Strong supersymmetric quantum effects on top quark production at the Fermilab Tevatron (Postprint)

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Full Text

Preamble

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Strong Supersymmetric Quantum Effects on Top Quark Production at the Fermilab Tevatron

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Abstract

The supersymmetric QCD corrections to top quark pair production by q$\bar{q}$ annihilation in p$\bar{p}$ collisions are calculated in the minimal supersymmetric model. We consider effects of the mixing of the scalar top quarks on the corrections to the total t$\bar{t}$ production cross section at the Fermilab Tevatron. We found that such correction is less sensitive to squark mass and gluino mass than in the no-mixing case, and in both cases the corrections can exceed 10% even if we consider the recent CDF limit on squark and gluino masses.

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The top quark has been discovered by the CDF (D0) collaboration at the Tevatron [?]. The mass and production cross section are found to be, respectively, 176 ± 8(stat) ± 10(syst) GeV (199^{+19}_{-21}(stat) ± 22(syst) GeV) and $6.8^{+3.6}_{-2.4}$ pb (6.4 ± 2.2 pb) by the CDF (D0) collaboration. At the Tevatron, the dominant production mechanism for a heavy top quark is the QCD process $t\bar{t}$ [?]. Recently, there has been a lot of interest in the one-loop radiative corrections to the top quark production cross section at the Tevatron, which arise from new physics beyond the Standard Model such as the two-Higgs-doublet model (2HDM) and the minimal supersymmetric model (MSSM) [?, ?].

In Ref.~[?], the supersymmetric QCD corrections to top quark production in $p\bar{p}$ collisions were calculated in the simplest case of unmixed squarks and degenerate masses. However, we need to know the impact from mixing squarks because of phenomenological interest, especially the sensitive dependence of the supersymmetric QCD corrections on the squark masses and gluino mass [?]. In that reference, only two cases of the gluino mass $m_{\tilde{g}} = 150$ GeV and $m_{\tilde{g}} = 200$ GeV were considered. The purpose of the present letter is to evaluate the supersymmetric QCD corrections to top quark production at the Tevatron in the general case of mixing squark masses and compare our results with those in the case of unmixed squark masses given in Ref.~[?]. We also discuss further the dependence of such corrections on the gluino mass.

In the MSSM, the strong supersymmetric interaction Lagrangian relevant to our calculation is given, in the presence of squark mixing, by [?]

$$\mathcal{L}_{\tilde{g}\tilde{q}q} = i\sqrt{2} g_s T^a (\bar{q} P_R \tilde{q}_L - \bar{q} P_L \tilde{q}_R) \tilde{g}^a + \text{H.C.},$$

where g_s is the strong coupling constant, T^a are $SU(3)_c$ generators, $P_{L,R} = \frac{1}{2}(1 \mp \gamma_5)$, $\tilde{g}^a$ are the Majorana gluino fields, and $\tilde{q}_{L,R}$ are the current eigenstate squarks, which are related to the corresponding mass eigenstates $\tilde{q}_{1,2}$ by

$$\tilde{q}_1 = \tilde{q}_L \cos\theta + \tilde{q}_R \sin\theta, \quad \tilde{q}_2 = -\tilde{q}_L \sin\theta + \tilde{q}_R \cos\theta.$$

The mixing angle θ as well as the masses $m_{\tilde{q}_{1,2}}$ of the physical squarks are determined by the following mass matrices [?]:

$$\begin{pmatrix} m_{\tilde{t}L}^2 + 0.35 D_Z \& m_t(A_t + \mu\cot\beta) & \\ m_t(A_t + \mu\cot\beta) & m_{\tilde{t}R}^2 + 0.16 D_Z \end{pmatrix},$$

where $D_Z = M_Z^2 \cos 2\beta$ with $\tan\beta$ being the ratio of Higgs vacuum expectation values, $m_{\tilde{t}L,\tilde{t}_R}$ are soft breaking masses, A_t is the soft breaking parameter describing the strength of trilinear scalar interactions, and μ is the supersymmetric Higgs mass parameter.

The relevant Feynman diagrams contributing to one-loop supersymmetric QCD corrections to the $q\bar{q} \to t\bar{t}$ amplitude are shown in Fig.~1 of Ref.~[**?**]. In our calculation we follow the same notation and adopt the same regularization and renormalization scheme as used in Ref.~[**?**]. Notice that the off-diagonal elements of the squark mass matrices are proportional to the quark mass. In the case of supersymmetric partners of light quarks, mixing between the current eigenstates can therefore be neglected. So we will take into account only the mixing between the squarks for the supersymmetric QCD corrections to the $gt\bar{t}$ coupling and neglect such mixing for the $gq\bar{q}$ couplings because these couplings mainly involve light quarks.

The renormalized amplitude for $q\bar{q} \to t\bar{t}$ can be written as

$$\mathcal{M}_{\text{ren}} = \mathcal{M}_0 + \delta\mathcal{M}_{\text{vertex}} + \delta\mathcal{M}_{\text{vertex}} + \delta\mathcal{M}_{\text{box}},$$

where $\mathcal{M}_0$ is the amplitude at tree level and $\delta\mathcal{M}$ represent the supersymmetric QCD corrections arising from the effective $gq\bar{q}$ ($gt\bar{t}$) vertex and box diagrams, which are given by

$$\mathcal{M}_0 = \bar{v}(p_2)(-ig_s T^a \gamma^\nu) u(p1) \frac{-ig_{\nu\mu}}{\hat{s}} \bar{u}(p_3)(-ig_s T^a \gamma^\mu) v(p_4),$$

$$\delta\mathcal{M}_{\text{vertex}} = \bar{v}(p_2)(-ig_s T^a \gamma^\nu) u(p_1) \frac{-ig_{\nu\mu}}{\hat{s}} \bar{u}(p_3)(-ig_s T^a \Gamma^\mu) v(p_4),$$

$$\delta\mathcal{M}_{\text{vertex}} = \bar{v}(p_2)(-ig_s T^a \gamma a'^{\nu}) u(p_1) \frac{-ig_{\nu\mu}}{\hat{s}} \bar{u}(p_3)(-ig_s T^a \gamma^\mu) v(p_4),$$

$$\delta\mathcal{M}_{\text{box}} = i\frac{7\alpha_s g_s^2}{48\pi} \left\{ \bar{u}(p_3) P_R u(p_1) \bar{v}(p_2) P_R v(p_4) f_1 + \bar{u}(p_3) P_R u(p_1) \bar{v}(p_2) P_L v(p_4) f_2 + \bar{u}(p_3) P_L u(p_1) \bar{v}(p_2) P_R v(p_4) f_\right.$$

where

$$\gamma a'^{\nu} = \gamma^\nu F_1 + \gamma^\nu \gamma_5 F_2 + k^\nu F_3 + k^\nu \gamma_5 F_4 + ik_\mu \sigma^{\nu\mu} F_5 + ik_\mu \sigma^{\nu\mu} \gamma_5 F_6,$$

$$\Gamma^\mu = \gamma^\mu F_1' + \gamma^\mu \gamma_5 F_2' + k^\mu F_3' + k^\mu \gamma_5 F_4' + ik_\nu \sigma^{\mu\nu} F_5' + ik_\nu \sigma^{\mu\nu} \gamma_5 F_6'.$$

Here, $k = p_1 + p_2$, p_1, p_2 denote the momenta of incoming partons, and p_3, p_4 are used for the outgoing top quark and its antiparticle. F_i, F_i' and f_i are form factors which are presented in the Appendix.

The renormalized differential cross section of the subprocess is given by

$$\frac{d\hat{\sigma}}{d\cos\theta} = \frac{1}{32\pi\hat{s}} \overline{|\mathcal{M}_{\text{ren}}|^2},$$

where

$$\overline{|\mathcal{M}_{\text{ren}}|^2} = \overline{|\mathcal{M}_0|^2} + 2\text{Re}\, \overline{\mathcal{M}_0^\dagger (\delta\mathcal{M}_{V1} + \delta\mathcal{M}_{V2} + \delta\mathcal{M}_b)},$$

$$\overline{|\mathcal{M}_0|^2} = \frac{8\pi^2\alpha_s^2}{9\hat{s}^2}[2\hat{s}m_t^2 + (\hat{t} - m_t^2)^2 + (\hat{u} - m_t^2)^2],$$

$$\delta\mathcal{M}_{V1} = \frac{8\pi^2\alpha_s^2}{9\hat{s}^2}\left\{[2\hat{s}m_t^2 + (\hat{t} - m_t^2)^2 + (\hat{u} - m_t^2)^2](F_1' - 1) + 2m_t^2\hat{s}^2F_5'\right\},$$

$$\delta\mathcal{M}_{V2} = \frac{8\pi^2\alpha_s^2}{9\hat{s}^2}\left\{[2\hat{s}m_t^2 + (\hat{t} - m_t^2)^2 + (\hat{u} - m_t^2)^2](F_1 - 1) + 2m_t^2\hat{s}^2F_5\right\},$$

$$\delta\mathcal{M}_b = \frac{8\pi^2\alpha_s^2}{216\hat{s}}\left\{m_t[\hat{s}^2 + \hat{s}\hat{t} - \hat{s}m_t^2 + (\hat{s} + \hat{t} - m_t^2)^2](f_6 + f_7 - f_{10} - f_{11}) + m_t[(\hat{s} + \hat{t} - m_t^2)(\hat{s} + \hat{u} - m_t^2) + 2m_t^2\hat{s}\right.$$

In the above equations, θ is the scattering angle between the quark and the top quark, $\hat{s}, \hat{t}$, and $\hat{u}$ are the kinematic invariants for the $2 \to 2$ subprocess with $\hat{s} + \hat{t} + \hat{u} = 2m_t^2$ and $\beta_t = \sqrt{1 - 4m_t^2/\hat{s}}$ is the velocity of the final quarks.

The total cross section for the production of top quark pairs can be written in the form

$$\sigma(s) = \sum_{i,j}\int dx_1 dx_2\, \hat{\sigma}_{ij}(x_1x_2s, m_t^2, \mu^2)\left[F_i^A(x_1,\mu)F_j^B(x_2,\mu) + (A \leftrightarrow B)\right],$$

where $s = (P_1 + P_2)^2$, $p_1 = x_1P_1$, $\hat{s} = x_1x_2s$, $p_2 = x_2P_2$, A and B denote the incident hadrons, P_1 and P_2 are their four-momenta, i, j are the initial partons, x_1 and x_2 are their longitudinal momentum fractions, and the functions $F_j^{A,B}$ are the parton distributions of the initial-state hadrons A and B. In our numerical calculations, we have used the MRS Set A' parton distribution functions [**?**], and do not consider SUSY corrections to the parton distribution functions since the principle of decoupling demands that these corrections are negligible. Introducing a convenient variable $\tau = x_1x_2$ and changing to x_1 and τ as independent variables, the total cross section expression becomes

$$\sigma(s) = \sum_{i,j}\int_{\tau_0}^1 d\tau\, \hat{\sigma}_{ij}(\tau s, m_t^2, \mu^2)\frac{d\mathcal{L}_{ij}}{d\tau},$$

where $\tau_0 = 4m_t^2/s$ and the quantity $d\mathcal{L}_{ij}/d\tau$ is the parton luminosity, which is defined as

$$\frac{d\mathcal{L}_{ij}}{d\tau} = \int_\tau^1 \frac{dx_1}{x_1}\left[F_i^A(x_1,\mu)F_j^B(\tau/x_1,\mu) + (A \leftrightarrow B)\right].$$

Now we present some numerical examples. In our numerical calculation, we input $m_t = 176$ GeV and 1-loop $\alpha_s(Q^2 = \hat{s})$ and use the phase space cuts: $|\eta| < 2.5$, $p_T > 20$ GeV. As for the supersymmetric parameters involved in our calculations, in general, once $\tan\beta$ and $m_{\tilde{t}L}$ are fixed, we are free to choose two independent parameters in the stop mass matrix: $(m_{\tilde{t}R}, A_t + \mu\cot\beta)$, which can also be transferred to $(m_{\tilde{t}R}, m_{\tilde{t}1})$. As explained above, we neglect the mixing of squark masses in the corrections to the $gq\bar{q}$ couplings and also neglect the mass splittings between squarks of different flavor for simplicity.

In Figs.~1-3 we give some numerical results for a simple case, in which we set $\tan\beta = 1$ and $m_{\tilde{t}L} = m_{\tilde{t}R} = m_{\tilde{t}1} = m_{\tilde{q}}$ and thus the mixing angle is $\pi/4$ and $A_t + \mu = m_t$. To compare the results in the no-mixing case with those in the mixing case, we show the results in both cases in Figs.~1-3. From these figures we can see that the corrections in the mixing case are smaller than in the no-mixing case. The plots versus squark and gluino masses in the mixing case become smoother than in the no-mixing case. Fig.~1 shows the dependence of the corrections on gluino mass for a fixed squark mass of 150 GeV. The corrections can be either positive or negative, depending on the gluino mass. The corrections are negative for gluino masses below 150 GeV and become positive above that. When the gluino mass is changed from 100 GeV to 200 GeV, the corrections vary from -2% (-10%) to 16% (23%) in the mixing (no-mixing) case. The corrections reach their positive maximum size at a gluino mass of 200 GeV. When the gluino mass is larger than 200 GeV, the corrections in both cases drop monotonically with the increase of gluino mass and tend to zero at 600 GeV, showing the decoupling effects. The recent CDF lower limit [**?**] on gluino mass is 160 GeV for arbitrary squark mass and 220 GeV when gluino mass equals squark mass. Therefore, the corrections are always positive if we consider the CDF limit on gluino mass.

Figs.~2 and 3 show the dependence of the corrections on squark mass for fixed gluino masses of 150 GeV and 200 GeV, respectively. The corrections in the mixing case differ significantly from those in the no-mixing case for low squark mass. For a gluino mass of 200 GeV, the CDF lower limit [**?**] for squark mass is about 220 GeV. At this lower limit, the corrections are 14% and 18% for the mixing and no-mixing cases, respectively. However, with the increase of squark mass, the difference between the corrections in both cases becomes negligibly small.

From Figs.~1-3 we have found that the corrections vary rapidly with gluino mass, especially for 150 GeV $< m_{\tilde{g}} <$ 200 GeV, though the mixed cases are smaller than the unmixed ones. This is due to the fact that we have set $m_t = 176$ GeV in the numerical calculation, and the threshold for open top quark production is crossed in that region. If we change the top quark mass, we find that such region is also shifted correspondingly, which provides a check on our calculation, especially on the treatment of the threshold.

We also perform numerical calculations for $\tan\beta = 10$. We find that the corrections are not sensitive to the $\tan\beta$ value. For example, with a gluino mass of 200 GeV and squark mass of 150 GeV, we get $\Delta\sigma = 0.482$ pb for $\tan\beta = 1$ and $\Delta\sigma = 0.490$ pb for $\tan\beta = 10$.

In conclusion, we have shown that the supersymmetric QCD corrections to top quark pair production by $q\bar{q}$ annihilation in $p\bar{p}$ collisions can exceed 10% in both the mixing and no-mixing cases of squark masses, even when we consider the recent CDF limit on squark and gluino masses.

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Appendix

The form factor F_1 is given by

$$F_1 = 1 + \sum_{i=1}^{2} \left[F_1^{(1)}(\tilde{t}_i \tilde{g} \tilde{t}_i) + 9F_1^{(2)}(\tilde{t}_i \tilde{g} \tilde{t}_i) \right],$$

where

$$F_1^{(1)} = 2m_t[m_{\tilde{g}} B_{ii}(c_0 + c_{11}) - m_{tA_{ii}}(c_{11} + c_{21})](p_3, k, m_{\tilde{g}}, m_{\tilde{t}i}, m_{\tilde{t}i}),$$

$$F_1^{(2)} = A_{ii}[2\bar{c}_{24} + \hat{s}(c_{22} - c_0) - \hat{t}(c_0 + 2c_{11} + c_{21})](p_3, k, m_{\tilde{t}i}, m_{\tilde{g}}, m_{\tilde{g}}) + 2B_{ii} m_t m_{\tilde{g}}(c_0 + c_{11})(p_3, k, m_{\tilde{t}i}, m_{\tilde{g}}, m_{\tilde{g}}),$$

with

$$A_{ii} = a_i^2 + b_i^2, \quad B_{ii} = a_i^2 - b_i^2,$$

$$a_1 = \cos\theta, \quad b_1 = \sin\theta, \quad a_2 = -\sin\theta, \quad b_2 = \cos\theta,$$

$$F(\tilde{t}_i \tilde{g} \tilde{t}_i) = \int_0^1 dy\, y \ln \frac{m_{\tilde{t}i}^2 y(1-y) + m_{\tilde{g}}^2 y}{m_{\tilde{t}i}^2 y(1-y) + m_{\tilde{g}}^2 y + m_t^2},$$

$$G(\tilde{t}_i \tilde{g} \tilde{t}_i) = \int_0^1 dy\, y^{n+1}(1-y) \ln \frac{m_{\tilde{t}i}^2 y(1-y) + m_{\tilde{g}}^2 y}{m_{\tilde{t}i}^2 y(1-y) + m_{\tilde{g}}^2 y + m_t^2},$$

and $\bar{c}_{24} = c_{24} - \Delta/2$ with $\Delta = 2/(4-d) - \gamma_E + \ln 4\pi$.

The form factor F_1' is obtained from F_1 by the replacement

$$A_{ii} \to K_{jj} = A_{jj}|_{\theta \to \theta'}, \quad B_{ii} \to L_{jj} = B_{jj}|_{\theta \to \theta'},$$

where

$$F_1^{(1)'} = 2K_{jj} \bar{c}_{24}(p_1, k, m_{\tilde{g}}, m_{\tilde{q}j}, m_{\tilde{q}j}),$$

$$F_1^{(2)'} = K_{jj}[2\bar{c}_{24} + \hat{s}(c_{22} - c_0)](p_1, k, m_{\tilde{q}j}, m_{\tilde{g}}, m_{\tilde{g}}).$$

The form factor F_5 is given by

$$F_5 = \sum_{i=1}^{2} \left[F_5^{(1)}(\tilde{t}_i \tilde{g} \tilde{t}_i) + 9F_5^{(2)}(\tilde{t}_i \tilde{g} \tilde{t}_i) \right],$$

where

$$F_5^{(1)} = [m_t A_{ii}(c_{11} + c_{21}) - m_{\tilde{g}} B_{ii}(c_0 + c_{11})](p_3, k, m_{\tilde{g}}, m_{\tilde{t}i}, m_{\tilde{t}i}),$$

$$F_5^{(2)} = [m_t A_{ii}(c_{11} + c_{21}) + m_{\tilde{g}} B_{ii} c_{11}](p_3, k, m_{\tilde{t}i}, m_{\tilde{g}}, m_{\tilde{g}}).$$

The f_i are given by

$$f_2 = m_{\tilde{g}}^2 \sigma_{ij}'^2 D_0 - m_t m_{\tilde{g}} \sigma_{ij} \sigma_{ij}'(D_{12} + D_{13} - D_{11}) + m_t^2 \sigma_{ij}^2(D_{26} - D_{12} - D_{24}),$$

$$f_6 = m_{\tilde{g}}\sigma_{ij}\sigma'_{ij}(D_0 + D_{11}) + m_t\sigma_{ij}^2(D_{12} + D_{24}),$$
$$f_{10} = m_{\tilde{g}}\sigma_{ij}\sigma'_{ij}(D_{13} - D_{11}) + m_t\sigma_{ij}^2(D_{12} + D_{23} + D_{24} - D_{13} - D_{26} - D_{25}),$$
$$f_{14} = \sigma_{ij}^2(D_{12} + D_{24} - D_{13} - D_{25}),$$
$$f_{18} = \sigma_{ij}^2 D_{27},$$

where

$$(f_3, f_7, f_{11}, f_{15}, f_{19}) = (f_2, f_6, f_{10}, f_{14}, f_{18})|_{\sigma_{ij}\to\lambda_{ij}},$$

with

$$\sigma_{ij} = (a_i + b_i)(c_j + d_j), \quad \sigma'_{ij} = (a_i - b_i)(c_j + d_j),$$
$$\lambda_{ij} = (a_i - b_i)(c_j - d_j), \quad \lambda'_{ij} = (a_i + b_i)(c_j - d_j),$$

and $(c_j, d_j) = (a_j, b_j)|_{\theta\to\theta'}$. Here $D_0, D_{ij}(p_2, p_4, m_{\tilde{g}}, m_{\tilde{q}j}, m_{\tilde{g}}, m_{\tilde{t}i})$ are the 4-point Feynman integrals [**?**].

Note: Figure translations are in progress. See original paper for figures.

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