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ON THE ASSOCIATE PRODUCTION OF
A NEUTRAL INTERMEDIATE - MASS HIGGS BOSON
WITH A SINGLE TOP QUARK AT THE LHC AND SSC

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**On the associate production of
a neutral intermediate-mass Higgs boson
with a single top quark at the LHC and SSC**

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Abstract

The typical behaviour of the semiweak single top quark production mechanism in hadron-hadron collisions in comparison with heavy quark pair production suggests that boson-gluon fusion may lead to important associated Higgs cross-sections as well. This is of particular interest for the interval in which the Higgs mass is close to the mass of the Z^0 and well below $2 M_{Z^0}$. In this region experimentally significant signals are known to be not very exhaustive while complicated to extract due to large backgrounds.

In two recent papers, studies have been presented in which contradicting results have been reported on the magnitude of Higgs production in association with a single top quark. We present an independent analysis and trace the origin of the discrepancies.

We find complete agreement with the results obtained by Stirling and Summers and therefore agree with their conclusion that this new channel is barely of relevance at the next generation of hadron colliders.

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Résumé

Le comportement spécifique de la production semi-faible du quark top suggère que la production associée d'un boson de Higgs pourrait conduire à une section efficace de l'ordre de grandeur du processus QCD. Nous mettons en évidence une compensation entre les deux graphes en jeu. Le nombre d'événements prédit est négligeable en accord avec les résultats de Stirling et Summers.

Keywords : Higgs boson, top quark.

1. Introduction

Although data in high energy particle interactions collected over the past decade show striking agreement with predictions from modern gauge theories, many principles essential to these theories remain as yet not probed. All models describing electroweak interactions through the mechanism of symmetry-breaking contain physical particles of the bosonic type. They predict that the boson sector will consist of at least one or more spinless so-called Higgs boson(s) which through the Higgs mechanism are responsible for mass generation. None of the models however, provide precise predictions on the mass of these scalar particles.

Direct searches for scalar particles at existing experimental facilities have been negative while the most stringent limit on the mass of the Standard Model [1] neutral scalar (H^0) has been derived by the ALEPH Collaboration at the Large Electron Positron (LEP) accelerator at Cern: $M_{H^0} \geq 57 \text{ GeV}/c^2$ [2]. No upper bound for H^0 can be derived from experimental data, however, there are strong theoretical indications that its mass should not exceed $\sim 1 \text{ TeV}/c^2$ [3] (or even: $M_{H^0} < 700 \text{ GeV}/c^2$ [4]). LEP200 will certainly improve the present lower bound, but its sensitivity is most likely to decrease rapidly above $M_{H^0} \sim 80 \text{ GeV}/c^2$.

Unless the Higgs is discovered at LEP200, present experimental facilities will not be sufficient to probe deeper into the symmetry breaking mechanism and new accelerator facilities are therefore indispensable. Although e^+e^- colliders have proven to be relatively clean and precise laboratories for new particle searches, with accelerator technologies as of today one is not able to construct electron machines which will both extend into the (multi) TeV centre-of-mass regime and provide enough luminosity to produce observable event rates at the same time. Hadron colliders such as the proposed Large Hadron Collider (LHC) at Cern and Superconducting Super Collider (SSC) in Texas provide laboratories for studying interactions that are less clean, but after careful analysis, turn out to allow for exploitation of parton interaction centre-of-mass energies into the multi TeV region. Luminosities at these colliders will be sufficiently large to allow for a scan for (neutral) scalars with a mass up to $\sim 1 \text{ TeV}/c^2$ [5, 6].

However, for an important mass window, namely $80 \text{ GeV}/c^2 \leq M_{H^0} \leq 120 \text{ GeV}/c^2$, which covers the lower part of the so-called intermediate Higgs mass interval, $80 \text{ GeV}/c^2 \leq M_{H^0} \leq 2 M_{Z^0}$, it will be extremely difficult to resolve an experimentally significant signal from the background. The rate for Higgs decays into four leptons is much too low in this region to be useful whereas the decay $H^0 \rightarrow \gamma \gamma$ puts severe constraints on photon detection resolution in order to survive the two-photon background from quark-antiquark annihilation and gluon-gluon scattering in the reconstruction of the Higgs mass [7].

Alternative production mechanisms have been studied for which backgrounds are sufficiently small. One should bare in mind that total signal event rates per unit of integrated luminosity are very limited and demand highly efficient detectors. The main two processes considered are [8]:

$$pp \rightarrow q\bar{q}' \rightarrow W^* \rightarrow W H^0 \quad (1)$$

where the W decays leptonically and $H^0 \rightarrow \gamma\gamma$ ($H^0 \rightarrow \tau^+\tau^-$ seems to be very difficult experimentally [9], but is interesting enough to justify further study) and a Higgs produced in association with a heavy quark-antiquark pair:

$$pp \rightarrow gg \rightarrow Q \bar{Q} H^0 \quad (2)$$

For a sufficiently large quark mass ($m_Q > M_W$), the heavy flavour will decay into a on-shell W and a quark with a substantially smaller mass than the decaying quark Q . Leptonic decays of the W will then provide additional tagging possibilities. Signal to background ratios for these mechanisms are of order 3 ~ 4 [10] at both the LHC and SSC operating at design luminosity, while the total number of events (in which decay branching ratios for H^0 , W and heavy quark have been properly included) does not exceed 50 events per year of data taking.

Single heavy flavour production has been discussed in length in the literature [11] and proceeds through the mechanism of boson-gluon fusion. The fact that the emerging event contains only one very heavy flavour reduces the overall hadronic activity with respect to heavy flavour pair production. A similar argument holds if one compares the production of a Higgs in association with a single heavy flavour or with a pair of heavy quarks. Obviously, also here the heavy flavours will generate a considerable amount of hadronic fragments. Diaz-Cruz and Sampayo [12] have therefore studied the process:

$$pp \rightarrow Wg \rightarrow q Q \bar{Q} H^0 \quad (3)$$

where $m_Q \gg m_q$, arguing that this process may lead to event rates comparable with the ones from the process in (2). Using the effective W approximation, they predict of order 30 $W\gamma\gamma$ events per running year at LHC.

This result seems, a priori not unreasonable if one follows the arguments presented in [11]. In comparing the shapes of the total cross-section for single top quarks and heavy flavour pair production as a function of the heavy quark mass, one concludes that for a relatively small top quark mass ($m_t \sim 100 \text{ GeV}/c^2$), pair production clearly dominates. However, increasing the top mass illustrates the different behaviour of the two mechanisms as a function of the invariant mass of the hard scattering process. The matrix element squared for the boson-gluon fusion mechanism varies with $1/\hat{s}$ (due to the t -channel exchange of the W), while the $Q\bar{Q}$ production cross-section is proportional to $1/\hat{s}$. Both mechanisms become equally important for $m_t \sim 250 \text{ GeV}/c^2$. One may argue that for associate Higgs production where $m_t \sim 100 - 200 \text{ GeV}/c^2$ and $M_{H^0} \sim 100 \text{ GeV}/c^2$, process (3) may compete with the heavy flavour pair production mechanism.

Stirling and Summers [13] have explicitly calculated the semiweak process as well and obtain a result in striking contradiction with [11]. They claim that W -gluon fusion leads to a negligible contribution to the heavy flavour associated Higgs production cross-section, while increasing the top quark mass at a fixed value of the Higgs mass decreases the cross-section even further.

Our paper is organised as follows. In the next section we discuss the calculation of the matrix elements and elucidate the large compensation occurring among Feynman graphs. In section 3, we pre-

sent cross-section estimations as a function of the Higgs mass, for different top quark mass values and illustrate the various contributions.

2. Process description

As the origin of the discrepancy between the results of [12] and [13] is not apparent, we have performed an independent calculation of the matrix element squared. Diaz-Cruz and Sampayo have applied the effective W approximation in order to obtain numerical results for the W-gluon scattering process. In ref. [13], the authors have used the helicity formalism extended to heavy quarks to calculate $pp \rightarrow qb \rightarrow q'H^0t$ [14]. This method is well suited to include the matrix element for the decay of the top quark as well [15]. In order to verify the total cross-section for producing a top quark and a Higgs and to understand the disagreement with [12], they have performed the same calculation using the effective W approximation. They find that they cannot reproduce the results of [12]. Since the discrepancies between the two calculations appear at the top quark production level, we have computed the matrix element for the process

$$pp \rightarrow qb \rightarrow q'tH^0 \quad (4)$$

using standard Feynman rules. Here we use that convolution with the bottom quark distribution function in the proton gives excellent agreement [16] with calculating the more complicated graph where a $b\bar{b}$ pair is produced from gluon splitting [11]. The number of Feynman diagrams is then small. They are depicted in figure 1.

Fig. 1: Feynman diagrams contributing to the process $q b \rightarrow q' t H^0$.

In the following we will show that although each of the graphs in fig. 1 give potentially large contributions to the cross-section, a huge interference term is arising which is destructive. Let us consider the sub-process $W(k) b(p_1) \rightarrow H^0(p_H) t(p_2)$ for which the diagrams are given in fig. 2.

Fig. 2: Feynman diagrams contributing to the process $W b \rightarrow t H^0$.

Conventional Feynman rules lead to the following expressions for the matrix elements:

$$M_a = -i (ig_W) (ig_{WWH}) \frac{\bar{u}(p_2) \gamma_\mu (1 - \gamma_5) u(p_1)}{(p_1 - p_2)^2 - M_W^2} \left(g^{\mu\nu} - \frac{(p_1 - p_2)^\mu (p_1 - p_2)^\nu}{M_W^2} \right) \varepsilon_\nu \quad (5)$$

$$M_b = i (ig_W) (ig_{t\bar{t}H}) \frac{\bar{u}(p_2) (\not{p}_1 + \not{k} + m_t) \gamma^\mu (1 - \gamma_5) u(p_1)}{(p_1 + k)^2 - m_t^2} \varepsilon_\mu \quad (6)$$

where the couplings among fermions, Higgs and W are specified through their indices (g_W is taken to be universal to all fermion doublets, thus ignoring off-diagonal Cabibbo-Kobayashi-Maskawa matrix elements). The high energy behaviour of these amplitudes becomes worse when the incoming Ws are longitudinally polarised. Then, for a W with 4-momentum k , the polarisation vector can be written as $\varepsilon_\mu = \frac{k_\mu}{M_W} + O(\frac{M_W}{k_0})$. Substitution in eqs. (5) and (6) leads to:

$$M_a = -i \frac{g_W g_{WWH}}{M_W} \bar{u}(p_2) \left(\frac{(M_W^2 - k \cdot p_H) m_t}{M_W^2 (M_H^2 - 2k \cdot p_H)} - \frac{k}{M_H^2 - 2k \cdot p_H} \right) (1 - \gamma_5) u(p_1) \quad (7)$$

$$M_b = -i \frac{g_W g_{t\bar{t}H}}{M_W} \bar{u}(p_2) \frac{M_W^2 + 2p_1 \cdot k + m_t k}{M_W^2 + 2p_1 \cdot k - m_t^2} (1 - \gamma_5) u(p_1) \quad (8)$$

In the high energy limit, masses may be neglected and the amplitudes reduce to

$$M_a = -i \frac{g_W g_{WWH}}{M_W} \left(\frac{\bar{u}(p_2) m_t (1 - \gamma_5) u(p_1)}{2M_W^2} - \frac{\bar{u}(p_2) k (1 - \gamma_5) u(p_1)}{M_H^2 - 2k \cdot p_H} \right) \quad (9)$$

$$M_b = -i \frac{g_W g_{t\bar{t}H}}{M_W} \left(\bar{u}(p_2) (1 - \gamma_5) u(p_1) + \frac{\bar{u}(p_2) m_t k (1 - \gamma_5) u(p_1)}{2p_1 \cdot k} \right) \quad (10)$$

Substitution of the couplings $g_{t\bar{t}H} = -g \frac{m_t}{2M_W}$ and $g_{WWH} = gM_W$ (with g defined as the SU_2 gauge coupling), finally shows that the magnitude of the leading terms in both amplitudes are equal and cancel satisfying the unitarity constraint [1, 17].

It is interesting to note that this result is also applicable to extensions of the Standard Model such as minimal SUSY, where the strength of the top-Higgs coupling depends on the choice of the mixing angle β (defined by the ratio of the two vacuum expectation values) and angle α (from the diagonalisation of the scalar mass matrix). Unitarity imposes [18]:

$$-\frac{m_t}{2M_W^2} g_{WWH} g_{Wtb} - g_{WHh} g_{Wtb} = g_{t\bar{t}H} g_{Wtb} \quad (11)$$

where h denotes the charged scalar field. We now turn to the computation of the cross-section for the process $pp \rightarrow qb \rightarrow qH^0 t$.

3. Results

After squaring the amplitudes of the graphs in figure 1, we obtain relatively comprehensive expressions for the contributions from both diagrams and their interference:

$$|M|^2 = 32 \left(\frac{N_a}{D_a} + \frac{N_b}{D_b} - \frac{N_{ab}}{D_{ab}} \right) \quad (12)$$

with the numerators and denominators referring to the two diagrams in fig. 1 and their interference given by:

$$N_a = |g_{WWH} g_W|^2 \left(2s_{12}s_{34} + m_t^2 \frac{2s_{12}s_{23} - 2s_{12}s_{34} + s_{13}s_{24} - s_{14}s_{23}}{M_W^2} + \right. \\ \left. m_t^2 \frac{-s_{13}s_{24}m_t^2 + 2s_{12}s_{23}s_{24} - 2s_{12}s_{24}s_{34} + 2s_{13}s_{24}^2 - 2s_{14}s_{23}s_{24} + 2s_{14}s_{24}s_{34}}{2M_W^4} \right) \\ N_b = |g_{tH} g_W|^2 \left(4s_{12}s_{35}s_{45} + 8s_{12}s_{34}m_t^2 + 8s_{12}s_{35}m_t^2 - 2s_{12}s_{34}M_H^2 \right) \quad (13)$$

$$N_{ab} = |g_{WWH} g_{tH} g_W^2| \left(-8s_{12}s_{34}m_t - 4s_{12}s_{35}m_t + \right. \\ \left. m_t^3 \frac{s_{13}s_{25} - 2s_{13}s_{24} - s_{15}s_{23} + 2s_{14}s_{23} - 4s_{12}s_{23} + 2s_{12}s_{34} + s_{12}s_{35}}{M_W^2} + \right. \\ \left. 2m_t \frac{s_{12}s_{25}s_{34} - s_{12}s_{24}s_{35} - s_{12}s_{23}s_{45} - s_{13}s_{24}s_{45} + s_{14}s_{23}s_{45} - s_{14}s_{25}s_{34} + s_{15}s_{24}s_{34}}{M_W^2} \right)$$

$$D_a = \left(\left(-2s_{13} - M_W^2 \right)^2 + M_W^2 \Gamma_W^2 \right) \left(\left(-2s_{24} + m_t^2 - M_W^2 \right)^2 + M_W^2 \Gamma_W^2 \right) \\ D_b = \left(\left(-2s_{13} - M_W^2 \right)^2 + M_W^2 \Gamma_W^2 \right) \left(2s_{45} + M_H^2 \right)^2 \quad (14) \\ D_{ab} = \left(\left(-2s_{13} - M_W^2 \right)^2 + M_W^2 \Gamma_W^2 \right) \left(2s_{45} + M_H^2 \right) \left(-2s_{24} + m_t^2 - M_W^2 \right)$$

where the scalar products of the 4-vectors are denoted by $s_{ij} = p_i \cdot p_j$ ($p_1 \equiv$ the momentum of the incoming light quark, $p_2 \sim$ bottom quark, $p_3 \sim$ outgoing light quark, $p_4 \sim$ t-quark and p_5 is defined as the Higgs 4-momentum). Convolution with structure function parametrisations and phase space integra-

tion is performed numerically using the Eurojet simulation package [19]. The scale at which the structure functions are evaluated is chosen to be $Q^2 = \hat{s}$, the total invariant mass of the partonic sub-process squared. We use the Eichten et al. [20] structure function parametrisations (set 1, $\Lambda = 0.2$ MeV). Since the amplitude squared does not contain any intrinsic divergences, integration can be carried out over full phase space.

In figures 3^{a+b}, we present total cross-sections as a function of the Higgs mass at LHC and SSC energies respectively. We have chosen three different values for the top quark mass in order to illustrate the cross-section dependence (lower set of curves: $m_t = 100$ (solid), 137 (dashed) and 200 (dotted) GeV/c^2 resp.). For comparison purposes, we have included cross-section estimations for the boson fusion mechanisms without associated heavy flavours as well (upper solid curve: WW-fusion, upper dashed curve: ZZ-fusion).

Fig. 3: Higgs production cross-sections at the LHC (a) and SSC colliders (b) ($\sqrt{s} = 16$ GeV and $\sqrt{s} = 40$ TeV resp.) through boson-boson fusion without (top curves) and with (bottom curves) associated top quarks.

To see more clearly at which centre-of-mass energies of the partonic sub-process the various contributions to the cross-section are important, we have plotted for two choices of the top quark mass (fig. 4) the individual contributions of the two graphs in fig. 1. The dotted curves show the behaviour of the diagram containing the WWH^0 coupling as the Higgs mass is increasing (and thus for increasing $\hat{s}$). The cross-section corresponding to the contribution of the second graph is given by the dashed curves. The total cross-section is then obtained by adding the two and the interference term.

Fig. 4: The cross-section for $pp \rightarrow qtH^0$, at the LHC (a, c) and SSC (b, d) as a function of the Higgs mass. Apart from the total cross-section (solid curves), the individual contributions of the amplitudes squared containing the $t\bar{t}H^0$ (dashed) and WWH^0 (dotted) couplings are drawn for two top quark masses, 100 (a, b) and 200 (c, d) GeV/c² resp..

For increasing M_{H^0} , the amplitude containing the $t\bar{t}H^0$ coupling becomes less important due to its $1/\hat{s}$ dependence and with it the interference term becomes smaller (in absolute value). At small M_{H^0} , the $t\bar{t}H^0$ amplitude takes over and a further decrease of the Higgs mass leads to a decreasing contribution from the interference term. In the intermediate region, that is at $m_t \sim M_{H^0}$, destructive interference reaches its peak value. Apart from being larger at the SSC (4^{b+d}), the cross-sections fall less rapidly for increasing Higgs mass than at the LHC (4^{a+c}). Evidently, for a small top quark mass, the cross-section is dominated by the process containing the WWH^0 coupling. Although we have argued that cancellations should occur naturally to preserve unitarity, it is quite striking to see that the magnitude of the interference term rapidly rises as a function of the top mass at a fixed Higgs mass. In figure 5 we show this behaviour for $M_{H^0} = 200 \text{ GeV}/c^2$, varying the top mass between 50 and 450 GeV/c^2 .

Fig. 5: Higgs production cross-section at the LHC with $M_{H^0} = 200 \text{ GeV}/c^2$ as a function of m_t . The increasing importance of the $t\bar{t}H^0$ coupling leads to an increase of the cross-section at larger values of m_t (the absolute value of the interference term is drawn).

Finally, if we compare our results with the ones obtained by Diaz-Cruz and Sampayc, they indeed are quite different. However, the fact that their cross-sections seem to increase for an increasing top quark mass and are even larger than the contribution from the graph with the WWH^0 alone, indicates either small or negligible interference contributions or an interference term with wrong sign*.

* It turns out that we can precisely reproduce their results by changing the sign of the interference term in our expression for the matrix element squared.

Conclusions

We have shown that the disagreement between the results of ref. [12] and [13] is (most certainly) due to the very large destructive interference which occurs between the two amplitudes describing the partonic scattering process. Naively, this compensation is not unexpected since, for sufficiently large top and Higgs mass, the two contributing Feynman diagrams will each give rise to events populating overlapping regions in phase space. Moreover, in order not to violate unitarity, compensation is bound to occur.

We find our results in perfect agreement with Stirling and Summers, while changing the sign of the interference term in the expression for the matrix element squared reproduces the results of ref. [12].

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