentially determined by the second term in the right-hand side of Eq. (1) if one can neglect, with respect to it, \(1/\alpha(\Lambda)\), i.e., if the theory is entering a non-perturbative regime (\(\alpha(\Lambda) \sim O(1)\)). The constant \(C\) appearing in Eq. (1) depends upon the gauge group and the particle content of the theory.

At "our" energies (of the order of the W mass) the SU(3) and SU(2) sectors are asymptotically free: one needs new particles in order to obtain the behaviour of an asymptotically divergent theory at large energies. The demand in terms of new quark and lepton generations within the standard model is very high (a total of 8-9 generations). A more economical solution is provided by the \(N = 1\) supersymmetry version of the standard model when only a total of five quark and lepton generations is required\(^2\). Indeed, in this case, one gets "for free"
the superpartners of ordinary particles. This also provides a solution of the well-known hierarchy problem. Without supersymmetry, the corrections to the curvature of the Higgs potential are of the form:

$$m^2 = m_0^2 + \alpha \mathcal{C} \Lambda^2$$  \hspace{1cm} (3)

with $\mathcal{C}$ a model-dependent constant, and extraordinary fine tuning is required to keep $m$ as low as the Fermi scale:

$$m \simeq \langle H \rangle \simeq \mathcal{G}_F^{-1/2} \ll \Lambda$$  \hspace{1cm} (4)

Supersymmetry turns Eq. (3) into:

$$m^2 = m_0^2 + \alpha \mathcal{C} \Delta M^2$$  \hspace{1cm} (5)

where $\Delta M$ is the mass gap between normal particles and their supersymmetric partners. With:

$$\Delta M \leq (\alpha \mathcal{G}_F)^{-1/2} \simeq 3 \text{ TeV}$$  \hspace{1cm} (6)

the inequality in Eq. (4) is not spoiled by higher order corrections.

The purpose of this lecture, which is based on a work done in collaboration with L. Maiani, is to analyze the low-energy values of gauge couplings when all gauge theories are divergent at a large energy scale within the $N = 1$ supersymmetric extension of the standard model. In particular, we consider an $SU(3) \times SU(2) \times U(1)$, supersymmetric gauge theory with $n$ supersymmetric generations and two Higgs chiral supermultiplets. We assume that supersymmetry is broken by soft mass terms, which split squarks and sleptons from their fermionic counterparts and give masses to the gauginos. All these mass terms are independent of each other, but, for simplicity, we will assume that they are of the same order of magnitude. Thus, particles and gauginos start contributing to the $\beta$-functions above a single threshold, $M_{SS}$. Quarks and leptons take masses from the Yukawa couplings to the Higgs doublets. The latter couplings cannot diverge too early, which gives a general bound of the order of 200-250 GeV to quark and lepton masses. To be definite, we assume that quarks and leptons of the yet unseen, $n-3$, generations have a mass above the $W$-mass and smaller than the "Fermi scale" $\Lambda_F$:

$$\Lambda_F \equiv 250 \text{ GeV}$$  \hspace{1cm} (7)
Similar bounds apply to the Higgs scalar masses. Excluding ad hoc cancellations, this implies the whole Higgs supermultiplet to be below $\Lambda_F$.

Present experiments determine the values of the gauge couplings at a mass scale of the order of $M_W$. For colour interactions, we take:

$$\alpha_3(M_W) = 0.12^{+0.04}_{-0.02}$$ \hfill (8)

corresponding to

$$\Lambda_{QCD} = 150^{+150}_{-100} \text{MeV}$$ \hfill (9)

and

$$\alpha_{\text{e.m.}}(M_W) = 0.00772$$ \hfill (10)

The latter value results from the average of the CDHS\textsuperscript{6}) and CHARM\textsuperscript{7}) results:

$$\sin^2 \theta_W(M_W)_{\text{CDHS}} = 0.227 \pm 0.005 \pm 0.005$$ \hfill (12)

$$\sin^2 \theta_W(M_W)_{\text{CHARM}} = 0.236 \pm 0.005$$

Evolution of the above values from $M_W$ to 250 GeV ($\Lambda_F$) has been computed with three or five quark and lepton generations, and the mean value of the two results has been taken as a reference value, to be compared with the $\alpha_i(\Lambda_F)$, $i = 1, 2, 3$, given by our calculation.

The reference values should be attained - within the errors - by the coupling constants running "backward" from the large energy scale down to $\Lambda_F$. Before they get to $\Lambda_F$ they pass the mass threshold $M_{ss}$ below which the supersymmetric partners do not contribute anymore to the beta function. Therefore, their values at $\Lambda_F$ depend upon $M_{ss}$. The evolution has been done numerically using formulas which are correct up to two loops. The values of the coupling constant at $\Lambda$ can be varied between 1 and 10 without affecting significantly the results. The relevant parameters are $n$ (the number of generations), $\Lambda$ the large energy scale and $M_{ss}$. $n$ is essentially fixed to the first two integers for which colour
interactions above $M_{ss}$ are not asymptotically free to one loop, $n = 5, 6$. Larger values would make colour interactions to diverge too early, for $M_{ss} \ll \Lambda$. For $n = 5$ or 6, we have varied independently $\Lambda$ and $M_{ss}$, in the range:

$$10^{14} \text{GeV} < \Lambda < 10^{19} \text{GeV}$$

$$\Lambda_{F} < M_{ss} < 10^{-7} \text{GeV}$$

Our results for $n = 5$ generations are illustrated in Figs. 1 to 3. A good, simultaneous fit to the gauge couplings is obtained for:

$$\Lambda = 1 \cdot 10^{16} \text{GeV}$$

$$M_{ss} = 2 \text{ TeV}$$

The no-threshold situation $^4$, $M_{ss} = \Lambda_{F}$, still leads to an acceptable solution. Increasing $\Lambda$, the curve giving $\sin^2 \theta_W$ as a function of $M_{ss}$ remains quite stable, while $\alpha_3$ and $\alpha_{e.m.}$ shift both downwards. As their steepness is quite different, the triple coincidence is lost at some value of $\Lambda$. The opposite effect on $\alpha_3$ and $\alpha_{e.m.}$ is produced when lowering $\Lambda$.

In conclusion, for $n = 5$, we can reproduce satisfactorily the values of the gauge couplings, for:

$$10^{16} \text{ GeV} \leq \Lambda \leq 5 \cdot 10^{16} \text{ GeV}$$

and for

$$M_{ss} \leq 5 \text{ TeV}$$

No lower bound to $M_{ss}$ results.

We find it quite satisfactory that a good fit to the low energy couplings requires a value of the mass-gap of $N = 1$ supersymmetry of the order given in Eq. (16), and well consistent with the value of $M_{ss}$ needed to stabilize the low energy scale, Eq. (6).

The case $n = 6$ also admits a solution but only for considerably larger values of $M_{ss}$.
\[ M_{ss} \sim 10^3 \text{TeV} \]  
(17)

The reason is that colour interactions diverge more rapidly above \( M_{ss} \). This must be compensated by delaying the mass-scale where all particles contribute, i.e., by increasing \( M_{ss} \). In contrast with the previous case, the value Eq. (17) seems too large to be reconciled with the bound in Eq. (6).

In conclusion, the idea that gauge interactions are not asymptotically free at large energy, but rather diverge at a common energy scale \( \Lambda \) and that \( N = 1 \) supersymmetry holds below, leads to a good fit to the low-energy gauge couplings for five supersymmetric generations. At the same time, a very significant upper bound can be put on the mass gap of \( N = 1 \) supersymmetry, \( M_{ss} \), less than a few TeV. This is remarkably consistent with what is needed for the hierarchy problem, and very promising for the next generation of accelerators. The range of values which we find for \( \Lambda \) is close enough to the Planck mass, not to exclude that the strong regime of the gauge interactions is related to a further unification which includes gravity.
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FIGURE CAPTIONS

Fig. 1 The strong coupling constant, $\alpha_3$, at $A_F = 250$ GeV as a function of the
supersymmetry threshold $M_{\tilde{g}}$ for five generations of quarks and leptons.
Values of the ultra-violet cut-off, $\Lambda$, are indicated. The shaded region
represents the range allowed by present experimental data; see the
Table.

Fig. 2 Same as Fig. 1, for the electromagnetic coupling, $\alpha_{e.m.}$.

Fig. 3 Same as Fig. 1, for the weak mixing parameter, $\sin^2\theta_W$.  
- Figure 3 -
CONCLUSIONS

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Despite the title of this talk, before getting to the conclusions, I would like to review a number of topics which were presented at the Conference.

1. HARD HADRON PHYSICS

From data on many hard reactions we derive impressive evidence that the parton model and QCD work well. We have to be aware, however, that the evidence for QCD is only semi-quantitative.

1.1. Direct Photons.

Direct photons are there and will possibly provide a handle to study strong interactions at short distances even at the highest Collider energies. SPS data are meager, still they indicate the phenomenon quite well. As done in the past by R806-807 at the ISR, NA24 shows the angular distributions of hadrons around a) two-converted photons of a π0, and b) one converted photon at transverse momentum and distance from detector such that the chance of missing a second one from π0 or η decay would be very small. This is shown in fig. 1.1. The absence of hadrons at small angles to the trigger in the single photon sample, as well as the small accompanying jet in the two photon spectrum is precisely what QCD would predict for direct photons and for π0s in a jet respectively. This consolidates the calculations indicating a limited background at these transverse momenta (\( P_T > 4 \text{ GeV} \)), and one is thus justified to try to check QCD more quantitatively. One way of doing this is to compare direct photon production in reactions with and without leading antiquarks in the primary hadrons. Two first order QCD processes, Compton and annihilation,

enter with different weights in these reactions, the annihilation graph being much more important at large \( x_T \) when leading antiquarks are available.

R808 compares \( \overline{p}p \) and pp at the ISR where at \( P_T \sim 6 \text{ GeV/c} \) an excess of ~ 23% due to the annihilation graph is expected in the ratio \( \gamma / \pi^0 \). The data is shown in fig 2a, and unfortunately the statistics is not sufficient to give a significant answer. The same conclusion is also reached when looking at the NA24 and WA70 data in figs. 2b and
2c. This is particularly disappointing since in this case one compares $\pi^-$ and $\pi^+$ direct photon production, where at $x_T = 1$ one expects (1st order QCD) the $\pi^-$ rate to be larger than the $\pi^+$ rate by as much as a factor of two.

UA2 reports an interesting measurement at the SPS collider. At these energies the $\gamma / \pi^0$ ratio is too small to be measurable at $p_T < 10 \text{ GeV}$. At larger $p_T$'s photon pairs from $\pi^0$, $\eta$ decay are not separated in space in the calorimeter, and one must rely on their different conversion probability in a given layer of material in order to see the difference from single photons. This so-called indirect method is less powerful in signalling photon pairs, still it can work if the $\gamma / \pi^0$ ratio is large enough, and indeed it works for UA2. Fig. 3 shows as full curves the computed conversion probabilities $\epsilon_\gamma$ for single photons and $\epsilon_\gamma$ for photon pairs, which are different by $\sim 25\%$. The data points indicate the measured conversion probability for unaccompanied showers - the direct photon candidates - and for showers with hadron tracks at small distances. Clearly the two samples behave differently, very much as expected for single photons and $\pi^0$ or $\eta$'s. The excess over the computed probabilities can be attributed to some $\pi^0$ contamination in the sample of single photon candidates, and to some contamination of $\pi^0$ pairs etc. in the $\pi^0$ sample. A best-fit to the data can be made in terms of the $\gamma / \pi^0$ ratio as follows. Call $\alpha$ the observed conversion probability for the direct photon candidates. The amount of background can be indicated by the parameter

$$b = \frac{\alpha - \epsilon_\gamma}{\epsilon_\gamma - \epsilon_\gamma},$$

which is 1 if all direct photon candidates are photon pairs and 0 if all direct photon candidates are single photons. UA2 finds in the central calorimeter

$$b = 0.23 \pm 0.09 \text{ (stat.)} \pm 0.04 \text{ (syst.)}$$

at $10 \leq p_T \leq 40 \text{ GeV}$.

The parameter $b$ is found to decrease with increasing $p_T$, so that one would expect $\gamma / \pi^0 \sim 1$ at $p_T \sim 60 \text{ GeV}$. All in all, the data seem to indicate that the signal rate decreases slowly with $p_T$, such as to be still measurable at unexpectedly large $p_T$. Since the $\pi^0$ rate apparently decreases faster, one might eventually be able to isolate large $p_T$ photon events with small background. For instance at $\sqrt{s} = 2 \text{ TeV}$ (Fermilab Collider) one might hope to be able to measure the direct photon cross-section at $30 \leq p_T \leq 90$ GeV, with in addition a rather clean sample of single photon events at $p_T \geq 60 \text{ GeV}$. The direct photon and $\pi^0$ inclusive cross-section as measured in UA2 up to $p_T \geq 45 \text{ GeV}$ are shown in fig. 4.

1.2. Jets.

Interesting results on jet production at large $p_T$ were reported by UA1 $^2$) and UA2 $^3$). We should always be aware of the importance of this beautiful and simple phenomenon that
dominates at large transverse energy. Fig. 5 shows the dependence of event sphericity on $E_T$ in UA2. The onset of the new regime is extremely fast, between $\sim 40$ and $\sim 70$ GeV. Can this plastic evidence for the parton structure of the proton be exploited to provide a firm test of QCD? Here is where things become difficult. One way would be to test if the energy dependence of the inclusive jet cross-section between the Collider energies $\sqrt{s} = 546$ and $\sqrt{s} = 630$ GeV is as expected in QCD. On top of the parton model prediction of $x_T$ scaling:

$$p_T^4 \ E \ \frac{d\sigma}{d^3p} = f(x_T), \ x_T = \frac{2p_T}{\sqrt{s}}$$

one should observe an additional (weak) energy dependence of $d\sigma/d^3p$, due to the $Q^2$-dependence of the QCD amplitudes and of the proton structure functions. The systematic errors do not allow to reach this goal. The ultimate determination of what the energy of a jet is, depends in UA1 on a sizable correction for spill-in and spill-out of energy from the jet fiducial cone. This correction is based on $I_s$ which cannot really be trusted in this role, and in average amounts to increase the jet energy by $\sim 20\%$. The uncertainty in calorimeter calibration is another important source of error. Altogether, the jet energy scale is uncertain by $\pm 70\%$. As a consequence, one can only conclude that the observed cross-section increase between $\sqrt{s} = 546$ and $\sqrt{s} = 630$ GeV is consistent with $x_T$ scaling. The UA1 data is shown in fig.6. One sees that even the $\sqrt{s} = 2$ TeV cross-section is not expected to violate $x_T$ scaling very much (lower broken curve).

UA2 is able to control its absolute energy scale better, but still uncertainties brought in by corrections based on Monte Carlo calculations, uncertainties in calorimeter calibration and collider luminosity etc. produce a $p_T$ independent scale error of $\pm 45\%$. A K-factor of 1.5 is applied to the leading order QCD cross-section, which is also estimated to be uncertain by $\pm 50\%$. In conclusion, $x_T$ scaling is again enough to explain the energy dependence of the inclusive cross-section between the two Collider energies, as seen in fig. 7a. On the other hand, one sees in the same figure that the R807 ISR data do not scale by a large amount. This effect (softening of the spectrum with increasing $\sqrt{s}$) is qualitatively as expected by QCD. This can be empirically expressed by a single parameter: if the $p_T$-exponent 4 is replaced by 4.74, $x_T$ scaling from ISR to collider is restored, as shown in fig.7b. Leading order QCD with $\Lambda_{QCD} = 200$ MeV can fit the UA2 jet-jet mass distribution very well, as shown in fig.8. At $\sqrt{s} = 630$ GeV, a fit was also made to see to what extent the large $M$ tail can accommodate a contribution of a new point-like amplitude due to direct interaction between parton substructures. Following Eichten et al.$^{[4]}$, this was represented as a four-particle Fermi interaction at a scale determined by a cut-off parameter $\Lambda_c$. A finite $\Lambda_c$ value would eventually produce an excess of events at large $p_T$ and $M$. With 90% confidence limit, the result is $\Lambda_c \geq 350$ GeV. UA1 performed a similar analysis and obtained the limit $\Lambda_c > 400$ GeV.
1.3. \( W^\pm, Z^0 \) Production.

The energy dependence of the \( W \)-production cross-section as well as \( d\sigma/dx \), \( d\sigma/dp_T \) as measured by UA1 \(^5\) are found to be consistent with leading order QCD predictions based on Drell-Yan amplitudes and \( Q^2 \)-dependent structure functions. The results are in qualitative agreement with QCD (the 1985 data are still preliminary). For example, the QCD prediction for

\[
\alpha = \frac{\sigma_w (\sqrt{s} = 630 \text{ GeV})}{\sigma_w (\sqrt{s} = 546 \text{ GeV})} = 1.26
\]

should be compared with the experimental results

\[
\alpha_e = \frac{\sigma_w B_e (630)}{\sigma_w B_e (546)} = 1.14 \pm 0.18
\]

\[
\alpha_\mu = \frac{\sigma_w B_\mu (630)}{\sigma_w B_\mu (546)} = 1.00 \pm 0.28
\]

The \( x_T \)-distribution of the \( W \)'s is compared with theory in fig. 9, and is found to be a little harder. From this distribution one can derive the primary parton structure functions \( u(x) \), \( d(x) \) (by analysing \( W^+ \)and \( W^- \)data separately), which are shown in fig. 10. One would find in this data an indication for \( d(x) \) being a little harder than \( u(x) \), which of course cannot be taken at face value. Clearly, this interesting approach is still being tuned, and the data are preliminary.

A straight-forward approach to check QCD is to study associated jet-production in \( W, Z^0 \) events. In a first approach one would simply expect that the relative rate \( N[ W + j ] / N[ W ] \sim \alpha_s \). Indeed, the event rate depends on the number of associated jets as shown in fig. 11. The curve drawn through the data indicates an approximately constant "duty factor".

\[
\frac{N[ W + n_j ]}{N[ W + (n-1)j ]} = 0.33
\]

which can be taken as an indication of the magnitude of \( \alpha_s \). Also, the \( W \)-transverse momentum is larger when jets are produced, and the \( p_T \)-distribution checks very well with QCD. This is shown in fig. 12 (UA2 data), where the shaded histogram shows the events with associated jets, and the full curve is the overall QCD prediction.

1.4. Determination of \( \alpha_s \).

The dream of hadron fans is to derive a real measurement of \( \alpha_s \) by studying some features of jet production. The ideal gymnasium for this is three jet events. Nowadays, these events stand out in the data beautifully more or less as two jet events. When one studies in three-jet events the angular distribution of the hardest jet relative to the beams, one finds a Rutherford scattering-like forward peak as expected in QCD, which is exactly the same as for the two-jet events. The UA2 results are shown in fig. 13 \(^3\).
great majority of cases one of the two weaker jets should be a radiated gluon. A hard effort to derive \( \alpha_s \) by an accurate comparison of data with a QCD Montecarlo was done by UA2 6). However, in order to do this one needs to put in the model

a) \( Q^2 \) - dependent structure functions. UA2 took a single " effective gluon "distribution \( G(x) \) which was derived by the observed two-jet data.

b) Matrix elements. The 2j and 3j amplitudes were computed to first order QCD. A simplifying approximation was done to compute the 4j amplitudes, by taking the 3j ones as computed to 1st order and allowing for an additional gluon radiation.

c) Fragmentation functions, which cannot be computed in perturbative QCD.

d) A detector model.

e) The choice must be made on which parameter should correspond to \( Q^2 \). UA2 took the maximum \( p_T \) among the three jets, \( Q^2 = (p_T^{\text{max}})^2 \).

One assumes that the three jet sample contains also four jets events where one is missed, and similarly that the two jet sample contains also three jet events. The simulated ratio in the model is thus

\[
\frac{R_{QCD}}{R} = \frac{\sigma_{3 \to 2} + \sigma_{4 \to 2}}{\sigma_{2 \to 2} + \sigma_{3 \to 2}}
\]

which is found to be a linear function of the coupling constant, say \( \alpha_s^{\text{1st}} \). To account for not-computed heigher order QCD effects one write

\[
\alpha_s^{\text{1st}} = \alpha_s \frac{R_3}{R_2}
\]

when \( R_2 \), \( R_3 \) are higher order corrections to the 2-jet and 3-jet cross-sections. On comparing with the experimental 3j/2j ratio one finally finds

\[
\alpha_s \frac{R_3}{R_2} = 0.23 \pm 0.01 \text{ (statistical)} \pm 0.04 \text{ (systematic)}
\]

The \( Q^2 \) - dependence of \( \alpha_s \) can also be studied, as shown in fig.14 where the \( j j j \) mass is taken as a varialoble. One finds that data are consistent with QCD expectations, but are unable to prove the existence of a \( Q^2 \) - dependence. This was an excellent work of analysis, ending unfortunately with the prove that the task is nearly impossible. Among the various assumptions and simplifications of rather unknown consequences that had to be made, it is worth indicating a very significant one. In three jet events, the choice of which parameter should be taken to fix the \( Q^2 \) - scale is not obvious. UA2 took \( (p_T^{\text{max}})^2 \). If one had taken the average \( p_T \) of the three jets, \( Q^2 = <p_T^2> \), \( \alpha_s \) would have decreased by 25%!

UA1 carried through a simplified analysis 7) along similar lines, getting the " golden rule" applicable to their jet definition and cuts

\[
\alpha_s \frac{R_2}{R_2} = \frac{1}{0.81} \frac{n_{3j}}{n_{2j}}
\]

The analysis of the new 1985 data is not completed yet, but an agreement with the
published result
\( \alpha_s \left( \frac{R_3}{R_2} \right) = 0.16 \pm 0.02 \) (statistical) \( \pm 0.03 \) (systematic) is anticipated. This number looks lower than the UA2 one – but of course is consistent with it within errors.

The statement of \( e^+e^- \) fans is that they can measure \( \alpha_s \). In principle indeed there is a simple approach in \( e^+e^- \), since the \( s \) dependence of \( R = \frac{\sigma_{h}}{\sigma_{\mu\mu}} \) determines \( \alpha_s \). To second order in \( \alpha_s \) one finds

\[
R_{\text{QCD}} = 1 + \frac{\alpha_s}{\pi} + \left( 1.985 - 0.11_0^{+0.05}_{-0.04} \right) \left( \frac{\alpha_s}{\pi} \right)^2.
\]

Since the series converges fast, this result should already be good to \( \sim 1\% \). The problem is that this prediction must be compared with data whose absolute precision is uncertain to a few percents, which is nearly the same size as the overall QCD effect, and this has an enormous impact on \( \alpha_s \). The curved marked QPM+QCD in fig.15 differs from the QPM line by an essentially \( s \)-independent step of \( \sim 5\% \). This is why one can only measure \( \alpha_s \) to \( \pm 40\% \) from this data (\( \alpha_s = 0.16 \pm 0.06 \)). This fact was made quite clear by Min Chen, who also described the hard work done by Mark J to derive \( \alpha_s \) from an accurate study of three jet events. In this study the usual thrust or \( p_t^2 \) distributions were not used, since they were found to be somewhat more sensitive to the fragmentation models used in the Monte Carlo calculations. Rather, they exploited internal correlations in energy flow. These were compared with full second order QCD predictions. The result is \( \alpha_s = 0.12 \pm 0.01 \). It is remarked by the authors that a first order QCD calculation would have given \( \alpha_s \sim 0.19 \). Certainly the correction made by passing from first to second order was essential; however one can hardly refrain from wondering at this point what \( \alpha_s \) would turn out to be, if somebody would compute the QCD prediction to full third order! Fig.16 shows the Mark J results compared with other \( e^+e^- \) results at lower energy. Two Monte Carlo models yield somewhat different \( \alpha_s \) values. All in all, the order of magnitude of \( \alpha_s \) is established, but not its \( Q^2 \) dependence.

Another approach to derive \( \alpha_s \) from \( e^+e^- \) data is to study the radiative decay of \( Y_{1S} \).

The rate

\[
\frac{Y_{1S} \rightarrow \gamma\gamma e^+e^-}{Y_{1S} \rightarrow ee e^-}
\]

is calculable in QCD and is related to \( \alpha/\alpha_s \). However, even here there are serious obstacles to a clean measurement of this kind. When looking for a direct photon, the background from \( \pi^0, \eta \rightarrow \gamma\gamma \) is unbearable at \( x_\gamma \leq 0.4 \). This reduces heavily the attainable statistical accuracy. A physical background to \( Y_{1S} \rightarrow gg \) decay is \( Y_{1S} \rightarrow qq \), as well as non resonant \( e^+e^- \rightarrow qq \) at the \( Y_{1S} \) mass (where the ratio signal / non-resonant background is \( \sim 1/1 \)). All in all, the result is only qualitative: \( 0.27 \leq \alpha_s \leq 0.40 \).

We can conclude this discussion on \( \alpha_s \) determination by quoting a statement made by L. Rosenberg who discussed this problem in his contribution at the Conference: data are
consistent, model calculations of $\alpha_s$...are not!

1.5. Dimuons in UA1

UA1 presents a nice $\mu^+\mu^-$ mass-spectrum, with data extending up to $m_{\mu\mu} = 35$ GeV\(^{10}\) (fig.17). The data is consistent with Y and Drell - Yan production. However, the background, which can be estimated by the observed rate of $\mu^+\mu^-$ and $\mu^+\mu^-$ pairs and is tentatively attributed to $\bar{b}b$ semileptonic decay, is also of some importance. Despite these clouds, this is the first lepton pair data available at very small $\tau = m^2/s$, where production occurs predominantly by interactions of soft gluons and the cross-section estimates are highly uncertain. The new data is shown in figs. 18a and 18b, and falls nicely on previous estimates of the cross-section at small $\tau$. UA1 also finds an anomously high rate of equal sign dimuons. This of course is interesting, given that $\bar{b}b$ mixing could produce an effect of this type. However, for the time being, given the numerous possible sources of these pairs and the uncertainty in the estimate of their relative weights, one cannot derive any safe result from this observation.

1.6. QCD versus experiment, more.

As it is well known, in QCD one expects gluon jets to be different from quark jets. Specifically, one expects the gluon fragmentation spectrum to be softer and the multiplicity to be substantially higher. Because of the $9/4$ ratio of the color factors, even after accounting for higher order QCD corrections one expects $<n_g> \sim 2 <n_q>$ when $E_j \to \infty$. Both Mark II \(^{11}\) and HRS \(^{12}\) have analysed a special sample of $e^+e^- \to h$ events at $\sqrt{s} = 29$ GeV, using "mercedes-like" events with three very similar jets at relative angles $\theta_{12} \sim \theta_{23} \sim \theta_{31} \sim 120^\circ$ and $E_j \sim \sqrt{s}/3 = 9.6$ GeV. The idea is that kinematically the three jets should be identical, except for the gluonic nature of one of them which should thus be easier to discover. Mark II finds the quark jet fragmentation function at $f_S = 2 \cdot 9.6 = 19.2$ GeV by interpolating lower energy data and defines a distribution ratio as

$$r(x) = \frac{\frac{1}{3\sigma_T} \frac{d\sigma}{dx} (\text{mercedes events at } \sqrt{s} = 29 \text{ GeV})}{\frac{1}{2\sigma_T} \frac{d\sigma}{dx} (\text{two-jet events at } \sqrt{s} = 19.2 \text{ GeV})}$$

The multiplicity excess in the numerator, where the gluon jet is supposed to contribute, is very small and consistent with $<n_g> \sim 1.2 <n_q>$. However, the gluon jet appears to be significantly softer than the quark jet since $r(x) > 1$ at small $x$ and $r(x) < 1$ at large $x$. This effect is illustrated in fig.19.

A similar analysis lead HRS to conclude that gluon and quark multiplicities are essentially the same:
\[ \frac{n_g}{n_q} = 1.29 \pm 0.21 \text{ (stat.)} \pm 0.20 \text{ (syst.)} - 0.41 \]

The result is presented in a provocative picture in fig.20. The data points represent the experimental gluon multiplicity distribution. The narrow peak on the left shows how data would be if gluon jets were identical to quark jets, the wide peak on the right shows the distribution if \( \frac{n_g}{n_q} = 9/4 \frac{n_q}{n_q} \). If there is an excess in gluon multiplicity, this is very small indeed.

A pretty result which checks with QCD prediction is obtained in an unusual experiment of \( \pi^0 \) production at threshold in \( \pi^- \) scattering on atomic electrons at 300 GeV/c \(^{13}\). The reaction is \( \pi^- e^- \rightarrow \pi^0 \pi^- e^- \) can occur at threshold by photon exchange, and the 3-pion vertex does not vanish because of the contribution of the quark-box diagram

\[
\begin{align*}
\text{which weights according to the number of quark colours. The NA7 result is presented in fig. 21, where one shows on the left how well one can detect forward } \pi^0 \text{s, while the cross-section plotted on the right checks well with the PCAC prediction with three colours.}
\end{align*}
\]

2. SOFT HADRON PHYSICS.

2.1. Total \( \bar{p}p \) cross-section.

Since an year or so, an accurate measurement by UA4 of \( \sigma_t \) ( \( \bar{p}p \)) at the CERN Collider is available. This result \(^{14}\) is shown in fig. 22, in comparison with \( pp \) and \( \bar{p}p \) data at lower energy, together with the dispersion relation fit to all real part and total cross-section data available at that time by Amaldi et al. \(^{15}\). One sees that the fit passes exactly through the UA4 point. This fit includes, in addition to effective Regge-pole terms which are important below the ISR, a dominant term at high energies \( \sigma_{\infty} \approx 0.5 \ln^2 s \text{ (GeV)} \text{ mb. Does this term represent the asymptotic behaviour? M.Block and R. Cahn}^{16} \text{ try a fit to the \( \bar{p}p \) total cross-section and real part data as well as to the } pp \text{ data, after including a C-odd term, from } \sqrt{s} = 6 \text{ GeV to Collider energy, by representing } \sigma_t \text{ as}
\]

\[
\sigma_t \approx \frac{\alpha \ln^2 s}{1 + \beta \ln^2 s} \text{ for } s \rightarrow \infty,
\]

and cannot find an acceptable fit if \( \beta = 0 \), which would correspond to \( \sigma_{\infty} \approx \ln^2 s \). On the other hand, an excellent fit can be obtained all over the above-mentioned energy range with \( \beta \neq 0 \) which would imply \( \sigma_{\infty} \approx \text{const. However, if one forgets about the energy region which is complicated by fast varying Regge Pole contributions and}
\]
considers only $\sigma_T (\bar{p}p)$ and $\rho (\bar{p}p)$ at $\sqrt{s} \geq 19.6 \text{ GeV}$ (the pp intermediate energy data can certainly be accomodated by including a C-odd amplitude), one can get a different answer. One can try a fit to 9 points with three parameters as $^{17)}$

$$\sigma_T = A + \alpha \left[ \ln^2 \left( \frac{\sqrt{s}}{\rho} \right) - \frac{\rho^2}{4} \right] , \quad \rho = \frac{\alpha n \ln^3 \sqrt{s}}{\sigma_T}$$

viz. adopting the same analytical form as in ref. $^{16)$, but with $b = 0$. An acceptable fit is found $\alpha = 0.44 \pm 0.06 \text{ mb}$, with $X^2 = 7.23 (6 \text{ d. o. f.})$. One is forced to conclude that the $\sigma_T$ rise is consistent with an asymptotic $\ln^2 s$ dependence. Have we proven that this is the case? Certainly not, the problem is fully open. It is interesting to discuss how we will possibly get to know more in the near future.

First of all, do we already know more based on the UA5 measurements at $\sqrt{s} = 900 \text{ GeV}$? UA5 knows the ratio of the Collider luminosity between 200 and 900 GeV to $\sim 1\%$ even without a direct L-measurement at these energies, through the known dependence of L on machine $\beta$-functions and energy. However, in order to derive $\sigma_T$ at $\sqrt{s} = 900 \text{ GeV}$ one needs a) to know $\sigma_T$ at $\sqrt{s} = 200 \text{ GeV}$, and b) to measure the total interaction rate at $\sqrt{s} = 900 \text{ GeV}$. UA5 derives $\sigma_T$ at $\sqrt{s} = 200 \text{ GeV}$ by interpolating existing data, from the ISR to the SPS measurement of UA4 at $\sqrt{s} = 546 \text{ GeV}$. This should provide a number with a few percent uncertainty (the UA4 error is $2\%$). The problem at $\sqrt{s} = 900 \text{ GeV}$ is much more serious. The total elastic cross-section is derived by assuming a smooth energy dependence of $\sigma_{el}/\sigma_T$ above $\sqrt{s} = 546 \text{ GeV}$. The inelastic cross-section is derived from the observed rate divided by the luminosity. However, since single and double diffractive events are in general (except for large masses) lost by the UA5 trigger, the total rate of these events must be estimated by a Montecarlo calculation, an essential input to which is the total diffractive cross-section. Only non-diffractive inelastic events are well accepted by the trigger. After all this, $\sigma_T (900)$ is found to be small relative to the $ln^2 s$ prediction. It is clear that this argument may be sufficient to raise more curiosity about what $\sigma_T$ really is, but not to produce a reliable number for it.

A measurement of the real part of the forward elastic scattering amplitude was performed at $\sqrt{s} = 546 \text{ GeV}$ by UA4 in 1985 and the results will be available soon. Since $\rho = d \ln \sigma_T / d \ln s$, one will thus measure the local speed of growth of the cross-section. The expectations of refs. $^{15)$ and $^{16)$ are shown in fig.23 and fig. 24. One sees that this $\rho$-measurement will be quite important.

Making the CERN interference experiment possible was not an easy task. A complete interference measurement, where Coulomb production may provide an absolute normalization to the cross-section, should reach $t_{\text{min}} \sim 10^{-3} (\text{GeV/c})^2$. This should be contrasted with $t_{\text{min}} \sim 3 \cdot 10^{-2} (\text{GeV/c})^2$, reached until now by UA4 in the special high $\beta$ ($\beta \sim 100\text{ m}$) runs. At CERN, a special SPS working point was worked out with a super high $\beta$ scheme, reaching $\beta \sim 1000\text{ m}$. The detector location was also optimized (by moving
the bunch-bunch crossing point), such that finally one was able to reach \( t_{\text{min}} \sim 10^{-3} \) (GeV/c)². In this condition a \( \rho \)-measurement is thus possible, although with little safety margin.

If one can rely on an external luminosity monitor to give an absolute cross-section scale, or if one already knows \( \sigma_T \), one can derive \( \rho \) from measurements at slightly larger \( t \). The critical \( t \)-value is where the Coulomb and nuclear amplitudes become equal, \( t_c = \frac{8 \pi \alpha}{\sigma_T} \), which for \( \sigma_T = 62 \) mb is \( t_c \sim 3 \cdot 10^{-3} \) (GeV/c)². If \( t_{\text{min}} < t_c \), one would have to derive \( \rho \) from the excess (or defect) of cross-section over the nuclear contribution. At still higher energies than the CERN Collider, a direct measurement of \( \sigma_T \) will be done up to \( \sqrt{s} = 2 \) TeV at Fermilab by R710 and CDF. Again a glimpse of what happens at still higher energies would be obtained if \( \rho \) could be measured at this energy. However, consider what the problem is for CDF. Recent tests to determine at which minimum distance from the beam high precision silicon detectors could be located inside the beam pipe concluded that \( d_{\text{min}} \geq 7 \sigma_{\text{beam}} \), where \( \sigma_{\text{beam}} \) is the local betatron width of the beam. Assuming this rule to be valid anywhere along the machine, one finds with the presently envisaged high \( \beta \) optics (\( \beta \approx 60 \) m), \( t_{\text{min}} \sim 2 \cdot 10^{-2} \) (GeV/c)². To measure \( \rho \), one would have to reach \( t_{\text{min}} = 10^{-3} \) (GeV/c)², or at least \( t_{\text{min}} = t_c = 2.5 \times 10^{-3} \) (GeV/c)² (for \( \sigma_T = 75 \) mb). One needs to implement also at Fermilab a high or super-high \( \beta \) scheme. At present, the possibility of operating with \( \beta \approx 300 \) m is being considered. At \( \sqrt{s} = 2 \) TeV this would give \( t_{\text{min}} \sim 4 \cdot 10^{-3} \) (GeV/c)², which is still not enough. Of course by reducing the beam energy one would do better. At \( \sqrt{s} = 1 \) TeV, \( t_{\text{min}} = 3 \cdot 10^{-3} \) (GeV/c)² and the measurement would be barely possible.

2.2. Multiplicity Distributions.

Thanks to the UA5 results, it is by now well known that the inclusive (non-diffractive) multiplicity distribution develops a long tail at Collider energies which does not fit with KNO scaling. On the other hand, this behaviour was predicted in the Dual Parton Model, which was at hand since some time \(^{18}\). In this model the soft interaction is attributed to the exchange of a few chains of quark-antiquark pairs between constituents of the projectiles. How well DPM can handle these distributions is illustrated in fig. 25 \(^{19}\). It's also very fashionable nowadays to fit these distributions by the Negative Binomial Distribution, whose at first sight involved expression can be predicted in several theoretical models \(^{20}\). An example showing that Negative Binomial can do equally well as DPM is shown in fig. 26 \(^{12}\). The data shown in this figure is from \( e^+ e^- \) annihilation at \( \sqrt{s} = 29 \) GeV, but NBD performs equally well in hadronic production at high energies. I believe that one should not pretend to learn from these data what they cannot give. The moments of multiplicity distributions are integral measurements of particle correlation functions, which in turn are average phenomena sensitive to many physical processes. These data show that strong
correlations exist in many-body events, including long range ones. Models which do not account for such correlations can be easily disproved by the data. On the other hand, since there are many ways to model these correlations and fit equally well the multiplicity distributions, it is hard to see in such type of agreements a strong argument in favour of a particular model.

2.3. Mini-jets?

UA1 reports more evidence for a number of effects which were around since some time, viz. that in inclusive soft events

i) the central multiplicity ($\frac{dN_{ch}}{dy}$) $y=0$ grows with $\sqrt{s}$.

ii) $<p_T>$ of the charged secondaries grows with $n_{ch}$.

iii) the above growth of $<p_T>$ is stronger at larger $\sqrt{s}$.

Fig. 27 shows the rise of the central multiplicity ($|y|<2.5$) as measured in the 1985 SPS (ramping) Collider run from $\sqrt{s} = 200$ to $\sqrt{s} = 900$ GeV. The effect is 25% over this energy range. Fig. 28 shows the dependence of $<p_T>$ on $n_{ch}$ from 200 (lower data) to 900 GeV (upper data).

UA1 reports also the results of a study of energy clustering in soft events, which looks very much like the calorimeter version of the older short range correlation study in rapidity. The observation that with increasing $\sqrt{s}$ $<n_{ch}>$ grows and that with increasing $n_{ch}$, $<p_T>$ grows originated the suspicion that the underlying mechanism could be the onset of small jets. After all, "real jets" should be present in some events, since in the hundreds of thousands of events collected with the inclusive trigger the tail of the $E_t$ distribution extends to $E_t \sim 40$ GeV. The standard UA1 cluster-search algorithm was applied lowering the threshold as low as $E_j > 5$ GeV, and of course low energy jets were found. Can this finding be turned into a cross-section? Not reliably, for the following reasons:

a) Some jet-like clusters, not generated by a single primary parton, are bound to exist in the data because of accidental overlap of a soft but fluctuating background with a single energetic hadron. This effect was removed from the data before deriving the parton cross-section, but it was estimated with Isajet which is not reliable in this application;

b) an alternative way to using Isajet was to assume that all observed growth of $<p_T>$ with multiplicity is due to the mini-jet component. This allows to derive the amount of "jetty events" at each energy, but is an essentially arbitrary assumption;

c) knowledge of jet energy and good energy resolutions are essential to derive a cross-section when the cross-section is varying fast with energy. At $E_h \sim 10$ GeV, in UA1 $\Delta E_h / E_h \approx 30\%$ for single hadrons, which is already bad. Moreover, a jet is composed in general of several particles of lower energies, and below 10 GeV there is in UA1 an anomalous response of the hadron calorimeter which makes a reliable calibration nearly impossible.

d) finally, the electromagnetic calorimeters, (the gondolas) integrate over $2\pi$ in azimuth
and is particularly hard to measure the electromagnetic component of small jets in large multiplicity soft events.

Still, the authors made their best effort, and quote a preliminary cross-section for inclusive jet production down to $E_j \geq 5$ GeV, which is found to join smoothly with the large $E_j$ data $^{21}$.

2. 4. Diffraction.

Fine data on $\bar{p}p \rightarrow \bar{p}X$ at $\sqrt{s} = 546$ GeV were reported by UA4 $^{22}$, for $50 \leq M_x \leq 140$ GeV. An important outcome - although not a new one any more - is the continued $1/M^2$ dependence of the cross-section which extends from $M \sim \text{a few GeV}$ to these large masses. Fig.29 illustrates the $M^2$ dependence at $t \sim 0.55$ $\text{(GeV/c)}^2$. The charged multiplicity was also studied in UA4, and was found to depend on $M_x$ as in inelastic non-diffractive $pp$ collision at $\sqrt{s} = M_x$, and also to be distributed in rapidity very much in the same way.

Fig.30 shows $\langle n_{ch} \rangle$ versus $M_x$, and fig.31 shows the pseudorapidity distribution at a number of representative masses. The data in this figure are compared with successful predictions of the Dual Parton Model.

It is important to realise that the dependence of the single diffraction cross-section as $d\sigma / dt \, dM^2 \sim (1 / M^2) \, f(t)$ implies that the cross-section for production of masses of hundreds of GeV could be of the order of millibarns at the higher energy Colliders. For example the small angle antiproton spectrometer of CDF would accept $\sim 0.5$ mb of proton diffractive excitation in the range $300 \leq M_x \leq 600$ GeV, at $\sqrt{s} = 2$ TeV. The products of the disintegration of these large masses are an essentially new field of investigation. The similarities with inclusive pp inelastic interactions noted above might justify the feeling that also rare events will still be rare, for instance large $p_T$ jets or heavy flavour production. On the other hand, one might also try to correlate the special dynamics underlying diffraction (Pomeron exchange) with a number of observables in order to learn more about it. For example, one might speculate that large $p_T$ jets will teach us about the parton structure of the Pomeron $^{23}$. Heavy flavours might also be very instructive.

There are some indications of an abundant production of heavy flavours at small angles at the ISR. If this was theoretically understood, for instance as a diffractive contribution, one would be in the condition of making model predictions at the higher energy Colliders. Hoping to account for such a possible anomalously large diffractive production of heavy flavours, K. Ellis has extended the QCD computations of $Q\bar{Q}$ production to order $\alpha_s^3$ $^{24}$. Probably to his disappointment, the $\alpha_s^3$ correction turned out to be distributed in rapidity as $\sigma(\alpha_s^2)$, and moreover to be slightly negative everywhere. Thus, no hope to understand an anomalously large forward production by including these additional terms. How much of the overall production of heavy flavours can be attributed to diffraction in QCD was estimated by E. Soper $^{25}$, by extending the
ideas of ref.23) to include a specific guess for the gluon structure function of the Pomeron. The result (fig. 32) indicates that one third of the overall production could be included in diffraction. Undoubtedly exploring this new world will be of great interest. CDF is in excellent position for doing this 26).

2.5. Heavy Flavours.

It is fair to mention this topics under the heading of soft collisions, because the available information is essentially limited to inclusive charm production. In his rapporteur review 27) S. Reuckroft reported that all hadroproduction cross-sections went recently down by factors of two because branching ratios were better measured and went up by the same amount. While these cross-sections as measured at the ISR cannot be relied upon very much because of the huge factors which had to be applied to go from the observed partial cross-section to the total one, and although beauty searches were so deceiving until now, one can base some hope on experiments of higher quality that are getting now on the floor. I would like to quote Fermilab experiment E 691 28), although the amount of data that was analysed until now is not significant. Fig. 33a shows how the \( D^0 \rightarrow K^- \pi^+ \) peak can be isolated with no background in the \( D^* \rightarrow D^0 \) decay sample. Even better, fig. 33b shows how clearly the \( D^* \) signal can be isolated in the inclusive \( K^- \pi^+\pi^+ \) sample. Such a good mass resolution will be of great help in the beauty search, as well as the excellent spatial resolution achieved by the silicon detectors at the vertex (15 µm in the transverse plane). This excellent space resolution is illustrated in fig. 34, which shows how \( D^* \) and \( D^+ \) decay vertices can be separated in the transverse plane.

2.6. Spectroscopy.

An impressive amount of accurate, hard searches for new heavy multiparticle states are going on which are proving rather stingy at the moment. Still, this is the type of work where fundamental results happen to be reached after hard periods of systematic filing-up of accurate data. To be productive in this field one needs to be first class and unbreakable - theorists might compete only with examples like J. Mandula 29). I shall quote only one example. Fig. 35 shows the inclusive photon spectrum as measured with the new BGO detector of CUSB 30). The tiny structures observed are significant, well below \( E\gamma \sim 100 \) MeV. Fig. 36 shows the two lines observed in the double radiative decay \( Y'' \rightarrow Y' \) through a \( \chi \)-state (\( Y' \rightarrow \mu^+\mu^- \)). The lines are at \( \sim 100 \) and \( \sim 80 \) MeV, and are beautifully separated. It is fair to base hopes on an experiment of such a quality.

CONCLUSIONS.
No discovery was announced this time, but we should not be disappointed: we must agree that discoveries are the exception, not the rule. In particular, we should not cheat ourselves and believe that we have discovered
i) that $\sigma_T \to \text{const as } s \to \infty$

ii) that there is a strong $b\bar{b}$ mixing

iii) that mini-jets are jets

iii) that $\alpha_s$ is $\alpha_s(Q^2)$, and less than ever

iii) that Montecarlo's can predict fragmentation in QCD!

All of us - we must admit - were at best when we were 20. This is, maybe, why Moriond was so great this year. However, I feel that in this case the rule could be upset and Moriond might become even better in the future, THANKS TO TRAN, who is working all the time so hard for us.

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27) S. Reuckroft, New Results on Charm Hadroproduction, Paper presented at this Conference.

28) K. Sliwa, A first look at new Data from tagged Photon Spectrometer charm Photoproduction, Paper presented at this Conference.
Angular distribution (relative to the trigger) of hadrons in large $p_T\pi^0$- events (crosses), and single photon events (dots) (ref. 1).

$\gamma/\pi^0$ ratio in pp versus pp, as measured by experiment R308 at the ISR, NA 24 $\pi^-p$ and $\pi^0p$ data, presented as $\sigma(\pi^-\to\gamma)/\sigma(\pi^-\to\gamma)$ and WA70 data, presented as $(\gamma/\pi^0)^{\pi^-}/(\gamma/\pi^0)^{\pi^+}$. The curves indicate the QCD expectations.

Computed photon conversion probabilities for single photons and photon pairs in the UA2 central detector (curves), and experimental data for unaccompanied (dots) and accompanied (crosses) showers.

Direct photon (crosses) and $\pi^-\pi^+$ inclusive cross-sections as functions of $l_T$ as measured by UA2 at $\sqrt{s}=1.4$.

Average event sphericity versus total transverse energy (UA2).

Inclusive jet production versus jet-jet mass compared with QCD predictions ($A_{QCD}=200$ MeV). The curve labelled with various values of $A_c$ corresponds to the expectations if a contribution by a direct interaction of parton substructures with this cut-off parameter is included.

$W$ distribution in $x_T$, as measured in UA1.

Structure functions of up (a) and down (b) quarks, as derived from observed $w^+$ and $w^-$ distributions in UA1.

Dependence of the number of $W$- events observed in UA1 on the number of associated jets.

$W$ distribution in $p_T$ observed in UA1, compared with QCD prediction. The events with associated jets are also separately indicated (shaded histogram).

Angular distribution of two-jet events (dots) and of hardest jet in three-jet events (crosses), as measured in UA2.

$\alpha_s$ as determined in UA2 in the study of three jet events, versus three-jet mass.

Compilation of $R=\sigma_{\mu\mu}/\sigma_{\mu\mu}$ data in $e^+e^-$, compared with Quark Parton Model, Quark Parton Model + QCD, Quark Parton Model + QCD + Weak contributions.

$\alpha_s$ results from the Mark J study of three jet events (Mark J data are above $\sqrt{s}=20$ GeV). The full curves are QCD predictions for $A_{QCD}=160$ and $55$ MeV, to be compared with data derived respectively with the Lund and Ali Montecarlos.

Dimuon mass spectrum from UA1.

$Y$ production cross-section (times $B_{\mu\mu}$) at the SPS Collider, compared with previous data at much larger values of $l_T = m_N/\sqrt{s}$; the same, for Drell-Yan.

Fragmentation functions ratio, as explained in the text, of mercedes events to two-jet events in $e^+e^-\to$ hadrons at $\sqrt{s} \leq 29$ GeV.

Charged multiplicity distribution of the "gluon jet" in mercedes events, at $\sqrt{s}=29$ GeV, as observed by ISR.

Reconstructed forward $\pi^0$ peak in NA 7 (left), and measured $\pi^-\pi^+	o\pi^0\pi^-\pi^+$ cross-section in comparison with the prediction of PCAC with three colours.

The Amaldi et al. fit to the pp and pp total cross-section data, up to ISR energies. The UA4 measurements at $\sqrt{s}=1.4$ GeV is also shown.

Fit to the real parts of ref.15, predicting $\rho<0.15$ at Collider energies.

Ref.16) fits to $\rho(pp)$ and $\rho(pp)$ under the assumptions of $\sigma_{\mu\mu}=ln^2s$ and $\sigma_{\mu\mu}=const.$

UA5 multiplicity data within pseudorapidity bins of variable width, fitted by DPM distributions.

Multiplicity distribution $e^+e^-\to$ hadrons within rapidity bins of variable width, fitted by Negative Binomial Distributions. $\Psi = (n_{ch})^{-1}dN/dn_{ch}$

Rise of charged multiplicity at $|y|<2.5$ between 200 and 900 GeV, as measured by UA1.

Rise of the average $p_T$ of charged secondaries as a function of $n_{ch}$, at $\sqrt{s}=200$ (lower data) and...
\( \sqrt{s} = 900 \) (upper data). The continuous curves are drawn to guide the eye through the data.

Fig. 29 1/\( M_{2} \) dependence of single diffraction, from the ISR to the SPS Collider\(^{22}\).

Fig. 30 Dependence of \( <n_{ch}> \) on \( M_{x} \) in single diffraction as measured by UA4. Inelastic non-diffractive pp data versus \( \sqrt{s} \) are also displayed.

Fig. 31 Pseudorapidity distributions of charged prongs in proton diffractive excitation. UA4 data\(^{22}\).

Fig. 32 Rapidity distribution of beauty production in proton diffraction at the SPS Collider, as estimated by E. Soper\(^{25}\).

Fig. 33 a) \( D^{0} \rightarrow K^{-}\pi^{+} \) mass peak after imposing the constraint of the \( D^{*} - D^{0} \) mass difference, and b) \( D^{+} \) signal in the inclusive \( K^{-}\pi^{+}\pi^{+} \) mass spectrum, from a first sample of analysed data of E 691\(^{25}\).

Fig. 34 Vertex structure on transverse event plane in a \( D^{*} \rightarrow D^{0}\pi^{+} \) event, as observed in E 691. The \( D^{*} \) is produced on the lower vertex where the \( \pi^{+} \) also originates, while the \( D^{0} \rightarrow K^{-}\pi^{+} \) secondary vertex is on the upper left.

Fig. 35 Inclusive photon spectrum from a CUSB run at the \( Y \) mass.

Fig. 36 The two photon lines in an sample of \( Y \rightarrow \gamma X \rightarrow \gamma \gamma Y' \rightarrow \gamma \gamma \mu^{+}\mu^{-} \) decays.
Fig. 1

NA 24

$E_T^{\text{track}} > 0.5\text{GeV}$

$p_T^{\text{trigger}} > 4\text{GeV/c}$

$1/N_{\text{event}} \frac{dN}{d\phi}$

\[\text{[radian}^{-1}]\]

\[\pi/2 \quad \Delta \phi \quad \pi\]

$\gamma$

$\pi^0$
Fig. 5

Fig. 6
\[ f(x_T) = A(1-x_T)^m / x_T^n \]

Best fit for

\[ m = 6.54 \pm 0.15 \]
\[ n = 4.74 \pm 0.06 \]

\[ X_T = \frac{2P_T}{\sqrt{s}} \]

Fig. 7
UA2 $\bar{p}p \rightarrow \text{jet+jet+x}$
($|\ln_{\text{jet}}| < 0.85$)

- $\sqrt{s} = 630$ GeV (1984)
- $\sqrt{s} = 546$ GeV (1983)

QCD
- $\sqrt{s} = 630$ GeV
- $\sqrt{s} = 546$ GeV

$\Lambda_c = 300$ GeV

Fig. 8

$\frac{1}{N} \frac{dN}{dx_w}$

- $\sqrt{s} = 546$ [GeV]
- $\sqrt{s} = 630$ [GeV]

Fig. 9
$\sqrt{s} = 630$ GeV

UA1

Eichten et al.

(a)

$0.2 \ 0.4 \ 0.6 \ 0.8$

$x_u$

Fig. 10

(b)

$0.2 \ 0.4 \ 0.6 \ 0.8$

$x_u$

Fig. 11
Fig. 14

merged values of all experiments from PETRA & PEP

\[ \frac{\alpha_s}{K_2} \]

\( M \text{ [GeV]} \)

1984

1985, preliminary

Fig. 15

merged values of all experiments from PETRA & PEP

\( e^+ e^- \rightarrow \text{hadrons} \)

Best fit with
\( \Lambda = 342 \text{ MeV} \)
\( M_{Z^0} = 93.0 \text{ GeV} \)
\( \sin^2 \theta_W = 0.22 \)

QPM

QPM + weak

QPM + QCD

\( \sqrt{s} \text{ [GeV]} \)
Strong Coupling Strength

\[ \alpha_s \]

\[ \sqrt{s} \text{ [GeV]} \]

Fig. 16

Isolated Unlike Sign Dimuons - '83-'84-'85

Drell-Yan (42.5±9.6 events)

D.Y.+b\bar{b} (9.8±7.8 events)

D.Y.+b\bar{b}+Y (36.8±6.3 events)

UA1

Dimuon mass [GeV/c²]

Fig. 17
Fig. 18
Fig. 19

Fig. 20
Fig. 25

\[ P_n \sim \frac{n}{\langle n \rangle} \]

- \(|h| < 0.5\) (x10^{-6})
- \(|h| < 1.5\) (x10^{-4})
- \(|h| < 3\) (x10^{-7})
- \(|h| < 5\)

\[ z = n / \langle n \rangle \]

Fig. 26

\[ \gamma(z) \]

- \(|\gamma| < 0.1\)
- \(|\gamma| < 0.25\)
- \(|\gamma| < 0.5\)
- \(|\gamma| < 1.0\)
- \(|\gamma| < 1.5\)
- \(|\gamma| < 2.0\)
- \(|\gamma| < 2.5\)
- \(|\gamma|\) FULL

\[ z = n / \langle n \rangle \]
Fig. 29

\[ \frac{d^3\sigma}{dt\,dM^2} [\text{mb}/\text{GeV}^4] \]

- \( t = 0.55 \text{ GeV}^2 \)

\[ \frac{1}{M^2} \]

Fig. 30

- Inelastic non-diffractive
- UA4 diffractive
Fig. 31

Fig. 32

b Quark Production
M=5.4 GeV
√s=540 GeV
Fig. 33

Fig. 34
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Do the Cross Sections for pp and p\bar{p} Continue to Rise

as log^2(s/s_0)?

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