Top decays into lighter stop and gluino

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ABSTRACT

We have calculated the decay width of the process $t \rightarrow \tilde{g}\tilde{t}_1$ including one loop QCD corrections. We found that the decay width of the such process could be comparable with that of the standard channel $t \rightarrow bW^+$, and, its QCD correction could enhance the widths over 30% in a very large mass range of the lighter stop.

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I. INTRODUCTION

Top quark play an important role in the search for new physics due to its huge mass compared to other fermions in the standard model (SM). Besides its large Yukawa coupling to Higgs sector, top quark also anticipates strong interaction (QCD) with its supersymmetric partner through the top-stop-gluino vertex in a particle model of supersymmetry (SUSY). As a supersymmetric counterpart, a large mass splitting can emerge between the mass-eigenstates of the stops and could lead to a lighter physical stop, so that stop quark may be produced in our top quark factories.

The experimental lower mass limit of stop is 86 GeV [1] and the theoretical one could be more smaller [2]. In most of conventional proposals stop will be produced in pair, then they need larger energy input in general. Comparing to stop pair production, one wonder whether stop could be produced associated with a light gluino in the decay of top quark. Such a channel can be accommodated in the single top events, and the top quark sector would become a nice place to study SUSY-QCD.

The key point lies in that, whether the gaugino of $SU(3)_c$ is also sufficiently light. We would like to make an analyses on the situation of light gluino bellow.

Historically, there had ever glimmered light gluinos in many scenarios with natural assumptions about the mechanics of SUSY breaking at higher energy scale. For example, the mass of gluino was resulted in $1 \sim 100$ GeV typically if the SUSY breaking was transmitted to the observable sector by gravity (mSUDRA); and the masses of gluinos and squarks will arise smaller ( bellow GeV ) as a loop correction (next leading order effect) when that breaking is gauge-mediated ( GMSB ) [3]. Those light gluino scenario can explain several well-established phenomena, such as the anomalously slow running of $SU(3)_c$ coupling, anomalies in jet production [4] as well as in the $\eta(1410)$ [5].

The difficulty is obviously that, present colliders do not favor light gluino. A lower mass bound 180 GeV has ever been published [6]. Nevertheless, a controversial assumption that there is no gluino lighter than 5 GeV as well as the state-of-art usage of perturbative theory [7], had made the constraint less convincing, especially leave a margin for a gluino light than 5 GeV. The concomitant predictions of a light Higgs mass and a light chargino, have also been (at least marginally) ruled out at LEP II [8], so has been the light gluino in mSUDRA.
models. However, light gluino in GMSB models can survive in such experiment. Anyway all the direct searches for a light gluino have turned to a negative results [9] since 1998.

Then light gluino is cornered into bound states (for example $R^0$ ) to live if it is light enough. For a light gluino in a bound state, the well known KTeV experiment is an rather important one to give a constraint on its mass. KTeV result indeed has closed most of the place for a light gluino [10] and has given an impression that light gluino had been ruled out. Although, there are still two shreds that are not excluded by KTeV data in fact. The gluino could weight as $0 \sim 1.6GeV$ or as $5GeV$ [10]. Such light gluino can survive through another important E761 experiment too [11].

Although recent re-analyses [12] showed a consistency with SM and really eroded the attractiveness of light gluino in explaining slow running of $SU(3)_c$ coupling, such re-analyses relied too much on the measurement of low energy QCD, for example $\alpha_s(M_\tau)$, which itself is difficult enough.

Through above investigation, we have make it clear that, there is no counter-indication which can conclusively convince one to give up a light gluino ( few GeV or so). Recently, more new suggestion on light gluino within the framework of GMSB have been proposed [13]. Both scenario allow a gluino at few GeV and predicted the standard missing energy signature of SUSY may become not suitable.

Then from the narrow window for light gluino, a noticeable partial width to stop and gluino might be detected in top quark decay when the statistics of the top events are improved.

In the present paper, we will consider the decay process $t \rightarrow \tilde{g}\tilde{t}_1$ including the QCD corrections, which possibility has been mentioned by Raby et. al [13]. In experiments, the stop may be followed by the decays $\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$ or $\tilde{t}_1 \rightarrow b\tilde{\chi}_1^+$ [16], then its typical signal may be one hadronic jet plusing a large missing $P_T$. Instead of the detailed analysing the signal against possible background, we emphasize the possible anomalous branch ratio of top decay. In fact, in the SM $BR(t \rightarrow bW^+)$ had been assumed as the dominant one for the observed top pair events at Tevatron [14]. Thus, the non-SM branching ratio of the top quark which could be as large as the SM one is not excluded, and the searching for non-SM top decay are undergoing in both CDF and $D\bar{O}$ [15].
In the literatures, there have been several works [17] discussing the decays of stop or gluino into top quark, which have the same dynamics (interaction vertex) as the process discussed here, and these calculations offered a good references to check our calculations.

II. CALCULATIONS

The tree level diagram of the process $t(p) \to \tilde{g}(k_1)\tilde{t}_1(k_2)$ is shown in Fig. 1 (a), and the decay width is

$$\Gamma_0 = \frac{N\alpha_s}{3m_t^2}(m_t^2 - m_{\tilde{g}}^2 + m_{\tilde{t}_1}^2 - 2m_{\tilde{g}}m_{\tilde{t}_1}\sin(2\theta)), \quad (1)$$

where

$$N = \sqrt{(m_t^2 - (m_{\tilde{g}} + m_{\tilde{t}_1})^2)(m_t^2 - (m_{\tilde{g}} - m_{\tilde{t}_1})^2)}, \quad (2)$$

and $\theta$ is the stop mixing angle which is to be defined in next section.

A. Virtual corrections

The $O(\alpha_s)$ virtual corrections arise from triangle as well as the gluino, quark and squark self energies diagrams as shown in Fig. 1 (b $\sim$ e) respectively. Through our calculations, we employed dimensional reduction regularization (DRR) to control all the ultraviolet divergences, which is widely adopted in the calculations of the radiative corrections in the minimal supersymmetric standard model (MSSM) since the conventional dimensional regularization violates supersymmetry with a mismatch between the numbers of degree of freedom of gauge bosons and gauginos [18]. At the same time, we adopted the on-mass-shell renormalization scheme [19,20].

The one loop decay width can be expressed as

$$\Gamma_1 = \frac{4\sqrt{2}\pi N\alpha_s^2}{3m_t^2} Re \left\{ f_1((m_{\tilde{g}} + m_t)^2 - m_{\tilde{t}_1}^2)(\cos \theta - \sin \theta) \\ - f_2((m_{\tilde{g}} - m_t)^2 - m_{\tilde{t}_1}^2)(\sin \theta + \cos \theta) \right\}, \quad (3)$$

where
where \( f_i^c \) are the contributions in the form of counterterms,

\[
f_1 = f_1^c + f_1^d, \quad f_2 = f_2^c + f_2^d,
\]

The renormalization constants in Eq. 5 are,

\[
f_1^c = \sqrt{2}(f_L \cos \theta - f_R \sin \theta),
\]

\[
f_2^c = \sqrt{2}(-f_L \cos \theta - f_R \sin \theta),
\]

\[
f_L = \frac{1}{2} Z_L^1 + \frac{1}{2} Z_R^9 + \delta g_s / g_s - \tan \theta \delta \theta + \frac{1}{2} Z_{t_1},
\]

\[
f_R = \frac{1}{2} Z_R^1 + \frac{1}{2} Z_L^9 + \delta g_s / g_s + \cot \theta \delta \theta + \frac{1}{2} Z_{t_1}.
\]
We found the on mass shell definition for the counterterm of stops mixing,

\[
\delta \theta = \frac{1}{2} \left( \delta Z^{12} + \delta Z^{21} \right) = \frac{1}{2} \left( \sum_{i=1}^{3} \frac{\tilde{t}_i}{m_{\tilde{t}_i}^2 - m_{\tilde{t}_i}^2} \right)
\]

\[
= \frac{\mu_1 \mu_2 \cos(2\theta)}{6\pi^2 (m_{\tilde{t}_2}^2 - m_{\tilde{t}_1}^2)} \left( (B_0(m_{\tilde{t}_2}^2, m_{\tilde{t}_1}^2, m_{\tilde{t}_2}) - B_0(m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2, m_{\tilde{t}_2})) \right) \tag{7}
\]
is appropriate for our calculation, which has been adopted in [21]. One can exam that such a $\delta \theta$ is UV convergent, so that it is renormalization group ( scale ) independent at one loop level. As to the counterterm of the Yukawa coupling $g_s^y$ of $SU(3)_c$, we prefer to the treatment in [17],

\[
\delta g_s^y = -\frac{\alpha_s(\mu_R)}{4\pi} \left\{ [\Delta - \log(\frac{\mu_R^2}{\mu^2})] \beta_0^{\text{susy}} \frac{\beta_0^{\text{susy}}}{2} - \left( \frac{2}{3} \frac{\beta_0^{\text{susy}}}{2} \right) \right\}. \tag{8}
\]

However there are fewer virtual particles decoupled in our scenario with light gluino and light squarks,

\[
\Delta g_s^y = \frac{\alpha_s(\mu_R)}{4\pi} \left[ \frac{1}{2} \frac{\log(\frac{\mu_R^2}{\mu^2})}{m_t^2} + \frac{1}{12} \frac{\log(\frac{\mu_R^2}{m_t^2})}{m_t^2} + \frac{1}{12} \frac{\log(\frac{\mu_R^2}{m_t^2})}{m_t^2} + \frac{1}{6} \frac{\log(\frac{\mu_R^2}{m_t^2})}{m_t^2} \right] \tag{9}
\]

The $f_d^i$ in Eq. 4 are the contributions from the 3-points diagrams,

\[
f_d^1 = \frac{\cos \theta - \sin \theta}{96\sqrt{\pi}^2} \left( (9 - \sin(2\theta)) B_0(m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2, m_{\tilde{t}_2}^2) + \sin(2\theta) B_0(m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2, m_{\tilde{t}_1}^2) \right)
\]

\[
+35 B_0(m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2, m_{\tilde{t}_2}^2) + 18 m_\tilde{g} (3 m_\tilde{g} + m_t) C_0^{(1)} - 18 (m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2) C_0^{(2)}
\]

\[
+(-m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2) C_0^{(3)} - 2 m_{\tilde{t}_1}^2 \sin(2\theta) C_0^{(4)} + 2 m_{\tilde{t}_1} (m_{\tilde{g}} + m_t + 2 m_{\tilde{g}} + m_t \sin(2\theta)) C_0^{(5)}
\]

\[
+18(-2k_{1,p} + 3m_{\tilde{g}}^2 + m_{\tilde{g}} m_t) C_1^{(1)} + 18(-2k_{1,p} + 2m_{\tilde{g}}^2 + m_{\tilde{g}} m_t) C_1^{(2)} - 2 m_{\tilde{g}} m_t C_1^{(3)}
\]

\[
-2 m_{\tilde{g}} (m_{\tilde{g}} + (m_{\tilde{g}} - m_t) \sin(2\theta)) C_1^{(4)} - 2 m_{\tilde{g}} (m_t + (m_t - m_{\tilde{g}}) \sin(2\theta)) C_1^{(5)}
\]

\[
+18(-4k_{1,p} + 3m_{\tilde{g}}^2 + m_{\tilde{g}} m_t) C_2^{(1)} + 18(-3k_{1,p} + 2m_{\tilde{g}}^2 + m_{\tilde{g}} m_t) C_2^{(2)} - 2(k_{1,p} - m_{\tilde{g}}) C_2^{(3)}
\]

\[
-2(m_{\tilde{g}}^2 - m_{\tilde{g}} m_t + (-k_{1,p} + m_{\tilde{g}}^2 - m_{\tilde{g}} m_t + m_t) \sin(2\theta)) C_2^{(4)}
\]

\[
-2(-m_{\tilde{g}}^2 + m_{\tilde{g}} m_t + (k_{1,p} - m_{\tilde{g}}^2 + m_{\tilde{g}} m_t - m_t^2) \sin(2\theta)) C_2^{(5)}
\]

\[
f_d^2 = \frac{\cos \theta + \sin \theta}{96\sqrt{\pi}^2} \left( -(9 + \sin(2\theta)) B_0(m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2, m_{\tilde{t}_1}^2) + \sin(2\theta) B_0(m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2, m_{\tilde{t}_2}^2) \right)
\]

\[
-35 B_0(m_{\tilde{t}_2}^2, m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2) + 18 m_\tilde{g} (-3 m_\tilde{g} + m_t) C_0^{(1)} + 18 (m_{\tilde{t}_2}^2 - m_{\tilde{t}_1}^2) C_0^{(2)}
\]

\[
-(-m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2) C_0^{(3)} - 2 m_{\tilde{t}_1}^2 \sin(2\theta) C_0^{(4)} + 2 m_{\tilde{t}_1} (m_{\tilde{g}} + m_t + m_t \sin(2\theta)) C_0^{(5)}
\]

\[
+18(2k_{1,p} - 3m_{\tilde{g}}^2 + m_{\tilde{g}} m_t) C_1^{(1)} + 18(2k_{1,p} - 2m_{\tilde{g}}^2 + m_{\tilde{g}} m_t) C_1^{(2)} - 2 m_{\tilde{g}} m_t C_1^{(3)}
\]
The corresponding matrix element as given by the Feynman diagrams (Fig. 1 (f)) is

\[
+2m_g(m_g - (m_\bar{g} + m_\ell) \sin(2\theta))C_1^{(4)} + 2m_\bar{g}(-m_\ell + (m_\ell + m_g) \sin(2\theta))C_1^{(5)} \\
-18(-4k_1.p + 3m_g^2 + m_\ell^2)C_2^{(1)} - 18(-3k_1.p + 2m_g^2 + m_\ell^2)C_2^{(2)} + 2(k_1.p - m_\ell^2)C_2^{(3)} \\
+2(m_g^2 + m_\ell.m_\ell + (k_1.p - m_g^2 - m_\ell.m_\ell - m_\ell^2) \sin(2\theta))C_2^{(4)} \\
+2(-m_\ell^2 - m_\ell.m_\ell + (-k_1.p + m_g^2 + m_\ell.m_\ell + m_\ell^2) \sin(2\theta))C_2^{(5)}.
\]

(10)

In the presentation of the formulas above, we have used notations,

\[
\Delta = \frac{2}{4-D} - \gamma_E + \ln(4\pi), \\
\beta_{0}^{\text{e.m.}} = 3N_c - n_f = \beta_{0}^{\text{light}} + \beta_{0}^{\text{heavy}} \\
k_1.p = \frac{m_\ell^2 - m_\ell^2 + m_\bar{g}^2}{2}, \\
\dot{B}(p^2, m_a^2, m_b^2) = \partial B_0(p^2, m_a^2, m_b^2)/\partial p^2,
\]

\[
(11)
\]

\beta_{\text{heavy}}^{\text{specific}} is for the decoupled virtual particles, which will not contribute to the scale evolution of \(\alpha_s\). The definitions of the scalar integrals \(B_s\) and \(C_s\) could be found in Ref. [19,22], and the indexes \((1) - (5)\) of \(C\) functions have the variables

\[
(m_g^2, m_\ell^2, m_\bar{g}^2, \lambda^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, \lambda^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, \lambda^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, \lambda^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, \lambda^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, m_\ell^2), (m_g^2, m_\ell^2, m_\bar{g}^2, m_\ell^2), \]

respectively.

**B. Real corrections**

The infrared divergence arise from the virtual massless gluon corrections are compensated by the real gluon bremsstrahlung corrections, i.e. the three-body decay

\[
t(p) \rightarrow \bar{g}(k_1)\bar{\ell}(k_2)g(k).
\]

(12)

The corresponding matrix element as given by the Feynman diagrams (Fig. 1 (f) ) is

\[
M_b = 4\sqrt{2}\pi\alpha_s U(k1)\{-i f_{abc} T_\epsilon \frac{1}{2k.k_1} f(k + k_1 + m_g)(\sin \theta P_R - \cos \theta P_L) \\
+ T_b T_a \frac{1}{-2p.k} (\sin \theta P_R - \cos \theta P_L)(k_1 + k_2 + m_\ell) f \\
+ T_a T_b \frac{1}{2k_2.k} (\sin \theta P_R - \cos \theta P_L)(2k_2 + k) \epsilon U(p),
\]

(13)
where $\epsilon$ denotes the polarization vector of gluon. Squaring the matrix element, performing the polarization and color sum over the square of the amplitude and integrating over the phase space yields the complete bremsstrahlung cross section as

$$\Gamma_b = \frac{\alpha^2}{24\pi m_t} \left\{ 48 \left[ 2m_g^2(m_g^2 - m_{t_i}^2 + m_I^2 - 2m_g m_t \sin(2\theta))I_{11} + 2m_g^2I_1 + I_1^0 \right] - 2m_g m_t \sin(2\theta)I_1 + \frac{64}{3} \left[ 2m_t^2(-m_g^2 + m_{t_i}^2 - m_I^2 + 2m_g m_t \sin(2\theta))I_{00} - 2m_t^2I_0 - I_0^1 + 2m_g m_t \sin(2\theta)I_0 \right] + \frac{64}{3} \left[ 2m_t^2(-m_g^2 + m_{t_i}^2 - m_I^2 + 2m_g m_t \sin(2\theta))I_{02} + 2(m_g^2 + m_{t_i}^2)I_2 + 2(m_g^2 - m_I^2)I_0 + I - 2m_g m_t \sin(2\theta)(I_2 - I_0) \right] + 24 \left[ -2(m_g^2 - m_{t_i}^2 + m_I^2)^2 - 2m_g m_t \sin(2\theta)(m_g^2 - m_{t_i}^2 + m_I^2)I_{01} + 2(m_g^2 - m_I^2)I_0 + 2(m_g^2 - m_{t_i}^2 + m_g m_t \sin(2\theta))(I_1 + I_0) - 2I \right] + 24 \left[ 2(-m_g^2 + (m_{t_i}^2 + m_I^2) - 2m_g m_t \sin(2\theta)(m_g^2 + m_{t_i}^2 - m_I^2))I_{12} - (m_g^2 + m_{t_i}^2 + m_I^2)I_2 + 2(m_g^2 - m_{t_i}^2)I_1 - I \right]. \quad (14)$$

Here the hard gluon has been included and the definition of function $I_{ab...}^{AB...}$ are

$$I_{ab...}^{AB...} = \frac{1}{\pi^2} \int \frac{d^3k_1}{2k_1^0} \frac{d^3k_2}{2k_2^0} \frac{d^3k}{2k^0} \delta^4(p - k_1 - k_2 - k) \frac{\pm k.P_A(\pm k.P_B)...}{(\pm k.P_A)(\pm k.P_B)...}, \quad (15)$$

where the minus corresponds top quark and plus corresponds gluino and stop$^*$. 

### III. NUMERICAL RESULTS

In the MSSM the mass eigenstates $\tilde{q}_1$ and $\tilde{q}_2$ of the squarks are related to the interaction eigenstates $\tilde{q}_L$ and $\tilde{q}_R$ by [23]

$$\begin{pmatrix} \tilde{q}_1 \\ \tilde{q}_2 \end{pmatrix} = R^\theta \begin{pmatrix} \tilde{q}_L \\ \tilde{q}_R \end{pmatrix} \quad \text{with} \quad R^\theta = \begin{pmatrix} \cos \theta_{\tilde{q}} & \sin \theta_{\tilde{q}} \\ -\sin \theta_{\tilde{q}} & \cos \theta_{\tilde{q}} \end{pmatrix}. \quad (16)$$

$^*$It should be noticed that (D.11) and (D.12) in Ref. [19] must be corrected as following: the first term in the square bracket of (D.11) must be replaced by the corresponding term in (D.12), and vice versa.
Following the notation of [23], the mixing angle \( \theta_q \) and the masses \( m_{\tilde{q}_1, 2} \) can be calculated by diagonalizing the following mass matrices

\[
M^2_{\tilde{q}} = \begin{pmatrix}
M^2_{LL} & M^2_{LR} \\
M^2_{RL} & M^2_{RR}
\end{pmatrix},
\]

\[
M^2_{LL} = m^2_{\tilde{Q}} + m^2_{q} + m^2_z \cos 2\beta(I^3_L q - e_q \sin^2 \theta_w),
\]

\[
M^2_{RR} = m^2_{\tilde{U}, \tilde{D}} + m^2_{q} + m^2_z \cos 2\beta e_q \sin^2 \theta_w.
\]

(17)

\( \omega \)From Eqs. 16 and 17, \( m^2_{\tilde{t}_1, 2} \) and \( \theta \) can be derived as

\[
m^2_{\tilde{t}_{1,2}} = \frac{1}{2} \left[ M^2_{LL} + M^2_{RR} \mp \sqrt{(M^2_{LL} - M^2_{RR})^2 + 4M^4_{LR}} \right]
\]

\[
\tan \theta = \frac{m^2_{\tilde{t}_1} - M^2_{LL}}{M^2_{LR}}.
\]

(18)

For the numerical calculations, the renormalization scale is set at the mass of decay particle, \( \mu_R = m_t = 176.0 \) GeV \( \alpha_S = 0.108 \). We also evaluated at \( \mu_R = m_t/2\alpha_s = .12 \) and found the scale dependence very weak. The soft SUSY breaking mass terms are always naturally assumed as \( m_{\tilde{U}} = m_{\tilde{D}} = m_{\tilde{Q}} = m_S \). For simplicity, we also define a variable \( r = M^2_{LL}/M^2_{LR} \).

In Fig. 2, we show the branching ratios and the QCD relative corrections as function of \( m_{\tilde{t}_1} \), assuming \( r = 0.9 \). The definition of branching ratio and relative correction are

\[
Br(t \to \tilde{g}\tilde{t}_1) = \frac{1}{1 + \Gamma(t \to bW^+)/\Gamma(t \to \tilde{g}\tilde{t}_1)} \quad \text{and} \quad \delta = (\Gamma_1 + \Gamma_0)/\Gamma_0,
\]

respectively. \( \omega \)From these curves, one can see that the branching ratios, which are not sensitive to the light gluino mass allowed by experiments, decrease with the increment of \( m_{\tilde{t}_1} \), which is the natural consequence of the shrink of the phase space. One also can see that, the decay width could be large enough to compare with the \( t \to bW^+ \) channel in a large range of parameter space. On the other hand, the QCD corrections always enhanced the decay widths, for \( m_{\tilde{t}_1} = 160 \) GeV, the corrections exceed 40%.

In Fig. 3, we exam the result for its dependence upon the decoupled particles, by alternating the way for the mass of the lighter stop to vary. In these figures, \( r \) was scanned with fixed soft SUSY breaking mass term \( m_S \) as 300 GeV. We can see these results are very close to Fig. 2, which is due to the fact that the contributions from other squarks (relatively heavy) are less important.
In summary, we have calculated the decay width of the process $t \to \tilde{g}\tilde{t}_1$ including the one loop QCD corrections.

From our numerical examples, we can see that the QCD corrections could enhance the decay widths over 30% in a very large mass range of the lighter stop, and the decay widths of the process could be larger than that of the channel $t \to bW^\pm$. As a supplement of stop pair production at LEP 2000 or Tevatron, such a single stop (top) channel might make the top phenomenology more rich if it is allowed kinetically.

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FIG. 1. Feynman diagrams for the process $t ightarrow 	ilde{g} 	ilde{t}_1$. 
FIG. 2. Branching ratios and relative corrections as function of $m_{\tilde{t}_1}$, where $r = 0.9$. (a), the solid (dashed) line represents the decay widths including QCD corrections for $m_{\tilde{g}} = 1.0$ (5.0) GeV, and the dotted (dot-dashed) line denotes the tree level decay widths for $m_{\tilde{g}} = 1.0$ (5.0) GeV; (b), the solid (dashed) line represent the relative corrections for $m_{\tilde{g}} = 1.0$ (5.0) GeV.
FIG. 3. see caption of Fig. 2 but fixed the soft SUSY breaking mass term $m_S$ as 300 GeV.