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Leading Supersymmetric Electroweak Corrections to Top Quark Decay into a Neutralino and Light Stop*

YANG YaSheng, LI ChongSheng, CAO QingHong and LIU HongXuan

Department of Physics, Peking University, Beijing 100871, China

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Abstract We calculate the leading supersymmetric electroweak corrections to $t \rightarrow \tilde{t}_1 \tilde{\chi}_j^0$ using dimensional reduction scheme within an approximation of low $\tan \beta$, which can easily be extended to $t \rightarrow \tilde{\chi}_j^+ \tilde{b}_i$. The numerical results show that such corrections can exceed -12% for $\tan \beta = 2$ and -6% for $\tan \beta = 11$, respectively. And these corrections vary insensitively with $m_{\tilde{t}_1}$ in the region allowed by the kinematics except for ones near the threshold.

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Among supersymmetric decay modes of the top quark, in addition to the top quark decay into possible charged Higgs bosons plus bottom quark, the most interesting decay channel is the decay into a light top squark (\tilde{t}_1) plus a neutralino ($\tilde{\chi}_j^0$), since both \tilde{t}_1 and $\tilde{\chi}_j^0$ can be quite light in the minimal supersymmetric standard model (MSSM) and the decay $t \rightarrow \tilde{t}_1 \tilde{\chi}_j^0$ could occur in a reasonably large volume of the parameter space with a sizable branching ratio.^[1] This exotic decay mode of the top quark has been extensively discussed at the tree level,^[1] and one-loop QCD corrections have also been calculated.^[2,3] Recent investigation^[4] indicated that the current bounds on the decay $t \rightarrow \tilde{t}_1 \tilde{\chi}_1^0$ from the available data at Fermilab Tevatron and CERN LEP are actually quite weak, and this supersymmetric decay mode of the top quark can be either discovered or placed much stronger upper limits (at the level of 1%) on the branching fraction $B(t \rightarrow \tilde{t}_1 \tilde{\chi}_1^0)$ at the Tevatron RunII by searching for the final state $t\bar{t} \rightarrow W b c \tilde{\chi}_1^0 \tilde{\chi}_1^0$. Thus a more accurate calculation of the decay mechanism is also necessary to provide a solid basis for experimental analysis of observing $t \rightarrow \tilde{t}_1 \tilde{\chi}_1^0$ at the Tevatron. In this letter, we present the calculation of the leading supersymmetric electroweak (SUSY-EW) corrections to the top quark decay into the lightest top squark plus a neutralino, which are induced by Yukawa couplings from the top quarks and the bottom quarks. Our results can be straightforwardly generalized to the decay $t \rightarrow \tilde{\chi}_j^+ \tilde{b}_1$, where \tilde{b}_1 is a light bottom squark.

The interactions of top and top squarks with neutralinos are given by the Lagrangian^[5]

$$\mathcal{L}_{t\tilde{t}, \tilde{\chi}_j^0} = -\sqrt{2}\bar{t}(L_{ij}P_L + R_{ij}P_R)\tilde{\chi}_j^0\tilde{t}_i + \text{h.c.}, \quad (1)$$

where

$$\begin{aligned} L_{1j} &= A_j \cos \theta - C_j \sin \theta, & L_{2j} &= -A_j \sin \theta - C_j \cos \theta, \\ R_{1j} &= B_j \cos \theta - D_j \sin \theta, & R_{2j} &= -B_j \sin \theta - D_j \cos \theta \end{aligned} \quad (2)$$

with

$$\begin{aligned} A_j &= D_j^* = \frac{h_t N_{j4}^*}{\sqrt{2}}, & B_j &= C_j^* + \frac{g N'_{j2}}{2C_W}, \\ C_j &= \frac{2}{3} e N'^*_{j1} - \frac{2}{3} \frac{g S_W^2}{C_W} N'^*_{j2} \end{aligned} \quad (3)$$

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and

$$N'_{j1} = N_{j1}C_W + N_{j2}S_W, \quad N'_{j2} = -N_{j1}S_W + N_{j2}C_W. \quad (4)$$

Here $h_t \equiv gm_t/\sqrt{2}m_w \sin \beta$ is the top Yukawa coupling, $S_W \equiv \sin \theta_w$, $C_W \equiv \cos \theta_w$, $P_{L,R} \equiv \frac{1}{2}(1 \mp \gamma_5)$, and N_{ij} are the elements of the 4×4 matrix N defined in Ref. [5], which can be calculated numerically. $T^a = \lambda^a/2$ are the Gell-Mann matrices and θ is the mixing angle between left- and right-handed stops which are related to the mass eigenstates \tilde{t}_i in Eqs (1) and (2) by

$$\begin{pmatrix} \tilde{t}_1 \\ \tilde{t}_2 \end{pmatrix} = \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} \tilde{t}_L \\ \tilde{t}_R \end{pmatrix}. \quad (5)$$

This rotation matrix, equation (5), diagonalizes the stop mass matrix^[5]

$$M_{\tilde{t}}^2 = \begin{pmatrix} M_{\tilde{t}_L}^2 + m_t^2 + 0.35 \cos(2\beta) M_Z^2 & -m_t(A_t + \mu \cot \beta) \\ -m_t(A_t + \mu \cot \beta) & M_{\tilde{t}_R}^2 + m_t^2 + 0.16 \cos(2\beta) M_Z^2 \end{pmatrix}, \quad (6)$$

where $M_{\tilde{t}_L}^2, M_{\tilde{t}_R}^2$ are the soft SUSY-breaking mass terms for left- and right-handed stops, μ is the supersymmetric Higgs mass parameter in the superpotential, A_t is the trilinear soft SUSY-breaking parameter, and $\tan \beta = v_2/v_1$ is the ratio of the vacuum expectation values of the two Higgs doublets.

The tree-level partial decay width of $t \rightarrow \tilde{t}_1 \tilde{\chi}_j^0$ is given by

$$\Gamma_0 = \frac{1}{16\pi m_t^3} \lambda^{1/2} (m_t^2, m_{\tilde{\chi}_j^0}^2, m_{\tilde{t}_1}^2) [(|L_{1j}|^2 + |R_{1j}|^2)(m_t^2 + m_{\tilde{\chi}_j^0}^2 - m_{\tilde{t}_1}^2) + 4 \operatorname{Re} (L_{1j}^* R_{1j}) m_t m_{\tilde{\chi}_j^0}], \quad (7)$$

where $\lambda(x, y, z) = (x - y - z)^2 - 4yz$. The leading SUSY-EW corrections to $t \rightarrow \tilde{t}_1 \tilde{\chi}_j^0$ arise from the diagrams of Fig. 1. In our calculation, we use the dimensional reduction technique,^[6] which preserves supersymmetry, and for regularization of all the ultraviolet divergencies, adopt the on-mass-shell renormalization scheme^[7] in which the coupling constant and the physical masses are chosen to be the renormalized parameters. The finite parts of the counterterms are then fixed by the renormalization conditions that the quark, squark and neutralino propagators have the poles at their physical masses. For the leading SUSY-EW corrections to the decay $t \rightarrow \tilde{t}_1 \tilde{\chi}_j^0$, which we are considering, only the top Yukawa coupling and the stop mixing angle in the bare couplings need to be renormalized, where the parameter $\tan \beta$ involved in the Yukawa coupling has to be renormalized, too. There are mainly two approaches to fixing renormalization constant Z_β of the parameter $\tan \beta$ in the literatures. One of them was presented by Mendez and Pomarol in Ref. [8], where the Z_β is fixed by the requirement that the on-mass-shell $H^+ \ell^- \bar{\nu}_\ell$ coupling keeps the forms of Eq. (3) of Ref. [8] to all orders of perturbation theory. The other one was recently introduced by Coarasa *et al.* in Ref. [9], where they chose as a physical observable the decay width of the charged Higgs boson into τ -lepton and associated neutrino. By measuring this decay width Coarasa *et al.* obtained a physical definition of $\tan \beta$ which can include the one-loop corrections to the decay $H^+ \rightarrow \tau^+ \nu_\tau$ in the MSSM. We will follow the approach of Ref. [9] to fix Z_β below. By introducing appropriate counterterms the renormalized amplitude can be expressed as

$$M_{\text{ren}} = -i\sqrt{2}\bar{u}(\tilde{\chi}_j^0)(aP_R + bP_L)u(t) \quad (8)$$

with

$$a = L_{1j}^* + \delta L_{1j}^* + L_{1j}^* \left(\frac{1}{2} \delta Z_t^R + \frac{1}{2} \delta Z_{\tilde{\chi}_j^0}^R + \frac{1}{2} \delta Z_{11} \right) + L_{2j}^* \delta Z_{12} + \Lambda_R^{\text{SUSY-EW}},$$

$$b = R_{1j}^* + \delta R_{1j}^* + R_{1j}^* \left(\frac{1}{2} \delta Z_t^L + \frac{1}{2} \delta Z_{\tilde{\chi}_j^0}^L + \frac{1}{2} \delta Z_{11} \right) + R_{2j}^* \delta Z_{12} + \Lambda_L^{\text{SUSY-EW}}, \quad (9)$$

where $\Lambda_{L,R}^{\text{SUSY-EW}}$ are the leading SUSY-EW vertex corrections from the irreducible vertex diagrams induced by the Yukawa couplings, expressions for which will be given below. δL_{1j}^* and δR_{1j}^* are the shifts from the bare couplings to renormalized couplings and, as mentioned above, can be found by renormalizing the top Yukawa coupling and the stop mixing angle,

$$\begin{aligned}\delta L_{1j}^* &= L_{2j}^* \delta\theta + L_{1j}^{*(h_t)} \frac{\delta h_t}{h_t}, & \delta R_{1j}^* &= R_{2j}^* \delta\theta + R_{1j}^{*(h_t)} \frac{\delta h_t}{h_t}, \\ L_{1j}^{*(h_t)} &= A_j^* \cos\theta, & R_{1j}^{*(h_t)} &= -D_j^* \sin\theta, \\ \frac{\delta h_t}{h_t} &= \frac{\delta g}{g} - \frac{\delta m_w}{m_w} + \frac{\delta m_t}{m_t} - \cos^2\beta \delta Z_\beta,\end{aligned}\quad (10)$$

where

$$\delta Z_\beta = \frac{\delta m_w}{m_w} - \frac{\delta g}{g} - \frac{1}{2} \delta Z_H - m_w \cot\beta \delta Z_{HW} + \Delta_\tau \quad (11)$$

with

$$\Delta_\tau = -\frac{\delta m_\tau}{m_\tau} - \frac{1}{2} \delta Z_L^{\nu\tau} - \frac{1}{2} \delta Z_R^\tau - F_\tau. \quad (12)$$

Here F_τ is the form factor describing the vertex corrections to the amplitude of $H^+ \rightarrow \tau^+ \nu_\tau$,^[9] δg can be written as:

$$\frac{\delta g}{g} = \frac{\delta e}{e} + \frac{1}{2} \frac{\cos^2\theta_w}{\sin^2\theta_w} \left(\frac{\delta m_w^2}{m_w^2} - \frac{\delta m_z^2}{m_z^2} \right).$$

The counterterms and the renormalization constants in Eqs (8) ~ (12) are defined by

$$\begin{aligned}m_t^0 &= m_t + \delta m_t, & m_{w0}^2 &= m_w^2 + \delta m_w^2, & m_{z0}^2 &= m_z^2 + \delta m_z^2, \\ m_\tau^0 &= m_\tau + \delta m_\tau, & \theta^0 &= \theta + \delta\theta, & \tan\beta_0 &= Z_\beta \tan\beta = (1 + \delta Z_\beta) \tan\beta, \\ t^0 &= Z_t^{1/2} t = (1 + \delta Z_t^L P_L + \delta Z_t^R P_R)^{1/2} t, & \tau^0 &= Z_\tau^{1/2} \tau = (1 + \delta Z_\tau^L P_L + \delta Z_\tau^R P_R)^{1/2} \tau, \\ \nu^0 &= Z_\nu^{1/2} \nu = (1 + \delta Z_\nu^L P_L)^{1/2} \nu, & \tilde{\chi}_j^{0(0)} &= Z_{\tilde{\chi}_j^0}^{1/2} \tilde{\chi}_j^0 = (1 + \delta Z_{\tilde{\chi}_j^0}^L P_L + \delta Z_{\tilde{\chi}_j^0}^R P_R)^{1/2} \tilde{\chi}_j^0, \\ \tilde{t}_1^0 &= (1 + \delta Z_{11})^{1/2} \tilde{t}_1 + \delta Z_{12} \tilde{t}_2, & H_\pm^0 &= Z_H^{1/2} H^\pm = (1 + \delta Z_H)^{1/2} H^\pm\end{aligned}\quad (13)$$

and

$$W_{\pm\mu}^0 = Z_W^{1/2} W_{\pm\mu} + i Z_{HW}^{1/2} \partial H^\pm. \quad (14)$$

We have fixed the wavefunction renormalization constants and the mass counterterms by the on-mass-shell renormalization scheme condition, and also fixed Z_β by following the approach of Ref. [9]. As for the mixing angle counterterm $\delta\theta$, we fix it by the requirement^[3] that the mixing angle counterterm simply is the negative of the counterterm δZ_{12} , that is,

$$\delta\theta = -\delta Z_{12}. \quad (15)$$

This condition ensures that all the ultraviolet divergencies will be canceled in the virtual corrections to the decay width, and is in agreement with the result of Ref. [10].

Calculating the self-energy diagrams and the irreducible vertex diagrams in Fig. 1, we can obtain the explicit expressions for all the renormalization constants and the vertex corrections. Since these expressions are very tedious, we will not show them in this limited letter, and only present the analysis formula of the virtual corrections of order $O(h_t^2)$ to the tree-level decay width, which gives the leading SUSY-EW corrections for the case of low $\tan\beta$ and is given by

$$\begin{aligned}\delta\Gamma_{\text{virt}} &= \frac{1}{16\pi m_t^3} \lambda^{1/2} (m_t^2, m_{\tilde{\chi}_j^0}^2, m_{\tilde{t}_1}^2) \text{Re} \left\{ 2(m_t^2 + m_{\tilde{\chi}_j^0}^2 - m_{\tilde{t}_1}^2) \left[(L_{1j} L_{1j}^{*(h_t)} + R_{1j} R_{1j}^{*(h_t)}) \frac{\delta h_t}{h_t} \right. \right. \\ &\quad \left. \left. + (|L_{1j}|^2 + |R_{1j}|^2) \frac{1}{2} \delta Z_{11} + \frac{1}{2} (|L_{1j}|^2 (\delta Z_t^R + \delta Z_{\tilde{\chi}_j^0}^R) + |R_{1j}|^2 (\delta Z_t^L + \delta Z_{\tilde{\chi}_j^0}^L)) \right] \right\}\end{aligned}$$

$$\begin{aligned}
& + L_{1j} \Lambda_R^{\text{SUSY-EW}} + R_{1j} \Lambda_L^{\text{SUSY-EW}} \Big] + 4m_t m_{\tilde{\chi}_j^0} \left[(L_{1j} R_{1j}^{*(h_t)} + R_{1j} L_{1j}^{*(h_t)}) \frac{\delta h_t}{h_t} \right. \\
& + (L_{1j} R_{1j}^* + R_{1j} L_{1j}^*) \frac{1}{2} \delta Z_{11} + \frac{1}{2} (L_{1j} R_{1j}^* (\delta Z_t^L + \delta Z_{\tilde{\chi}_j^0}^L) + R_{1j} L_{1j}^* (\delta Z_t^R \\
& \left. + \delta Z_{\tilde{\chi}_j^0}^R)) + L_{1j} \Lambda_L^{\text{SUSY-EW}} + R_{1j} \Lambda_R^{\text{SUSY-EW}} \Big] \Big\}. \quad (16)
\end{aligned}$$

We have checked analytically that all the ultraviolet divergencies are indeed canceled in the above virtual corrections to the decay width.

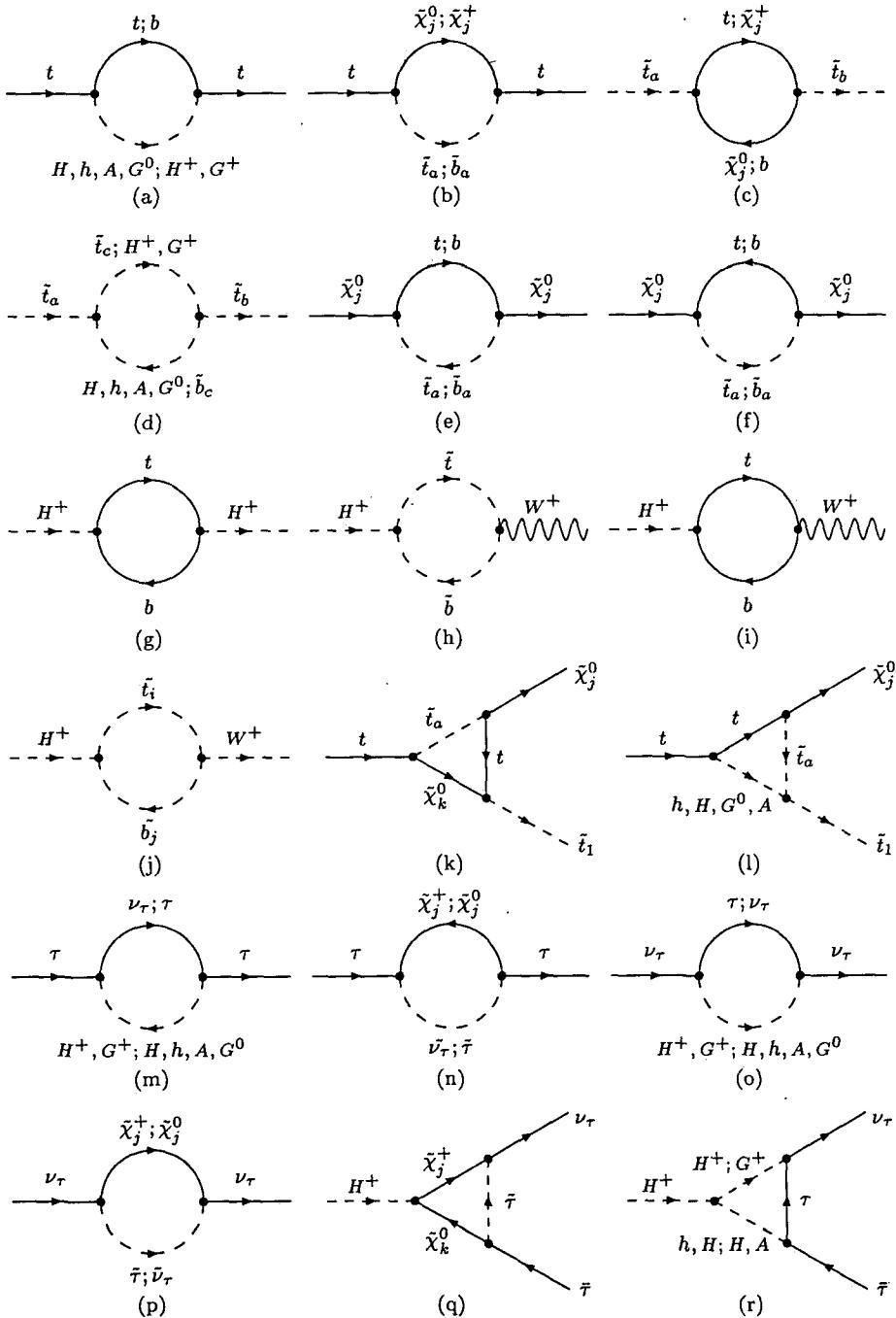


Fig. 1. Feynman diagrams for the leading SUSY-EW corrections.

In the following we give the numerical results for the leading SUSY-EW corrections to $t \rightarrow \tilde{t}_1 \tilde{\chi}_1^0$, where $\tilde{\chi}_1^0$ is the lightest neutralino. The dominant part of such corrections arises from the top Yukawa coupling and is of order $O(h_t^2)$ for small $\tan \beta$. So, we will limit the value of $\tan \beta$ to be in the range $1.5 \leq \tan \beta \leq 11$, which is consistent with that required by the most popular MSSM model with scenarios motivated by current low energy data (including α_s, R_b , the $e e \gamma\gamma + E_T$ event and the branching ratio of $b \rightarrow s\gamma$). Although the larger values of $\tan \beta$ are also allowed by the current data, for simplicity, as an approximation, we will only consider the case of small $\tan \beta$ as mentioned above, and neglect the effects of the bottom and τ Yukawa couplings in the numerical calculations below.

First of all, in order to compare with the one-loop QCD corrections, we use the same input parameters as those in Ref. [3], i.e., we fix $M = 200$ GeV, $\mu = -100$ GeV and we use the relation $M' = \frac{5}{3}(g'/g^2)M^{[5]}$ to fix M' . For the parameters in stop sector we assume $M_{\tilde{t}_R} = M_{\tilde{t}_L}$ and take the combination $A_t + \mu \cot \beta$ to be one parameter. There are then three free parameters in the stop sector and we choose $m_{\tilde{t}_1}$, $\tan \beta$, and $(A_t + \mu \cot \beta)$ as the three independent parameters. Moreover, we use the relation^[11] between the Higgs boson masses m_{h_0, H, A, H^\pm} and parameters β at one loop, and choose m_A and $\tan \beta$ as two independent input parameters. The SM input parameters are $m_z = 91.188$ GeV, and $G_F = 1.166372 \times 10^{-5}$ (GeV)⁻². The W mass was determined from^[12]

$$m_w^2 \left(1 - \frac{m_w^2}{M_Z^2}\right) = \frac{\pi \alpha}{\sqrt{2} G_F} \frac{1}{1 - \Delta r}, \quad (17)$$

where, for a heavy top, Δr is given by^[13]

$$\Delta r \sim -\frac{\alpha N_C c_w^2 m_t^2}{16\pi^2 s_w^4 m_w^2}. \quad (18)$$

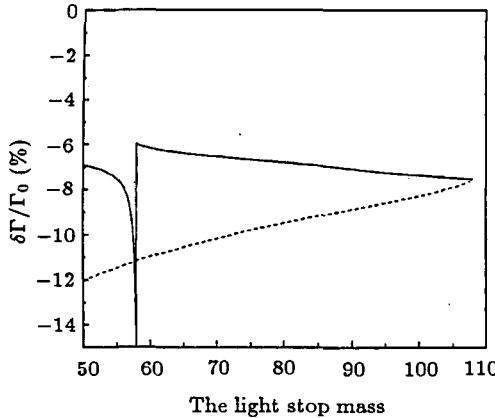


Fig. 2. The relative correction $\delta\Gamma/\Gamma_0$ to the decay rate as a function of the light stop mass assuming $M = 200$ GeV, $\mu = -100$ GeV, $\tan \beta = 11$ and $A_t + \mu \cot \beta = 100$ GeV. The solid and dashed curves correspond to the leading SUSY-EW corrections and SUSY-QCD corrections, respectively.^[3]

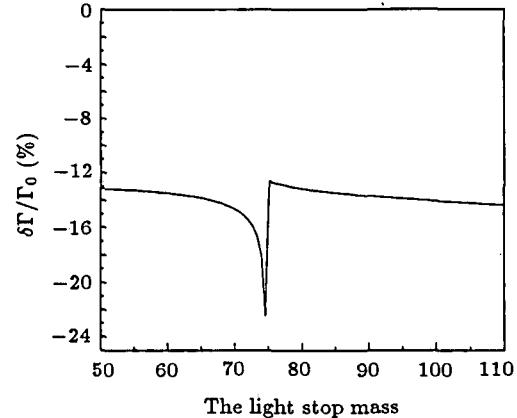


Fig. 3. The relative correction $\delta\Gamma/\Gamma_0$ to the decay rate as a function of the light stop mass assuming $M = 100$ GeV, $\mu = -60$ GeV, $\tan \beta = 2$ and $A_t + \mu \cot \beta = 200$ GeV.

Figure 2 shows the leading SUSY-EW corrections of order $O(h_t^2)$ to the decay rate $\delta\Gamma/\Gamma_0$ as a function of the light stop mass assuming $\tan \beta = 11$. The corrections are typically about $-6\% \sim -7\%$ and are not sensitive to the light stop mass except that $m_{\tilde{t}_1} = 57$ GeV, where there is a peak due to the fact that $m_t = 175$ GeV, $m_{\tilde{\chi}_1} = (68, 116, 118, 233)$ GeV for the

parameters used above, and the threshold for open top quark decay into a neutralino and a light stop is crossed in this region, which is caused by Fig. 1k. Since the QCD corrections^[3] are about $-8\% \sim -12\%$ in this case, the combined effects of the leading SUSY-EW and one-loop QCD can exceed -15% in a light stop mass range allowed by the kinematics.

Figure 3 gives the leading SUSY-EW corrections as a function of $m_{\tilde{t}_1}$ with $\tan\beta = 2$. Since the corrections are proportional to $m_t^2/m_w^2 \sin^2\beta$, they can be very large for small $\tan\beta$. From Fig. 3 one sees that the leading SUSY-EW corrections exceed -12% for $50 \leq m_{\tilde{t}_1} \leq 120$ GeV, and the corrections have a large peak at about 74 GeV, which also comes from the threshold effect of $t \rightarrow \tilde{t}_1 + \tilde{\chi}_j^0$ caused by Fig. 1k. In addition, comparing with the one-loop QCD corrections^[3] in the case of $\tan\beta = 2$, we find that the leading SUSY-EW corrections dominate over the QCD corrections.

In conclusion, we calculated the leading SUSY-EW corrections of order $O(h_t^2)$ to $t \rightarrow \tilde{t}_1 + \tilde{\chi}_j^0$ for the case of low $\tan\beta$, and found that such corrections can exceed -6% for $\tan\beta = 11$ and -12% for $\tan\beta = 2$, respectively. And these corrections vary insensitively with $m_{\tilde{t}_1}$ in the region allowed by the kinematics except for those near the threshold.

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