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Heavy quarks and squarks from W-gluon fusion

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A B S T R A C T

We discuss Wg-fusion as a source of heavy quark and squark pairs at very high energy hadron colliders. Effective W approximation is used to calculate the cross-sections analytically in the forward scattering configuration ; good agreement is obtained with exact numerical calculations. W-gluon fusion is found to be not nearly as important a production mechanism of heavy squarks as it is of heavy quarks. This is especially true when the mass-splitting within the $SU(2)_L$ doublet is small.

1. Introduction

It has been noted that some weak interaction processes can have production rates comparable to those of strong interactions at energies much larger than the W boson mass. Although the weak interaction coupling strength $\alpha_W = \alpha/\sin^2\theta_W = g^2/4\pi \simeq 0.037$ is small, there is a possibility for effective "strong" weak interactions between

(a) longitudinally polarized gauge bosons $W_L^{\pm}$ and Z_L^0 and the Higgs boson H[1] (the scalar sector of the standard model) with coupling $g_H = gm_H^2/2m_W$,

(b) heavy fermion F and $W_L^{\pm}$, Z_L^0 and H[2] with coupling $g_F = gm_F/2m_W$,

when the masses m_H or m_F are large compared to m_W. In fact partial wave unitary bounds indicate that perturbation theory in g_H or g_F is no more valid if $m_H \geq \mathcal{O}(1\ \mathrm{TeV})$[1] or if the mass splitting between $SU(2)_L$ fermion doublet $|m_U - m_D| \geq \mathcal{O}(0.5\ \mathrm{TeV})$[2].

At very high energies $E >> m_W$ $W^{\pm}$ and Z^0 can be treated as partons and it is sensible to talk about the $W^{\pm}$ or Z^0 content of colliding hadrons. The corresponding distribution functions for the longitudinal momentum fraction of W's in quarks are obtained in the effective W approximation (EWA)[3,4] and they are given separately for longitudinal (W_L) and transverse (W_T) polarization states of $W^{\pm}$

$$f_{W_L/q}(z) = \alpha_W/4\pi \cdot (1-z)/z$$
$$f_{W_T/q}(z, Q^2) = \alpha_W/8\pi\ [1+(1-z)^2]/z \cdot \log(Q^2/m_W^2) \qquad (1)$$

Now in the high energy limit W_L initiated processes can scale like $1/m_W^2$ (compared to $1/\hat{s}$ for typical QCD processes), which reflects the effective couplings g_H, g_F being proportional to $1/m_W$. Consequently at high energies and with $m_H, m_F > m_W$ W_L and Z_L initiated processes can become quite important; known examples are heavy Higgs production by WW-fusion[1,5,6], $W_L^+W_L^-, Z_L^0Z_L^0 \to H$, and heavy quark production by W-gluon fusion[3,7], $W_L^+g \to U\bar{D}$. Here we generalize the Wg-fusion process to include the production of scalar quarks, $W_L^+g \to \tilde{U}\tilde{D}^*$, in the minimal supersymmetric extension of the standard model.

The main mechanisms of heavy quark pair production in hadron collisions are the QCD processes $gg \to Q\bar{Q}$ and $q\bar{q} \to Q\bar{Q}$; the cross-sections are evaluated and presented numerically if ref.8). However, Willenbrock and Dicus[7] find the weak production $Wg \to U\bar{D}$ to be as important as strong production in very high energy $pp/p\bar{p}$ collisions if $m_U - m_D$ is large enough : Wg-fusion exceeds QCD production for $\sqrt{s}$ = 10-40 TeV if $m_{\tilde U} - m_{\tilde D} > 300$-400 GeV, and even at $\sqrt{s}$ = 2 TeV top-quark production by $Wg \to t\bar b$ is larger than that by $q\bar q \to W \to t\bar b$ if $m_t > 100$ GeV. The surprising importance of the Wg-fusion process is explained by several enhancement factors which compensate for the small ratio of couplings of $\mathcal{O}(10^{-4})$ for Wg vs. QCD production mechanisms[7]. The first factor is the $1/m_W^2$ -scaling of the weak process already discussed above. Secondly, Wg-luminosity L_{Wg} is the convolution of qg-luminosity L_{qg} with the distribution functions of eq.(1) and L_{qg} is almost as large as gg-luminosity L_{gg} in $pp/p\bar p$-collisions for center-of-mass energies in the TeV region. Willenbrock and Dicus calculate the $Wg \to U\bar D$ cross-section including the initial quark leg $q \to Wq'$. We find good agreement between our calculations using EWA and results by Willenbrock and Dicus, whereas the result in Ref. 3) clearly underestimates the cross-section.

Heavy squark production in hadronic collisions is discussed in Ref. 9), where cross-sections are evaluated and presented numerically for the main supersymmetric QCD production mechanisms $gg \to \tilde Q_i \tilde Q_i^*$, $q_i \bar q_j \to \tilde Q_i \tilde Q_j^*$ and $q_i q_j \to \tilde Q_i \tilde Q_j$. Although only the second of these processes can have the same final state as in $Wg \to \tilde U \tilde D^*$, the $\tilde U$, $\tilde D$, $\tilde U^*$ and $\tilde D^*$ squark decays may be difficult to distinguish. Following Eichten et al. (Ref. 8) we compare the Wg-fusion process with the sum over the different final states in supersymmetric QCD processes of squark pair production. The heavy quark mass range of interest here is 100 GeV $< m_{U,D} <$ 700 GeV. Below m_Q = 100 GeV the Wg-process is unimportant, and the upper limit reflects the unitarity bounds[2]. We also keep in mind constraints on the mass splitting $\Delta m = |m_U^2 - m_D^2|^{1/2} \le 350$ GeV from neutral current phenomenology[8]. Similarly, we take the heavy squark masses to be in the range 100 GeV $< m_{\tilde U,\tilde D} <$ 700 GeV. As W couples only to left-handed quarks Q_L and their scalar partners $\tilde Q_L$, the scalar masses $m_{\tilde Q}$ discussed here refer to the masses of the squarks $\tilde Q_L$. For strong production we use estimates with equal masses for the left- and right-handed squarks. Basically the masses of squarks (as of all scalar particles) are free parameters of the theory, which must be introduced explicitly in the scalar superpotential. In a class of models based on N=1 supergravity coupled to a grand unified theory (GUT) all scalar masses are assumed equal (= m) at the GUT unification scale[10]. In the renormalization group evolution to low energies the top-quark Yukawa coupling can drive one or more (Higgs) scalar masses squared to a negative value giving rise to spontaneous symmetry breaking of supersymmetry. Typical masses for the squarks are then[11] $m_{\tilde U}^2 = m^2 + (-\frac{1}{2} + \frac{2}{3}\sin^2\theta_W)m_Z^2$ and $m_{\tilde D}^2 = m^2 + (\frac{1}{2} - \frac{2}{3}\sin^2\theta_W)m_Z^2$ i.e. the mass difference $\Delta m = |m_{\tilde U}^2 - m_{\tilde D}^2|^{1/2}$ is expected to be at most of $\mathcal{O}(m_Z)$. Present experimental lower limits on the squark masses are $m_{\tilde Q} \ge 20$ GeV from e^+e^--collisions[12] and $m_{\tilde Q} \ge \mathcal{O}(40$-60 GeV) from monojet events in $p\bar p$-collisions[13].

2. Evaluation of $Wg \to U\bar{D}$

We calculate the cross-section $\sigma(Wg \to U\bar{D})$ from Feynman diagrams of Fig. 1 using EWA structure functions for W in the initial state. Sum over the W polarization vectors $\epsilon_\mu^\lambda(p_W)$ gives

$$\sum_{\lambda=T,L} \epsilon_\mu^\lambda (\epsilon_{\mu'}^\lambda)^* = -g_{\mu\mu'} + p_{W\mu} p_{W\mu'}/m_W^2 \tag{2}$$

where in the high energy limit the first term can be identified with transversely polarized W_T and the last term with longitudinally polarized W_L. More generally a covariant polarization vector can be defined for W_L [14)]

$$\eta_\mu^L = p_{W\mu}/m_W - p_{g\mu}\, m_W / p_W\cdot p_g \tag{3}$$

with $\eta_\mu^L \eta_{\mu'}^L$ projecting out the contribution of longitudinal W_L. This definition is used in our numerical calculations, but in the analytical expressions below we call for simplicity the $p_{W\mu} p_{W\mu'}/m_W^2$ contribution as coming from W_L and the $-g_{\mu\mu'}$ contribution from W_T.

The spin-averaged squared matrix elements are given by

$$\sum_{\text{spin ave}} |\mathcal{M}(W_T g \to U\bar{D})|^2 = 8\pi^2 \alpha_s \alpha_W \cdot \left[U/T + T/U + 2(m_U^2 + m_D^2 - m_W^2)\cdot F \right]$$

$$\sum_{\text{spin ave}} |\mathcal{M}(W_L g \to U\bar{D})|^2 = 8\pi^2 \alpha_s \alpha_W \cdot \left[(m_U^2 + m_D^2)/2m_W^2 \cdot (2 + U/T + T/U) + (\Delta^4/m_W^2 - m_U^2 - m_D^2)\cdot F \right] \tag{4}$$

where we introduced the kinematical variables

$$T = \hat{t} - m_D^2 = (p_W - p_U)^2 - m_D^2$$
$$U = \hat{u} - m_U^2 = (p_W - p_D)^2 - m_U^2 \tag{5}$$
$$S = \hat{s} - m_W^2 = (p_W + p_g)^2 - m_W^2$$

which obey $S+T+U = 0$, and where $\Delta^2 = m_U^2 - m_D^2$. Note factorization of the term

$$F = 1/T + 1/U + m_D^2/T^2 + m_U^2/U^2 - (m_U^2 + m_D^2 - m_W^2)/UT$$
$$= -\left(1 + m_D^2/T + m_U^2/U + m_W^2/S\right) \cdot S/UT \tag{6}$$

In the limit of massless quarks the W_L-contribution vanishes and we retain the result for $\sigma(\gamma\gamma \to e^+e^-) \propto \hat{u}/\hat{t} + \hat{t}/\hat{u}$. Also note the presence of both $(m_U^2 + m_D^2)/m_W^2$ and $(m_U^2 - m_D^2)^2/m_W^2$ terms in the W_L-contribution, which dominates over the W_T-contribution by factors of $\sigma(m_{U,D}^2/m_W^2)$ in the limit $m_{U,D} >> m_W$.
The subprocess cross-section

$$d\hat{\sigma}/d\hat{t} = \overline{|\mathcal{M}|^2} / 16\pi S^2 \tag{7}$$

enters directly the observable cross-section when the transverse momentum k_T and the rapidities y_U, y_D of the heavy quarks are fixed

$$d\sigma/dy_U dy_D d^2k_T = x_1 x_2\, g(x_1)\, W_{L(T)}(x_2) \cdot 1/\pi\, d\hat{\sigma}/d\hat{t} \tag{8}$$

Here $g(x_1)$ is the gluon distribution and $W_{L(T)}(x_2)$ is the convolution of the EWA structure function of eq.(1) with the singlet quark distribution $q^S(x)$. The longitudinal momentum fractions are given by

$$x_1 = (M_U^T e^{y_U} + M_D^T e^{y_D})/\sqrt{s}$$
$$x_2 = (M_U^T e^{-y_U} + M_D^T e^{-y_D})/\sqrt{s} - m_W^2/x_1 s \tag{9}$$

where $M_i^T = \sqrt{m_i^2 + k_T^2}$ is the quark transverse mass. Now $\hat{s}$, $\hat{t}$ and $\hat{u}$ and consequently $d\hat{\sigma}/d\hat{t}$ depend only on the rapidity difference $\Delta y = y_U - y_D$. For $\Delta y = 0$ we get especially simple expressions for the cross-sections : for example, when $m_W << m_{U,D}$ F is given by

$$F = -k_T^2/(M_U^T M_D^T)^2 \tag{10}$$

which term favors k_T-values of order $\sqrt{m_U m_D}$. In the same limit the subprocess cross-sections are

$$
\begin{aligned}
d\hat{\sigma}/d\hat{t}\,(W_T g \to U\bar{D}) &= \pi/2 \cdot \alpha_s \alpha_w\, (M_U^T + M_D^T)^{-4} / M_U^T M_D^T \\
&\quad \cdot \left[(M_U^T - M_D^T)^2 + 2\,(m_U^2 m_D^2 + k_T^4) / M_U^T M_D^T \right] \\
d\hat{\sigma}/d\hat{t}\,(W_L g \to U\bar{D}) &= \pi/2 \cdot \alpha_s \alpha_w\, (M_U^T + M_D^T)^{-4} / M_U^T M_D^T \\
&\quad \cdot \left[(m_U^2 + m_D^2)/2m_W^2 \cdot (M_U^T + M_D^T)^2 - \Delta^4/m_W^2 \cdot k_T^2 / M_U^T M_D^T \right]
\end{aligned}
\tag{11}
$$

Keeping only the leading terms when $\hat{s} \geq (m_U + m_D)^2 \gg m_W^2$ integration over $\hat{t}$ in eq.(7) gives for $W_L g \to U\bar{D}$

$$
\begin{aligned}
\hat{\sigma}(W_L g \to U\bar{D}) &= \pi/2 \cdot \alpha_s \alpha_w \cdot \hat{s}^{-3} \cdot \\
&\quad \Big\{ \Big[(m_U^2 + m_D^2)/2m_W^2 \cdot \hat{s}^2 - \Delta^4/m_W^2 \cdot (\hat{s} - m_U^2 - m_D^2) \Big] \\
&\quad \cdot (L_{UD} + L_{DU}) + 2\Delta^4/m_W^2 \cdot \sqrt{\lambda} \Big\}
\end{aligned}
\tag{12}
$$

where

$$
\begin{aligned}
L_{UD} &= \log \left[(\hat{s} + m_U^2 - m_D^2 + \sqrt{\lambda}) / (\hat{s} + m_U^2 - m_D^2 - \sqrt{\lambda}) \right] \\
L_{DU} &= L_{UD}\,(m_U \leftrightarrow m_D) \\
\lambda &= (\hat{s} + m_U^2 - m_D^2)^2 - 4\hat{s}\, m_U^2
\end{aligned}
\tag{13}
$$

Finally the total cross-section for the dominating longitudinal W_L scattering is $(\tau = (\hat{s} - m_W^2)/s)$

$$
\begin{aligned}
\sigma(pp/p\bar{p} \to W_L g \to U\bar{D}) &= \int_{\tau_{min}}^{1} dx \int_{\tau_{min}/x}^{1} dy \int_{\tau_{min}/x/y}^{1} dz \\
&\quad \cdot g(x)\, q^s(y)\, f_{W_L/q}(z) \cdot \hat{\sigma}(\hat{s} = xyz \cdot s + m_W^2)
\end{aligned}
\tag{14}
$$

In Fig. 2 we present $\sigma(W_L g \to U\bar{D})$ for $pp/p\bar{p}$ -collisions at $\sqrt{s} = 10$ TeV and $m_D = 100$ GeV as a function of m_U, using Duke-Owens set 1 structure functions[15] with $Q^2 = m_W^2$ and using the $\eta_\mu^L \eta_\nu^L$ -projector for W_L.

Using the $q^S \hat{=} F_2(x,Q^2)$ structure function we implicitely sum over the two processes $W^+g \to U\bar{D}$ and $W^-g \to \bar{U}D$, and to compare with ref. 7) our results must be divided by 2. The resulting curve lies about 25% above that of Willenbrock and Dicus; this small difference in normalization can be attributed to the softer gluon distribution used here (Duke-Owens set 1 vs. EHLQ set 2[8] in ref. 7)).
The $W_L g$-fusion process dominates over the QCD production processes of heavy quarks at large enough $\Delta m = \sqrt{m_U^2 - m_D^2}$; the cross-over with the $\sigma(QCD \to Q\bar{Q})$ curve (dashed line in fig. 2 : from ref. 8, fig. 163) happens at $\Delta m = 250$ GeV.

3. Evaluation of $Wg \to \tilde{U}\tilde{D}^*$

In addition to the $\hat{t}$- and $\hat{u}$-pole diagrams of Fig. 1 we also have the point-like interaction between W, gluon and scalar quarks in the minimal supersymmetric standard model[16]. The $-g_{\mu\mu'}$ and $p_{W_\mu}p_{W_{\mu'}}/m_W$ projections of the squared matrix element are

$$\sum_{\text{spin ave}} |\mathcal{M}(W_T g \to \tilde{U}\tilde{D}^*)|^2 = 8\pi^2 \cdot \alpha_s \alpha_W \cdot \left[2 + \left(2(m_U^2+m_D^2) - m_W^2\right)\cdot F\right] \tag{15}$$

$$\sum_{\text{spin ave}} |\mathcal{M}(W_L g \to \tilde{U}\tilde{D}^*)|^2 = -8\pi^2 \cdot \alpha_s \alpha_W \cdot \Delta^4/m_W^2 \cdot F$$

In this section the subscripts U and D refer to the scalar quarks $\tilde{U}$ and $\tilde{D}$ i.e. *we drop the twiddle to simplify notation. As there is a complete factor*-ization of the F-term also for the $\eta^L_\mu \eta^L_{\mu'}$ projection, we give this in analytical form

$$\sum_{\text{spin ave}} |\mathcal{M}(W_L g \to \tilde{U}\tilde{D}^*)|^2 = -8\pi^2 \alpha_s \alpha_W \left[\Delta^2/m_W + m_W(T-U)/(T+U)\right]\cdot F \tag{16}$$

In contrast to the case with $W_L g \to U\bar{D}$ there is no $(m_U^2 + m_D^2)/m_W^2$ -term for longitudinal W_L; therefore we can expect a stronger dependence on Δ^2 for the Wg-fusion process into squarks. The subprocess cross-section as a function of y_U, y_D and k_T is given for $\Delta y = 0$ and $m_{U,D} \gg m_W$ by

$$d\hat{\sigma}/d\hat{t}\,(W_T g \to \tilde{U}\tilde{D}^*) = \pi/2\ \alpha_s \alpha_W \,(M_U^T + M_D^T)^{-4} \cdot 2\left(m_U^2 m_D^2 + k_T^4\right)/(M_U^T M_D^T)^2 \tag{17}$$

$$d\hat{\sigma}/d\hat{t}\,(W_L g \to \tilde{U}\tilde{D}^*) = \pi/2\ \alpha_s \alpha_W \,(M_U^T + M_D^T)^{-4}\, \Delta^4/m_W^2 \cdot k_T^2/(M_U^T M_D^T)^2$$

The ratio $\hat{R} = d\hat{\sigma}(W_L g \to \tilde{U}\tilde{D}^*)/d\hat{\sigma}(W_T g \to \tilde{U}\tilde{D}^*)$ vanishes at $k_T = 0$ and reaches its maximum value at $k_T = \sqrt{m_U m_D}$

$$\hat{R}_{max} = \Delta^4 / (4 m_W^2 m_U m_D) \tag{18}$$

Integration over $\hat{t}$ gives the following leading terms when $\hat{s} \geq (m_U + m_D)^2 >> m_W^2$

$$\hat{\sigma}(W_L g \to \tilde{U}\tilde{D}^*) = \pi/2 \cdot \alpha_s \alpha_W \cdot \hat{s}^{-3} \cdot \Delta^4/m_W^2 \cdot \left[(\hat{s} - m_U^2 - m_D^2) \cdot (L_{UD} + L_{DU}) - 2\sqrt{\lambda} \right] \tag{19}$$

Note that the Δ^4/m_W^2 -term in eqs (15), (17) and (19) for $W_L g \to \tilde{U}\tilde{D}^*$ enters also the expressions for $W_L g \to U\bar{D}$ in eqs (4), (11) and (12) but with opposite sign; this is related to the sign difference between a fermion and a scalar loop.

Again we integrate numerically the exact expression for the $\tilde{U}_L\tilde{D}_L^*$ projection of the $W_L g \to \tilde{U}\tilde{D}^*$ process. In Fig. 3 we present $\sigma(W_L g \to \tilde{U}\tilde{D}^*)$ for $pp/p\bar{p}$ collisions at $\sqrt{s} = 10$ TeV and $m_D = 100$ GeV as a function of m_U. Because of the proportionality to Δ^4/m_W^2 of the $W_L g \to \tilde{U}\tilde{D}^*$ cross-section the result looks qualitatively very different from that for $W_L g \to U\bar{D}$. $\sigma(W_L g \to \tilde{U}\tilde{D}^*)$ falls down rapidly at smaller $|m_U - m_D|$ values and reaches 25% of $\sigma(W_L g \to U\bar{D})$ for corresponding heavy quark masses only at very large values of Δm.

Comparison with the supersymmetric QCD production processes of heavy squarks shows that no cross-over happens below $\Delta m \leq 700$ GeV with the $\sigma(s\text{-QCD} \to \tilde{Q}\tilde{Q}$, $\tilde{Q} = \tilde{U},\tilde{D},\tilde{U}^*,\tilde{D}^*)$ curve (dashed line in Fig. 3; from ref. 8, fig. 198).

4. Conclusions

We have presented a detailed calculation for the W-gluon fusion processes of heavy quark and squark pair production. Analytical expressions for the cross-sections help us to understand the dependence on the $SU(2)_L$ mass splitting $\Delta m = |m_U^2 - m_D^2|^{1/2}$; for squark production this dependence is especially strong as the $W_L g \to \tilde{U}\tilde{D}^*$ cross-section is proportional to Δm^4. This result is in fact more general, and implies that unitarity limits for scattering processes involving squarks and W_L, Z_L and H_0 give constraints only on the mass splittings $m_{\tilde{Q}1}^2 - m_{\tilde{Q}2}^2$ and not on the absolute scale of squark masses $m_{\tilde{Q}}$. This is to be compared with the case for heavy quarks, where unitarity constrains the mass scale m_Q above which weak interactions cannot be treated perturbatively.

The equivalent W approximation is tested against calculations for the $Wg \to U\bar{D}$ process including the initial state quark line. We find agreement within 25%, and taking into account possible differences arising from the use of different set of structure functions, we conclude that EWA agrees within a factor 2 with the exact calculation at $\sqrt{s}$ and m_Q values discussed above.

The importance of the $Wg \to U\bar{D}$ process is confirmed; it dominates over the QCD production processes of heavy quarks in $pp/p\bar{p}$ collisions at $\sqrt{s} = 10$ TeV if $|m_U - m_D| \gtrsim 250$ GeV. However, the strong production of squarks overwhelms the $Wg \to \tilde{U}\tilde{D}^*$ process in the mass and energy range discussed here; no cross-over is found in $pp/p\bar{p}$ collisions at $\sqrt{s} = 10$ TeV below $|m_U - m_D| \leq 700$ GeV. Keeping in mind the small value of $\Delta m \leq \mathcal{O}(m_Z)$ expected from supersymmetry model building, as discussed in the introduction, we must conclude that the W-gluon fusion process will not give nearly as important a contribution to squark pair production as it does give to heavy quark production.

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Figure Captions

Fig. 1 Feynman diagrams for W-gluon fusion into quark pairs ($U\bar{D}$) or squark pairs ($\tilde{U}\tilde{D}^*$).

Fig. 2 Total cross-section for $W^{\pm}g \to U\bar{D}, \bar{U}D$ in pp or $p\bar{p}$ collisions at $\sqrt{s}$ = 10 TeV as a function of heavy quark mass m_U (m_D = 100 GeV). Also shown is the heavy quark pair production cross-section from QCD (dashed line, refs 7, 8).

Fig. 3 Total cross-section for $W^{\pm}g \to \tilde{U}\tilde{D}^*, \tilde{U}^*\tilde{D}$ in pp or $p\bar{p}$ collisions at $\sqrt{s}$ = 10 TeV as a function of heavy squark mass m_U (m_D = 100 GeV). Also shown is the heavy squark pair production cross-section from supersymmetric QCD (dashed line, ref. 8).

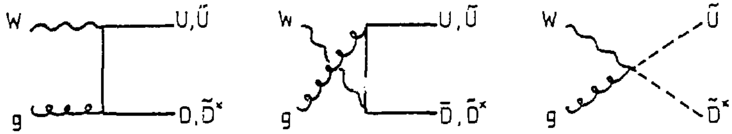
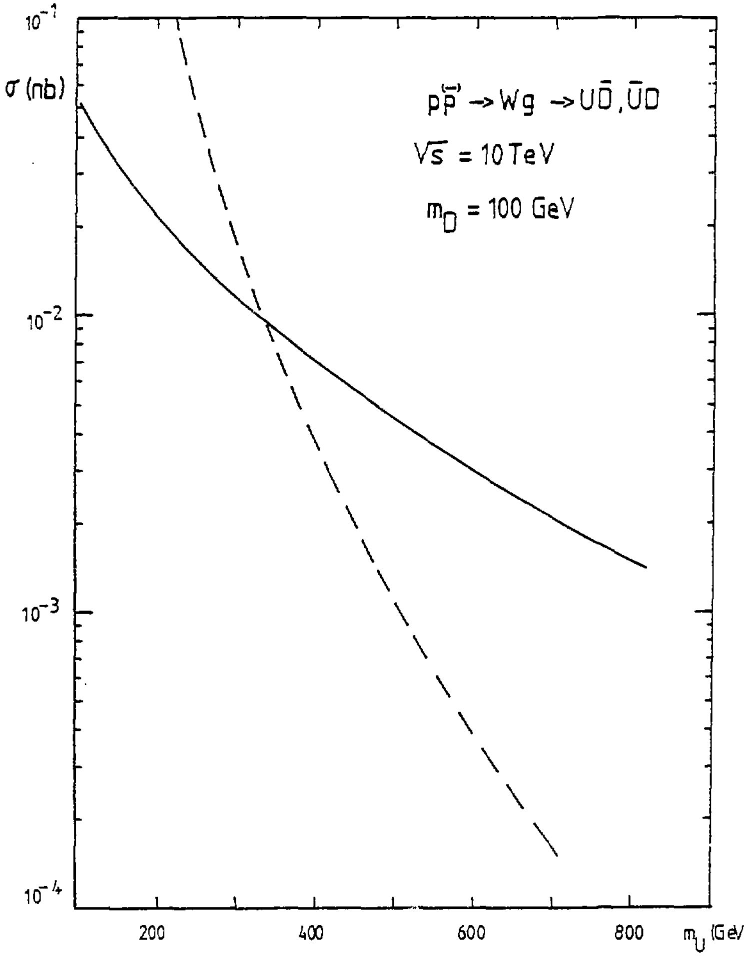
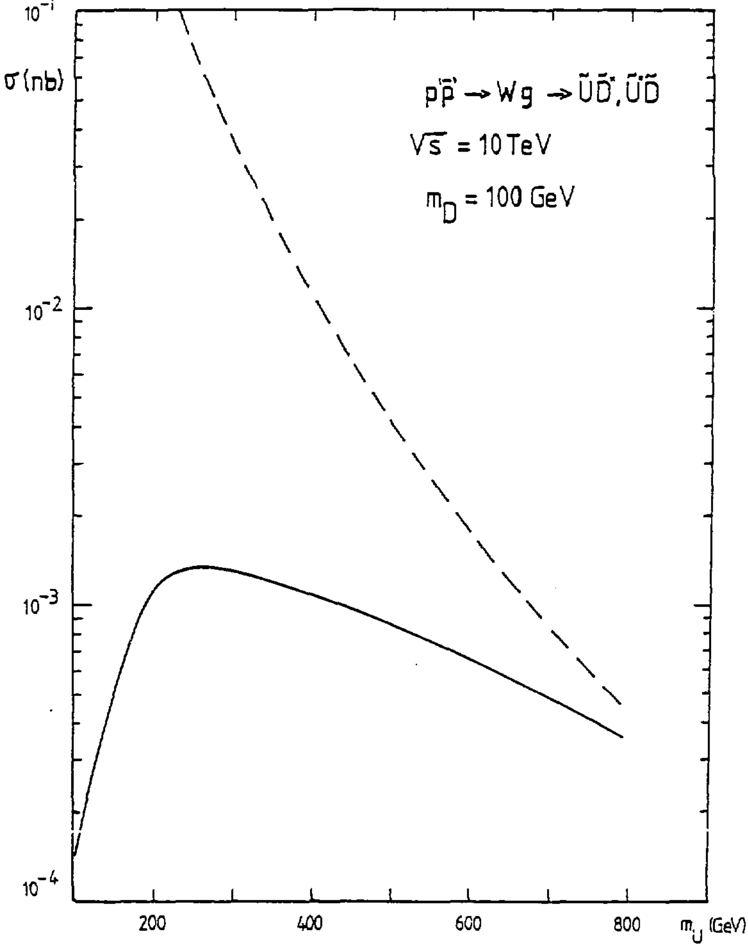


FIG. 1



10^{-1}
σ (nb)
$p\overset{(-)}{p} \to Wg \to U\bar{D}, \bar{U}D$
$\sqrt{s}$ = 10 TeV
m_D = 100 GeV
10^{-2}
10^{-3}
10^{-4}
200
400
600
800
m_U (GeV)

FIG. 2



10^{-1}
σ (nb)
$p\overset{(-)}{p} \to Wg \to \tilde{U}\tilde{D}^{*}, \tilde{U}^{*}\tilde{D}$
$\sqrt{s} = 10$ TeV
$m_D = 100$ GeV
10^{-2}
10^{-3}
10^{-4}
200
400
600
800
m_U (GeV)

FIG. 3