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Heavy flavor production at large transverse momentum through the boson-gluon fusion mechanism

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We present alternative calculations of the matrix elements for the process of heavy flavor production through the $\gamma/W/Z$-gluon fusion mechanisms in hard hadronic interactions. The expressions for the matrix elements squared have been implemented in the EUROJET Monte Carlo event generator providing a flexible tool to study event distributions for both partonic and hadronic final states. In case of a charged mediator, we find complete agreement between our $O(\alpha^2\alpha_s)$ calculations and recent calculations by Zerwas *et al.* Results on the factorization of the $O(\alpha^2\alpha_s)$ matrix elements, which allows an independent treatment of $O(\alpha^2)$ and $O(\alpha^2\alpha_s)$ contributions to cross sections, are presented as well. We have performed an analysis of the main heavy flavor sources in hadronic interactions and have studied their relative importance as a function of accelerator energies. Heavy flavor masses have been varied over a wide range of possible values. The pure QCD contributions dominate for only part of the parameter combinations, while the distinct event topology arising in boson-gluon fusion processes enhances the observability of heavy flavors. We briefly discuss implications of the fusion mechanisms for production of heavy E_6 fermions.

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I. INTRODUCTION

In this paper we present studies of semiweak boson-gluon fusion mechanisms in hadron-hadron interactions. Boson-gluon fusion is potentially important for the description of heavy flavor production at future colliders and may even have some relevance for the observation of the top quark at the Fermilab Tevatron. In the case of W-gluon fusion, for large quark masses ($m_Q > 250$ GeV/c^2) the exchange of the W provides cross sections competitive with QCD pair production. In addition, the event topology is quite different from the one resulting from pure QCD pair production of heavy quarks. This is mainly due to the large mass difference between the top and bottom emerging from W-gluon fusion. It has been argued [1] that the semileptonic decay of a single top quark may then considerably ease the precise mass determination of the top quark. In Sec. II, we discuss the structure of the $O(\alpha^2)$ and $O(\alpha^2\alpha_s)$ matrix elements [2], which form the basis of our calculations of cross sections, in some detail. We compare pure QCD and semiweak cross sections for different top-quark masses, at different collider energies.

Another source of heavy flavor pair production besides pure QCD is obtained when the charged weak mediator in the W-gluon fusion process is replaced by either a neutral weak boson or a photon. Although the production of two heavy objects is now reduced due to phase-space suppression and different behavior of the expression for the transition amplitude, the exercise is interesting since the diagrams containing a photon interfere with the ones having a Z^0 propagating. Results on bottom-quark production at present hadron colliders (CERN $Sp\bar{p}S$ and the Tevatron) and contributions to the top-quark signal at future colliders are presented in Sec. III.

New physics, as a consequence of extensions to the minimal standard model, may give rise to the production of exotic fermions. It is straightforward to apply our calculations to extensions of the standard model which give rise to, for instance, fourth-generation quarks with or without heavy neutrinos. Even though the ratio of the masses of the two partners in the same (new) doublet is not *a priori* expected to be as large as is the case for top and bottom quarks [3]. In Sec. IV we describe the application of our formulas for singlet down quarks whose masses are not subject to common standard-model restrictions and which acquire a coupling with weak bosons and ordinary quarks through mixing. Those quarks have been introduced in an attempt to understand the origin of fermion masses and their mass hierarchy using the seesaw mechanism [4] and find a raison d'être in some compactification of superstring theories [5]. We have explicitly calculated cross sections for a wide range of E_6 fermion masses. We conclude that both W-gluon and Z^0-gluon fusion mechanisms may serve as powerful "tools" in the observation and analysis of exotic phenomena as well.

II. TOP QUARKS: QCD VERSUS CHARGED WEAK CURRENTS

As the mass difference between the two members of the yet incomplete third generation of quarks increases, the production of top quarks through gluon fusion becomes competitive with QCD pair production and even dominant for very large top-quark masses [6]. Provided the top quark decays as prescribed by the standard model, a lower limit of 89 GeV/c^2 for the top-quark mass is obtained by the CDF Collaboration at the Fermilab $p\bar{p}$ collider [7]. At the CERN e^+e^- collider LEP, precise measurements of the Z^0 pole (that is M_Z and Z^0 decay properties) have initiated studies to determine the top-quark mass from weak (loop) corrections [8]. The interpretation that weak interactions are responsible for the experimentally observed mixing effects in the B-meson sector (both B^0-$\bar{B}^0$ and B_s^0-$\bar{B}_s^0$ mixing) [9] allows for another independent determination of the top-quark mass. The size of these weak effects is directly proportional to the mass of the t quark. However, although the calculation of the weak box diagrams is theoretically based on solid grounds, the latter method contains relatively large uncertainties on how to correct for hadronic effects (bag parameter for instance) since the b quarks appear in bound states.

Different analyses of experimental data obtained by a variety of experiments present standard-model restricted values for the top-quark mass: $m_t=139\pm38$ GeV/c^2 [10], $m_t=124^{+28+20}_{-34-15}$ GeV/c^2 [11], and $m_t=142^{+26}_{-49}$ GeV/c^2 [12]. These values are not in disagreement with a recent summary by Altarelli [13], who concludes $m_t=140\pm45$ GeV/c^2. With increasing LEP statistics and a more precise determination of the W mass [by the Collider Detector at Fermilab (CDF)], in the near future one may expect the error bars to become significantly smaller. Provided the top quark will be discovered, we face two different possibilities. Either the mass of the top lies in the range $\sim$100–200 GeV/c^2, which will be strongly in favor of the standard model and will at most leave room for some minimal extensions, or the top-quark mass may be found to be far outside the expected range, which may indicate completely new physics. (Note, however, that new vector bosons may, after some fine-tuning between top and W', Z' contributions to radiative corrections, remove the present upper bound on the top-quark mass [13].)

In any case, two aspects of the W-gluon fusion process become important in comparison with the usual QCD production of heavy fermions. At first, the distinct event topology arising in the fusion process may provide a more precise determination of the mass of the heavy fermion (top) as advocated in Ref. [1]. Second, if the ratio between cross sections for the weak production of t quarks and the QCD pair production at the lower bound of the mass interval is "only" of order $\frac{1}{10}$, both processes become competitive for increasing fermion masses ($m_f\sim250$ GeV/c^2) and eventually, the fusion mechanism will dominate. However, analyses presented in Ref. [6] have somewhat tempered these optimistic conclusions. The preceding considerations have led us to perform calculations assuming two different values for the top-quark mass: $m_t=137$ GeV/c^2 and $m_t=350$ GeV/c^2. Our computations provide an alternative to the calculations using a parton shower approach (PYTHIA [14]) presented in Ref. [6].

Encouraged by the wealth of data taken at hadron colliders over the past decade, predictions for heavy flavor production cross sections as from perturbative QCD calculations have been carried out in detail and have shown remarkable agreement with experiment. Both UA1 [15] and CDF [16] Collaborations have presented extensive lepton analyses showing evidence for huge charm- and bottom-quark production. Heavy flavors are mainly probed through the observation of charged leptons resulting from semileptonic decays of heavy hadrons. Knowledge on semileptonic branching ratios enables "order of magnitude" extrapolations to obtain cross sections for open heavy flavor production, while B spectroscopy as pursued in several e^+e^- experiments has provided important input to probe heavy flavor bound states in hadronic interactions [16,17]. Bearing in mind the top-quark mass limits as reviewed above, we will concentrate on open heavy flavor production processes only, since bound states of heavy fermions ($\sim$150 GeV/c^2) will not appear as very narrow resonances and will even completely vanish when the fermion mass increases [18]. More importantly, because of its large mass the quark will decay weakly before hadronization can take place. The decay signature will then be dominated by so-called single quark decays, which makes them indistinguishable from directly produced heavy flavors. Thus pure leptonic decay modes are highly suppressed [19]. The enhancement in production cross section due to rescattering forces and resonance formation is also much less dramatic compared to e^+e^- interactions [20] since the convolution with the structure functions largely washes out the threshold behavior.

Lowest-order cross sections for heavy flavor production involving $2\to2$ subprocesses were calculated by Combridge [21]. Contributing Feynman diagrams are given in Fig. 1.

After the usual trace calculations and integration over phase space, one obtains a finite (due to $m_Q\neq0$) expression for the partonic $O(\alpha_s^2)$ cross section:

$$\begin{aligned}\sigma(q\bar{q}\to Q\bar{Q})&=\frac{8\pi\alpha_s^2}{27\hat{s}^2}(\hat{s}+2m_Q^2)\beta\,,\\ \sigma(gg\to Q\bar{Q})&=\frac{\pi\alpha_s^2}{3\hat{s}}\left[-\left[7+\frac{31m_Q^2}{\hat{s}}\right]\frac{\beta}{4}\right.\\ &\quad\left.+\left[1+\frac{4m_Q^2}{\hat{s}}+\frac{m_Q^4}{\hat{s}^2}\right]\ln\left[\frac{1+\beta}{1-\beta}\right]\right],\end{aligned}\tag{2.1}$$

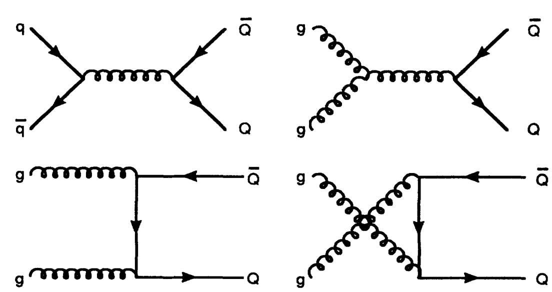

FIG. 1. $O(\alpha_s^2)$ quark-antiquark annihilation and gluon fusion contributions to heavy flavor pair production in hadronic interactions.

where α_s is the running strong coupling constant, m_Q is the mass of the heavy quark, and

$$\hat{s}=x_1 x_2 s \ ,$$
$$\beta=\left[1-\frac{4m_Q^2}{\hat{s}}\right]^{1/2} , \tag{2.2}$$

with x_1 and x_2 the partonic energy fractions taken from the beams with s the total center-of-mass energy. The observed cross section is then obtained by the usual convolution of the partonic cross section with the appropriate structure functions.

Although the formulas in Eq. (2.1) account for large part of the observed charm and bottom rates, higher-order terms give rise to important contributions as well. If we restrict ourselves to the discussion of tree-level diagrams only, the next-to-leading-order diagrams in α_s as presented in Fig. 2 may, depending on the mass of the heavy flavor, even dominate at present collider energies. For a large region of parton energy fractions and interaction scales, the most important contribution arises from the process

$$gg \to gg^* \to gQ\bar{Q} \tag{2.3}$$

for which one can find numerous discussions in the literature [22]. The $O(\alpha_s^3)$ tree-level calculations including massive fermion lines were first carried out by Kunszt and Pietarinen [23]. A more comprehensive calculation is presented in [24]. Both $O(\alpha_s^2)$ and $O(\alpha_s^3)$ matrix elements for massive fermions have been introduced in the Eurojet event generator for hard hadronic interactions [25], whereas parton shower models such as implemented in PYTHIA [14] are based on the $O(\alpha_s^2)$ matrix element to describe the core of the hard interaction. For an extensive discussion of the complete $O(\alpha_s^3)$ result (including virtual corrections) we refer the reader to [26] and references quoted therein. Introduction of virtual corrections into a Monte Carlo event generator is a highly nontrivial exercise in which numerical instabilities should be carefully evaluated. Nevertheless, by choosing an appropriate scale and cutoffs, finite results can be obtained based

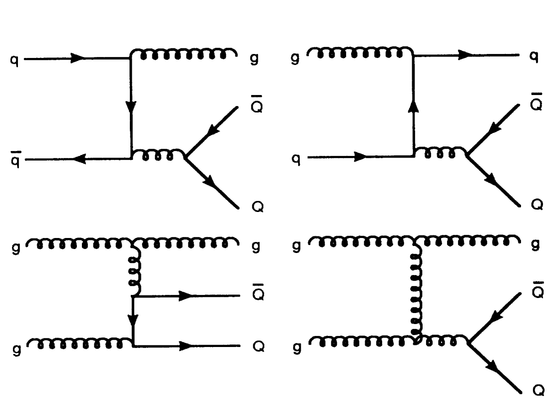

FIG. 2. A subset of $O(\alpha_s^3)$ quark-antiquark annihilation, (anti)quark-gluon scattering and gluon fusion contributions to heavy flavor pair production in hadronic interactions.

on tree-level contributions only. This procedure gives good agreement with the complete calculation. Compared to the $O(\alpha_s^3)$ corrections, electroweak corrections to the lowest-order QCD diagrams as a result of internal $W^{\pm}$, Z^0, H^0 lines (loops and vertices) are small ($\sim$10–20% [6]) and depend on both the Higgs-boson and top-quark masses, while the topology of the event is hardly affected.

The production of a single heavy quark through the exchange of a W boson in $pp(\bar{p})$ and ep scattering has already received quite some attention [27,28]. While the early studies were mainly aiming at clarifying charm production, higher beam energies make the contribution of boson-gluon fusion increasingly important in the study of far heavier fermions ($\gg m_b$). Lowest-order Feynman diagrams for single heavy flavor production in hadronic interactions are depicted in Fig. 3. After replacing the q, q' lines by e, ν lines, respectively, genuine single heavy flavor production in ep interactions can be obtained. Quark Q' couples to incoming quark Q with a probability proportional to the Cabibbo-Kobayashi-Maskawa (CKM) matrix element squared. Since the off-diagonal matrix elements are small in comparison with the diagonal elements, main contributions arise from the coupling among quarks belonging to the same doublet. For top quarks at the CERN Large Hadron Collider (LHC) ($\sqrt{s}=16$ TeV), we have explicitly calculated the Cabibbo-Kobayashi-Maskawa (CKM) suppressed production channels assuming $|V_{td}|=0.015$ and $|V_{ts}|=0.06$. Contributions to the total cross section are rather small: $\sim$0.1% and $\sim$1%, respectively. In the following computations we will therefore assume $|V_{tb}|=1$ and ignore off-diagonal contributions. The partonic differential cross sections in hadronic interactions, expressed in Lorentz-invariant Mandelstam variables, read

$$\frac{d\sigma(qQ\to qQ')}{d\Omega_2}=\frac{\alpha^2}{2\sin^4\theta_W}\frac{\hat{s}-m_{Q'}^2}{(\hat{t}-M_W^2)^2+M_W^2\Gamma_W^2} ,$$
$$\frac{d\sigma(q\bar{Q}\to q\bar{Q}')}{d\Omega_2}=\frac{\alpha^2(\hat{u}-m_Q^2)}{2\sin^4\theta_W(\hat{s}-m_Q^2)}\frac{\hat{u}-m_{Q'}^2}{(\hat{t}-M_W^2)^2+M_W^2\Gamma_W^2} \tag{2.4}$$

with $\hat{t}$ (M_W^2) defined as the invariant (on-shell) mass squared of the weak boson, while m_Q and $m_{Q'}$ are the masses of the incoming and outgoing heavy quarks, respectively. The width of the W, the Weinberg angle, and the (running) electromagnetic coupling constant are denoted by Γ_W, θ_W, and α. Light-quark masses can be safely neglected.

The dominant tree-level $O(\alpha_s)$ contributions (boson-

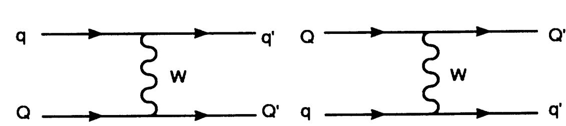

FIG. 3. $O(\alpha^2)$ quark-quark scattering contributions to heavy flavor production in hadronic interactions. To obtain the complete set of diagrams the charge conjugates and Cabibbo-Kobayashi-Maskawa mixings are to be included as well.

gluon fusion) can be easily obtained by replacing the lepton lines by (anti)quark lines in the expression of the matrix element for $ep \to egX \to \nu_e Q\bar{Q}'$ [27,28]. The corresponding $O(\alpha^2\alpha_s)$ Feynman graphs are given in Fig. 4.

Using the same conventions and variable definitions as in *ep* scattering—this choice is also sensible for the study of the singular structure of the expression for the short distance cross section—the following formula is obtained:

$$\frac{d\sigma}{d\Omega_3}=\frac{\alpha_s\alpha^2}{128\pi^2\sin^4\theta_W\hat{s}}\frac{R_V+R_A}{(q^2-M_W^2)^2+M_W^2\Gamma_W^2}\ , \tag{2.5}$$

where R_V (R_A) is even (odd) under the exchange $q\to\bar{q}$, $q'\to\bar{q}'$ and q^2 is the invariant mass of the virtual boson. The expressions for R_V and R_A are relatively simple:

$$\begin{aligned} R_V&=H_1^1\times G_1^1+H_2^1\times G_2^1+H_3^1\times G_3^1+H_4^1\times G_4^1\ , \\ R_A&=H_1^2\times G_1^2+H_2^2\times G_2^2 \end{aligned} \tag{2.6}$$

with the factors H_j^i and G_j^i only depending on kinematical variables. For a detailed discussion on the structure of these terms, we refer the reader to [2].

If we accommodate for the "*k* factors" in pure QCD production mechanisms by imposing appropriate cuts on the tree-level expression, a cross-section estimation for inclusive heavy flavor production can be obtained by simply adding the $O(\alpha_s^2)$ and $O(\alpha_s^3)$ contributions provided the scale and cutoffs are correctly chosen as discussed above [in Ref. [29] an analysis is presented in which $O(\alpha_s^4)$ tree-level contributions are included assuming massless fermions in the calculation of the $O(\alpha_s^4)$ matrix elements; mass effects affecting phase-space distributions have nevertheless correctly been taken into account]. In the case of single heavy flavor production we are able to explore a less drastic procedure, which turns out to be very suitable for Monte Carlo applications as well. If we ignore the α_s bremsstrahlung corrections to the diagrams in Fig. 3, which are partly suppressed due to the relatively large fermion masses involved [30], a cross-section estimation may be obtained by summing (after integration) Eqs. (2.4) and (2.5), provided the mass singularity in (2.5) is properly subtracted. In Eq. (2.5) we can isolate a logarithmic divergence, which is due to the internal bottom-quark propagator

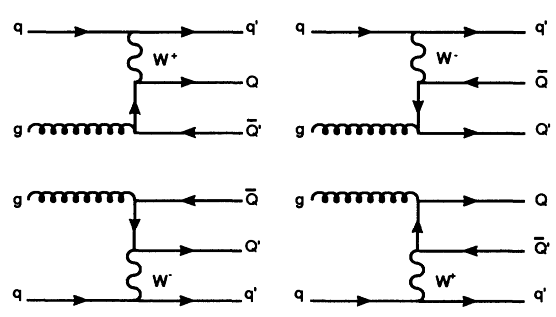

FIG. 4. $O(\alpha^2\alpha_s)$ (anti)quark-gluon scattering contributions to heavy flavor production via *W*-gluon fusion in hadronic interactions.

$$\frac{1}{(p_b-p_g)^2-m_b^2}=-\frac{1}{2(E_bE_g-|p_b||p_g|\cos\theta)}\equiv\frac{1}{z}\ , \tag{2.7}$$

and is regulated by the mass of the *b* quark (m_b). We can rewrite (2.5) such that this collinear singularity factorizes and we obtain an expression containing the convolution of the lowest-order result (2.4) with the first order in α_s distribution function of the *b* quark in a gluon. Not surprisingly, the collinear part describes just this domain of phase space where the bottom-quark distribution functions are applicable. In conclusion, lowest order and next-to-lowest order in α_s results can be added provided the divergence

$$\frac{d\sigma}{d\Omega_3}\sim\frac{\alpha^2\alpha_s}{8\pi^2\sin^4\theta_W}\frac{x^2+(1-x)^2}{(q^2-M_W^2)^2+M_W^2\Gamma_W^2}\frac{\hat{s}-m_t^2}{z} \tag{2.8}$$

is subtracted from the $O(\alpha^2\alpha_s)$ result in Eq. (2.5). Expression (2.8) is obtained from (2.5) after taking the limit $z\to 0$, where z is defined in (2.7) and x is the momentum fraction taken by the internal *b* quark from the gluon. In Fig. 5, we display curves representing the cross sections as obtained from the different mechanisms as function of the top-quark mass and collider energies. The scale at which the running electromagnetic coupling constant is calculated has been defined as $\hat{t}$, the invariant mass of the *W*. The scale for the strong coupling constant is defined as the sum of the outgoing parton transverse momenta and parton masses divided by the number of partons. We have used the structure function parametrizations as derived by Eichten *et al.*, set 1 ($\Lambda=200$ MeV) [31]. Apart from the Sp$\bar{p}$S, where phase-space suppression plays an important role, both $O(\alpha_s^2)$ and $O(\alpha_s^3)$ QCD processes dominate for top-quark masses below ~ 250 GeV/c^2. Except for the $O(\alpha_s^3)$ process, which contains divergences not regulated by quark masses, all cross sections are obtained after numerical integration over full phase space. Divergences in the $O(\alpha_s^3)$ calculation are removed by requiring a cutoff on the transverse momentum of the light quark (gluon). This cutoff is chosen such that, for different choices of structure functions and scales, there is a reasonable agreement with the complete $O(\alpha_s^3)$ calculations [26,32].

The enhancement of the *W*-gluon fusion mechanism over the pure QCD processes finds its main origin in the different $\hat{s}$ dependence of the amplitudes as becomes clear from comparing Eqs. (2.1) and (2.8), which is also emphasized in Ref. [33]. We have checked that our $O(\alpha^2\alpha_s)$ computation (*before* subtraction of the divergence) agrees with calculations by Zerwas *et al.*, who kindly provided us semianalytical formulas as well [34].

In Fig. 6 the transverse momentum (a) and pseudorapidity (b) distribution for $m_t=137$ GeV/c^2 are given at $\sqrt{s}=16$ TeV. The shape of both distributions is not very sensitive to changes in collider energy.

However, for the weak and semiweak production processes, increasing the top-quark mass moves the bumps in the rapidity distribution to larger values, while they become more pronounced as well [2]. Figure 6 is in good agreement with the results obtained using the $O(\alpha^2)$ matrix elements imposing a parton shower algorithm [6]. Nevertheless, one should emphasize that in order to analyze a more complete event topology, the matrix-element approach as pursued here can only be sufficient if higher orders are taken into account. One of the main advantages of our approach, we believe, is that we have obtained an improved estimation of the absolute cross section. Calculations based on lowest-order matrix elements together with parton showers can then be improved by introduction of "k factors." The calculation of the next order in α_s tree-level matrix elements is in progress [35].

III. WEAK AND ELECTROMAGNETIC NEUTRAL CURRENTS

As we discussed in the preceding section, one of the reasons why $W^{\pm}$-boson–gluon fusion takes over from

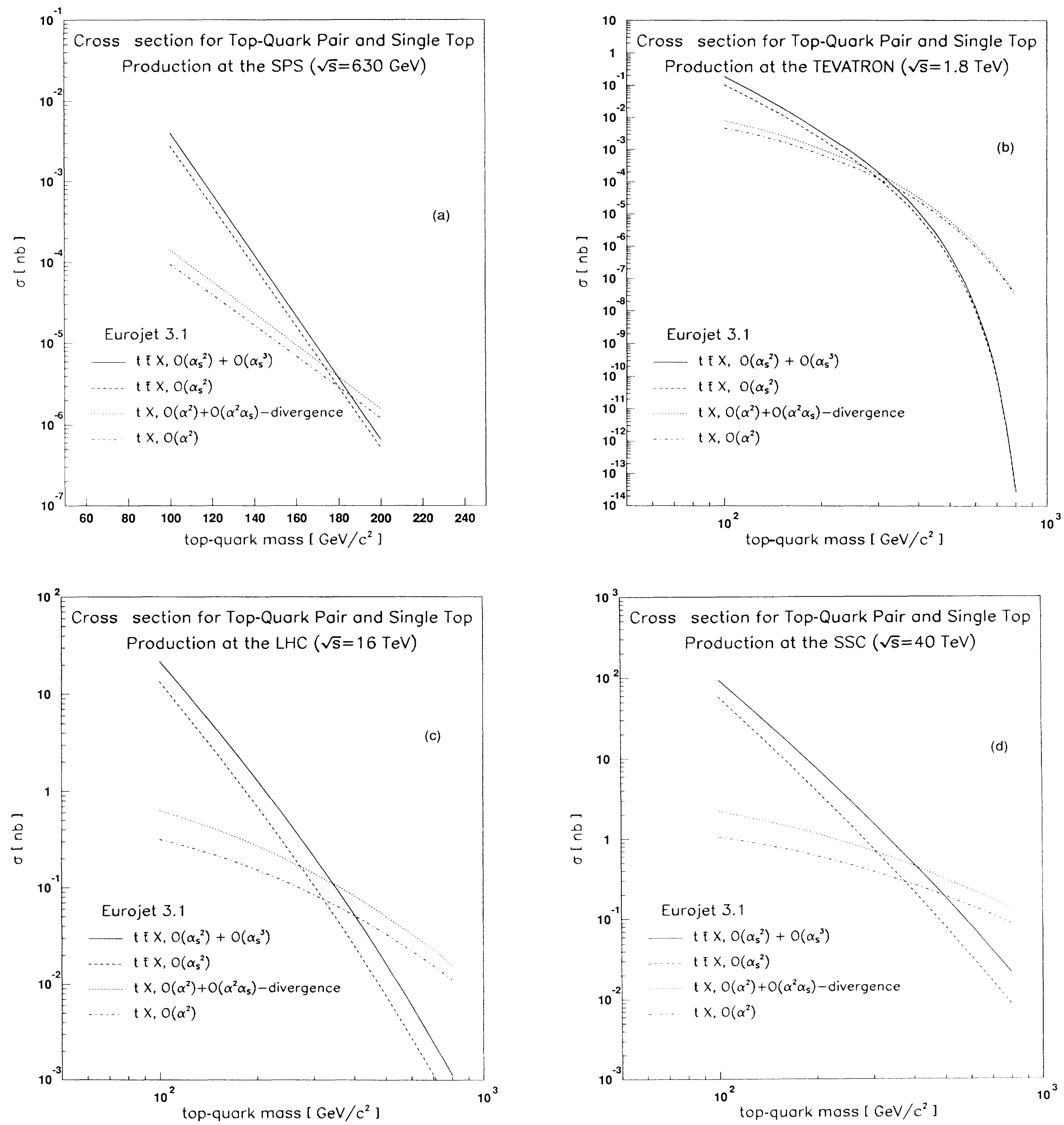

FIG. 5. Cross-section comparison for top-quark production in $pp(\bar{p}) \rightarrow tX$ as a function of the top-quark mass through pure QCD, weak, and boson-gluon fusion mechanisms at $\sqrt{s} = 630$ GeV (a), $\sqrt{s} = 1.8$ TeV (b), $\sqrt{s} = 16$ TeV (c), and $\sqrt{s} = 40$ TeV (d).

QCD processes when the mass of the heavy quark increases is the appearance of only one heavy quark in the final state. This also explains why the corresponding neutral process (Z^0-gluon fusion), which has been studied extensively in *ep* interactions [36,37], is not considered to be an important heavy flavor source in $p(\bar{p})$ interactions. Nevertheless, neutral-current processes may indeed become important when one allows for nonstandard couplings. In Sec. IV, we will discuss in detail a specific extension of the standard model in which neutral currents couple to two nonidentical heavy flavors (with large mass difference), without violating experimental constraints. Changing the couplings leads to expressions for the matrix elements squared, which are quite similar to the ones for $W^{\pm}$ exchange. We will not repeat them here (see, for instance, [28]).

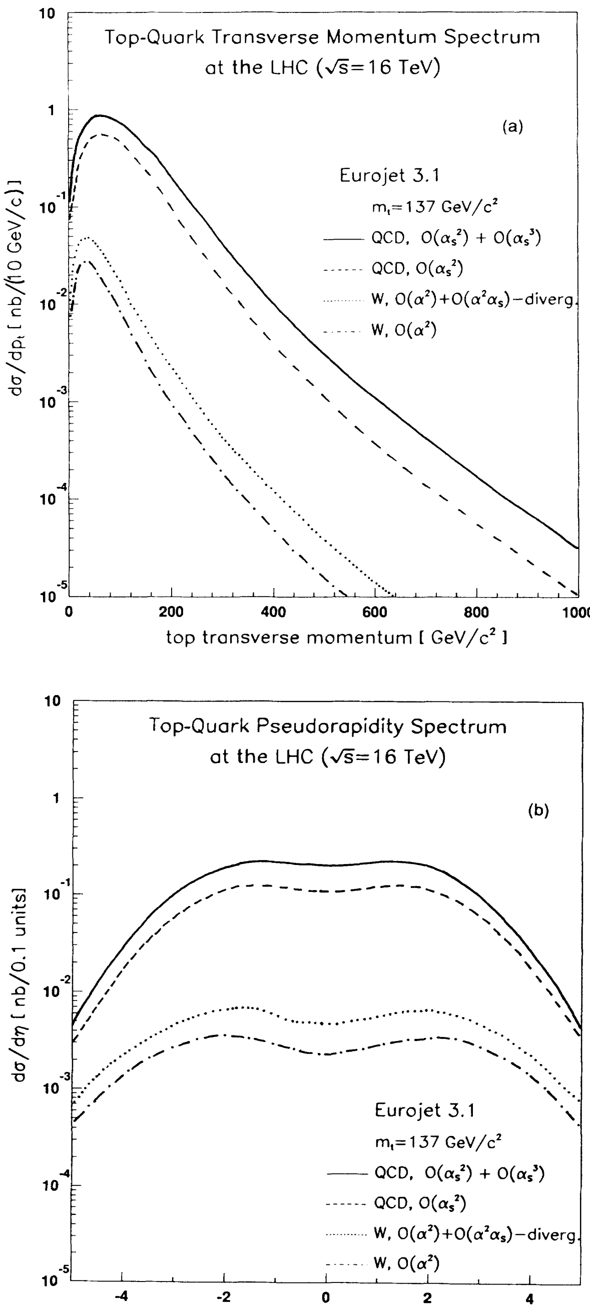

FIG. 6. Transverse-momentum (a) and pseudorapidity (b) distributions for the top quark as obtained from the different production mechanisms at $\sqrt{s} = 16$ TeV.

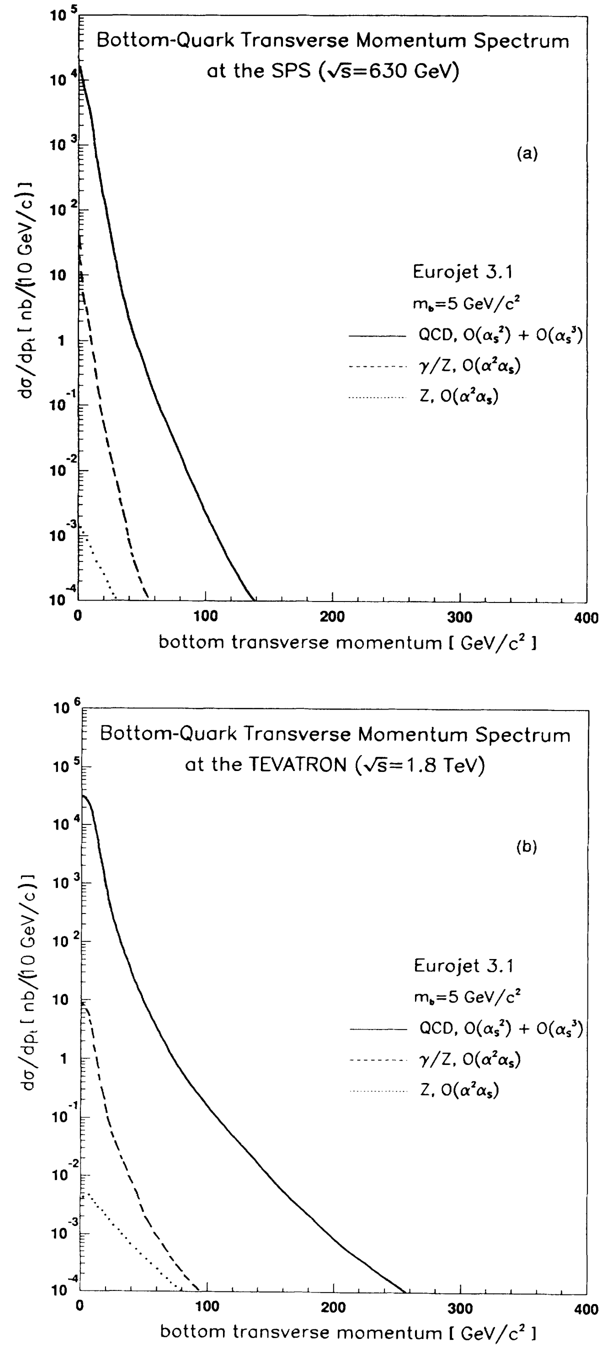

FIG. 7. Inclusive transverse-momentum distributions for bottom quarks pair-produced at the $Sp\bar{p}S$ (a) and Tevatron colliders (b) ($\sqrt{s} = 630$ GeV and $\sqrt{s} = 1.8$ TeV, respectively).

TABLE I. Bottom and top-quark pair production cross sections in $p\bar{p} \to Q\bar{Q}\ X$ (a) at the S$p\bar{p}$S ($\sqrt{s}=630$ GeV) and (b) at the Tevatron collider ($\sqrt{s}=1.8$ TeV). Top-quark pair production cross sections in $pp \to Q\bar{Q}\ X$ (c) at the LHC ($\sqrt{s}=16$ TeV) and (d) at the SSC ($\sqrt{s}=40$ TeV).

	$O(\alpha_s^2)$	$O(\alpha_s^3)$	$O(\alpha^2\alpha_s)$ (Z^0)	$O(\alpha^2\alpha_s)$ (γ/Z^0)
		(a)		
$b\bar{b}$	0.90×10^4	0.91×10^3	1.17×10^{-2}	0.3×10^1
$t\bar{t}(m_t=100\ \mathrm{GeV}/c^2)$	0.276×10^{-2}	0.13×10^{-2}	$\ll10^{-6}$	$\ll10^{-6}$
$t\bar{t}(m_t=200\ \mathrm{GeV}/c^2)$	0.8×10^{-6}	0.14×10^{-6}	$\ll10^{-6}$	$\ll10^{-6}$
		(b)		
$b\bar{b}$	0.289×10^5	0.65×10^4	0.13×10^{-1}	0.8×10^1
$t\bar{t}(m_t=137\ \mathrm{GeV}/c^2)$	0.191×10^{-1}	0.14×10^{-1}	0.11×10^{-4}	0.4×10^{-4}
$t\bar{t}(m_t=350\ \mathrm{GeV}/c^2)$	0.331×10^{-4}	0.16×10^{-4}	0.44×10^{-8}	0.7×10^{-8}
		(c)		
$t\bar{t}(m_t=137\ \mathrm{GeV}/c^2)$	0.365×10^1	0.27×10^1	0.93×10^{-2}	0.19×10^{-1}
$t\bar{t}(m_t=350\ \mathrm{GeV}/c^2)$	0.489×10^{-1}	0.52×10^{-1}	0.95×10^{-3}	0.11×10^{-2}
		(d)		
$t\bar{t}(m_t=137\ \mathrm{GeV}/c^2)$	0.178×10^2	0.13×10^2	0.46×10^{-1}	0.9×10^{-1}
$t\bar{t}(m_t=350\ \mathrm{GeV}/c^2)$	0.394×10^0	0.45×10^0	0.82×10^{-2}	0.09×10^{-2}

Another interesting aspect is that the neutral current receives contributions from photon exchange. Imagine that we replace the W lines in Fig. 4 by either a Z^0 or γ; we easily see that both processes will interfere. In order to get cross-section estimates, we have carried out a complete calculation and have carefully studied the effects of switching on and off γ/Z^0 interference.

In Fig. 7 we have depicted the transverse-momentum distributions of the bottom quarks at the S$p\bar{p}$S (a) and Tevatron (b). Contributions from pure QCD, electroweak, and weak production mechanisms are indicated separately.

In Table I, we present bottom and top cross sections (given in nb, statistically significant digits only) at various center-of-mass energies. Although the γ/Z^0 interference considerably enhances the neutral current signal, pure QCD pair production of either bottom or top quarks clearly dominates.

Increasing both quark mass and collider energy diminishes the effect of γ/Z^0 interference, as can be derived from the transverse-momentum distributions for top quarks given in Figs. 8(a)–8(c). In conclusion, although pure QCD pair production remains the dominant source for heavy flavors within the standard model, semiweak cross sections presented in Table I are non-negligible.

IV. E_6 FERMIONS

Within the standard model, constraints on the absolute mass scale and mass differences between individual members of doublets have been derived [3]. Although very elegant in its unification of electromagnetic, weak, and strong forces, the standard model does not answer major fundamental questions about the hierarchy problem, CP violation, etc. Therefore, in the past, several attempts have been made to extend the model in order to improve our understanding of unification and at the same time provide solutions for the as yet unresolved questions. Many distinct approaches have been advocated. A detailed discussion of the different models is beyond the scope of this paper; however, for a recent and rather complete overview we refer to Ref. [38]. An interesting modification of the standard model, which introduces heavy fermions in a natural way, is one where exotic quarks are vectorlike, avoiding usual limitations on quark masses. Recall that mass terms such as $m_0\bar{u}u = m_0(\bar{u}_R u_L + \bar{u}_L u_R)$ are scalars under any symmetry of the unbroken Lagrangian [39]. For the sake of argument, we limit ourselves to the discussion of a well-founded extension of the standard model based on superstring theories.

A coherent quantification of string theory can be obtained[1] by extending the number of space-time dimensions to $n>4$ [41]. We briefly describe the construction of compactified string theories. If we restrict ourselves to a minimal extension of the standard model, e.g., minimize the number of degrees of freedom, and if we demand that anomalies are absent, the obtained model is based on an internal $E_8 \otimes E_8$-symmetry group. In the low-energy limit, that is after compactification, the number of dimensions reduces to four. An $N=1$ supersymmetry and an effective E_6 grand unified gauge group are emerging [42]. The grand unification model based on the group E_6 has a fundamental representation in 27 dimensions. Then $SU(3)\otimes SU(2)\otimes U(1)$ spans a subspace within the E_6 group. The quantum numbers of the standard-model particles fix

[1]However, superstring theories have been derived in four dimensions by assigning the supplementary dimensions to an internal gauge group [40].

those of the remaining members of the E_6 multiplets. In practice this leads to the appearance of an isosinglet-color-triplet with charge-$\frac{1}{3}$ for each family [43] and one single new heavy neutral boson (we will not discuss any phenomenological implication of the coupling of these heavy quarks with the new boson [38]).

The behavior of these downlike quarks (D) can be described as follows. As far as strong interactions are concerned, D quarks couple to gluons in exactly the same way as standard model quarks. Differences appear in the electroweak sector. The isosinglet has pure vector couplings with the Z^0. However, having the same quantum numbers as the standard model downlike quarks, D quarks may mix with family members and, as some theories advocate, may provide an explanation for the mass differences among members of the same family [44]. Diagonalizing the mass matrix, we obtain the following mixing between the ordinary left-handed d and exotic D

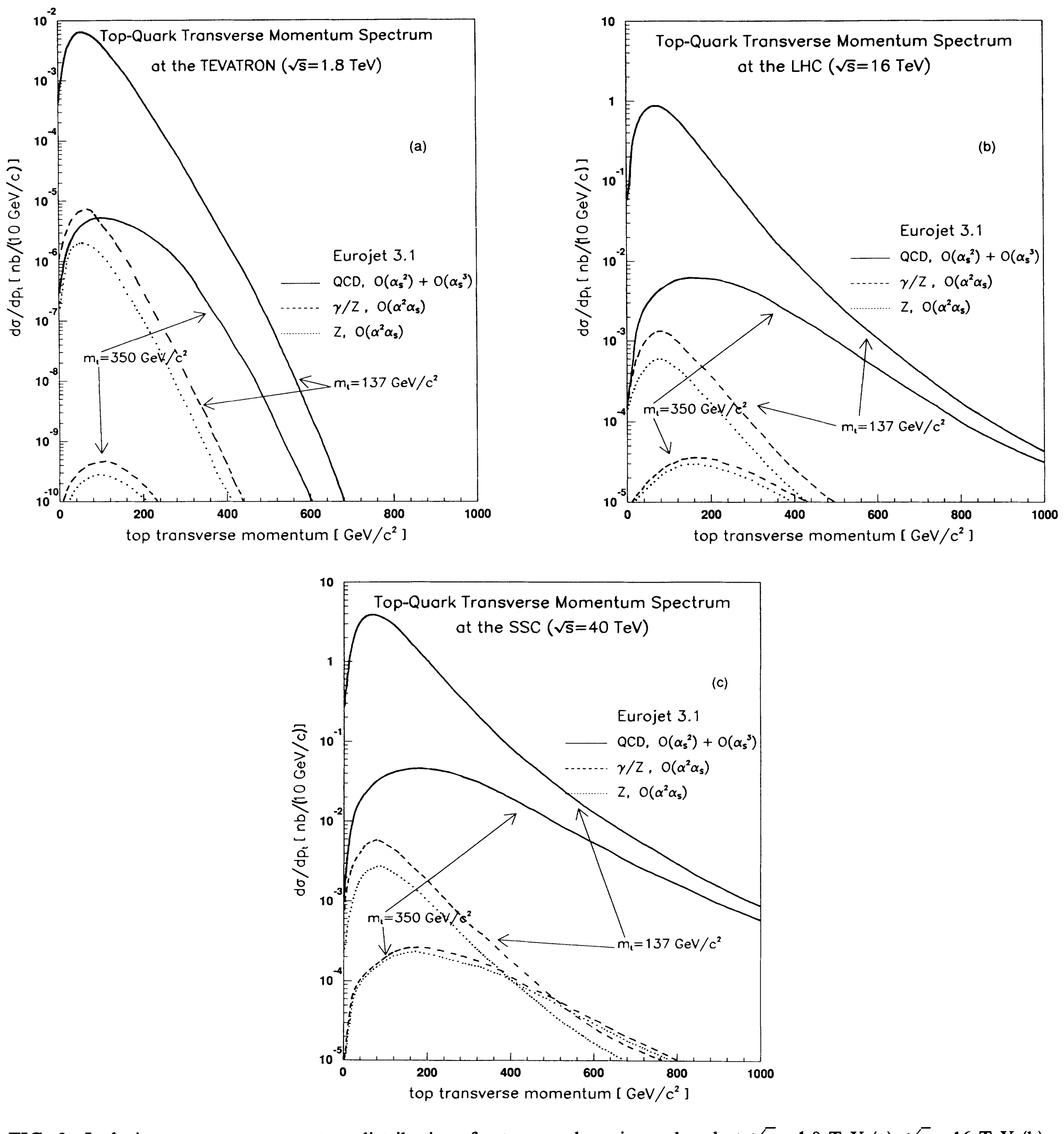

FIG. 8. Inclusive transverse-momentum distributions for top quarks pair produced at $\sqrt{s}=1.8$ TeV (a), $\sqrt{s}=16$ TeV (b), and $\sqrt{s}=40$ TeV (c).

quarks:

$$\begin{bmatrix} d \\ D \end{bmatrix}_L = \begin{bmatrix} \cos\phi & -\sin\phi \\ \sin\phi & \cos\phi \end{bmatrix} \begin{bmatrix} d_0 \\ D_0 \end{bmatrix}_L \tag{4.1}$$

with ϕ the mixing angle. The weak Lagrangian now contains terms describing the interaction between ordinary u, d, and exotic D quarks, and the obtained couplings with the gauge bosons are [39]

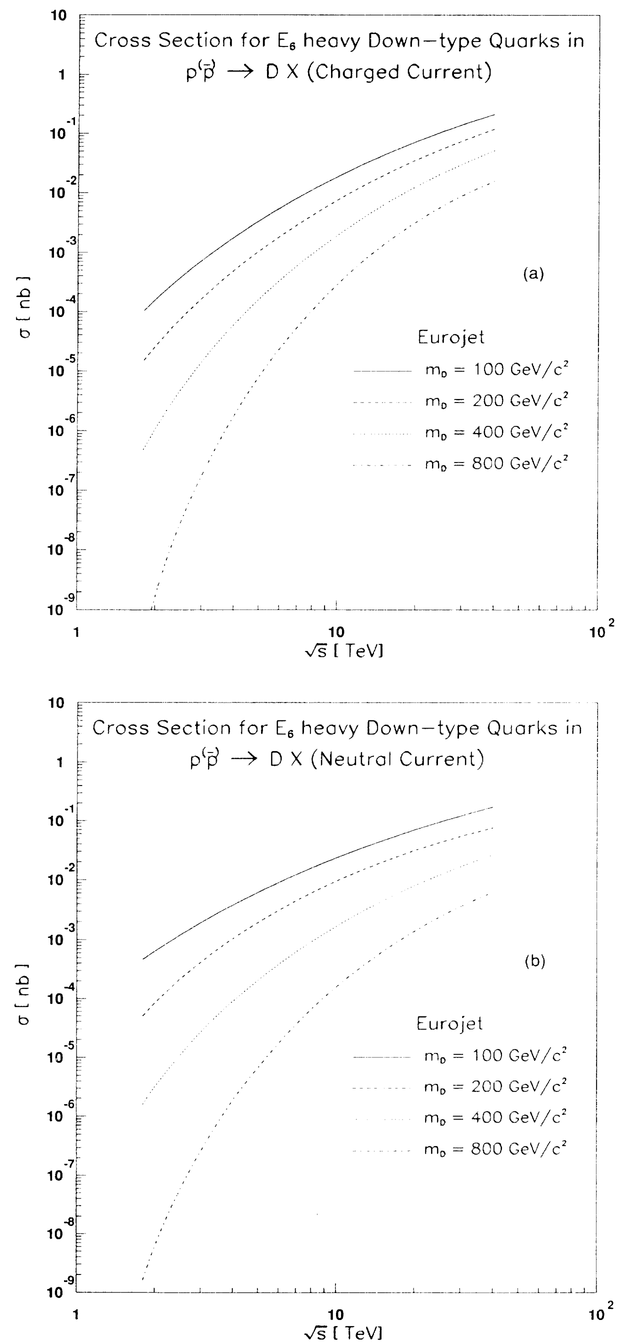

FIG. 9. Cross sections for E_6 D-quark production in $pp(\bar{p})\rightarrow DX$ as a function of the D-quark mass and center-of-mass energy through charged (a) and neutral (b) currents.

$$\mathcal{L}_{WDu} = g_W \sin\phi \bar{D}\gamma_\mu(1-\gamma_5)uW^\mu + \text{H.c.}\ ,$$
$$\mathcal{L}_{ZDd} = g_Z \sin\phi \cos\phi \bar{D}\gamma_\mu(1-\gamma_5)dZ^\mu + \text{H.c.}\ , \tag{4.2}$$

where

$$g_W = \frac{e}{2\sqrt{2}\,\sin\theta_W},\quad g_Z = \frac{e}{2\,\sin 2\theta_W}\ . \tag{4.3}$$

The ZDd term finds its origin in the isodoublet and isosinglet nature of the ordinary and exotic quarks, respectively. Several studies have been undertaken varying the unknown parameters in the model [43,45]. The mass of the D quarks is not subject to constraints from electroweak radiative corrections. This is due to the fact that large radiative corrections to Z^0 quantities are caused by the W coupling to heavy quarks which both belong to the same doublet. Since the E_6 quarks are singlet states, such couplings are absent and contributions from internal Du and Dd lines as from (4.2) to the ρ parameter are suppressed by the mixing strength and negligible in comparison with the corresponding $t\bar{b}$ term. In order to avoid experimentally unobserved flavor-changing neutral currents (FCNC's), one has to assume that each ordinary fermion only mixes with one unique heavy partner. Limits on the mixing strength are obtained from constraints imposed by experimental data on FCNC's. CP violation, $b\bar{b}$ mixing, and upper bounds on the branching ratios of several rare decays. Combining the limits leads to the following estimation: $\sin^2\phi = 0.01$ [45,46]. In addition, (linear or quadratic) "seesaw" models [4] provide relations between the masses of D quarks and their coupling with light quarks. Obviously, in contrast with the QCD cross sections derived in ([39], see also Ref. [6]), the

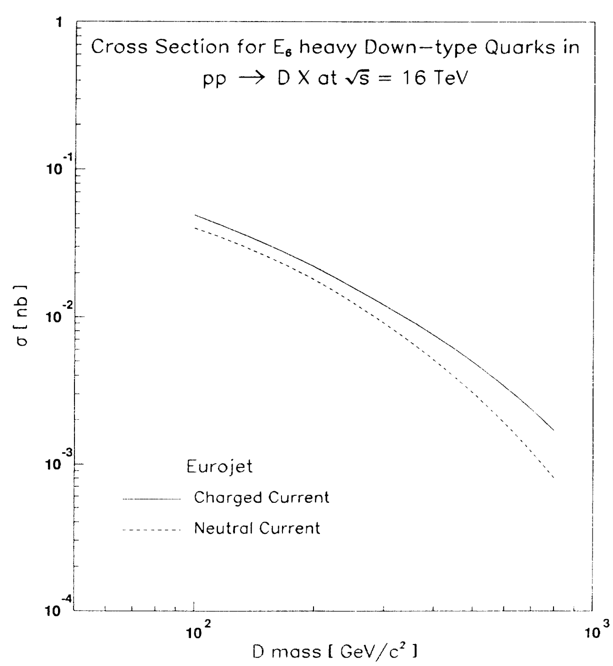

FIG. 10. Cross sections for E_6 D-quark production in $pp(\bar{p})\rightarrow DX$ as a function of the D-quark mass at $\sqrt{s}=16$ TeV through charged and neutral currents.

semiweak production mechanisms are therefore proportional to this badly estimated mixing angle.

In the following analysis, for simplicity, we consider the mixing of a third family D quark, b and t quarks only. In analogy with the calculations presented in Secs. II and III, we derive production cross sections for D quarks assuming different D masses using the similar production mechanisms. The relevant basic subprocesses are

$$qb \to qD, \quad qg \to q'D\bar{b}\ ,$$
$$qt \to qD, \quad qg \to qD\bar{t}\ , \tag{4.4}$$

where, since a direct coupling between a photon, D, and d quark is absent, the signal will not receive contributions from γ/Z^0 interference. Since the mass ratio between the top and D quark is taken to be smaller than the mass difference between the bottom and top quark, the subtraction scheme as pursued in Sec. II must be applied to the neutral-current processes in (4.4). From (4.2), it is straightforward to derive an expression for the cross section, which is quite similar to the formulas obtained for charged-boson exchange. $Dq(Q'q)$ production through a neutral boson reads [following the definitions given for Eq. (2.4)]

$$\frac{d\sigma}{d\Omega_2}=\frac{\alpha^2}{2\sin^4 2\theta_W(\hat{s}-m_Q^2)}\frac{1}{(\hat{t}-M_Z^2)^2+M_Z^2\Gamma_Z^2}$$
$$\times(4G_1\{\hat{t}[\hat{t}+2\hat{s}-(m_Q^2+m_{Q'}^2)]+2\hat{s}(\hat{s}-m_Q^2-m_{Q'}^2)+2m_Q^2m_{Q'}^2\}-8\varepsilon G_2\hat{t}(\hat{t}+2\hat{s}-m_Q^2-m_{Q'}^2)) \tag{4.5}$$

with $\varepsilon=1(-1)$ when one of the incoming partons is an (anti)quark. The $O(\alpha^2\alpha_s)$ mechanism for producing a $D\bar{d}q(Q\bar{q}'q)$ final state becomes

$$\frac{d\sigma}{d\Omega_3}=\frac{\alpha_s\alpha^2}{32\pi^2\sin^4 2\theta_W\hat{s}}\frac{G_1R_V+G_2R_A}{(q^2-M_Z^2)^2+M_Z^2\Gamma_Z^2} \tag{4.6}$$

with

$$G_1=v^2+a^2, \quad G_2=va \tag{4.7}$$

and R_V and R_A as in Eq. (2.6). The usual definitions apply for v and a, e.g., $v_u=1-\frac{8}{3}\sin^2\theta_W$ and $a_u=1$. Charged- and neutral-current cross sections for different D-quark masses, as a function of accelerator energies, are presented in Fig. 9.

In Fig. 10, we display the individual charged- and neutral-current contributions to single D-quark production at $\sqrt{s}=16$ TeV, as a function of the D-quark mass. We have not included the coefficients $\sin^2\phi$ and $\sin^2\phi$ $\cos^2\phi$. Semiweak cross sections should therefore be scaled with the appropriate mixing strength. The experimental observability of those exotic quarks depends on their decay properties. An analysis on tracing decays of E_6 fermions at LHC energies is presented in Ref. [39]. Our curves may serve as conservative estimates for the production of heavy downlike quarks which mix with either one of the first two families of the standard model as well. However, one should stress that for incoming light quarks the boson-gluon fusion calculations presented here cover only part of those cross sections. Inclusion of the $O(\alpha^2\alpha_s)$ boson-gluon fusion diagrams is indeed more delicate due to contributions from bremsstrahlung graphs [26].

V. CONCLUSIONS

We have demonstrated that boson-gluon fusion mechanisms may play an important role in the description of heavy flavor physics at hadron colliders, in particular if the top-quark mass exceeds the present bounds imposed by standard-model expectations.

Since superstring-inspired models provide hints on mass hierarchy and related puzzles, we have studied possible production of exotic fermions through the boson-gluon fusion mechanism. We have shown that the subtraction scheme, in which both $O(\alpha^2)$ and $O(\alpha^2\alpha_s)$ calculations are introduced to obtain a more precise determination of single top-quark production cross sections, can be applied for E_6 fermions as well. However, these string theories still have intrinsic parameters whose magnitude is largely unpredictable. The main uncertainty remains the mass of the exotic fermions, whereas the mixing strength with ordinary quarks plays an important role in the mass regime where the semiweak production mechanisms become dominant.

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