PHYSICAL REVIEW D, VOLUME 60, 035007

Flavor changing top decay $t\to c\tilde{\nu}$ or $\tilde{\nu}\to t\bar{c}$ in the MSSM without R parity

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(Received 31 December 1998; published 7 July 1999)

Widths for the new flavor-changing top quark decay $t\to c\tilde{\nu}$ or for the reversed sneutrino decay $\tilde{\nu}\to t\bar{c}$ are calculated in the minimal supersymmetric standard model without R-parity conservation. For large $\tan\beta$, e.g., $\tan\beta\sim m_t/m_b\sim 40$, $\mathrm{Br}(t\to c\tilde{\nu})>10^{-5}$ or $\mathrm{Br}(\tilde{\nu}\to t\bar{c})>10^{-6}$ in a relatively wide range of the supersymmetric parameter space as long as there is more than one nonzero R-parity violating coupling. In the best cases, with a typical squark mass around 100 GeV, we find that $\mathrm{Br}(t\to c\tilde{\nu})\sim 10^{-4}-10^{-3}$ or $\mathrm{Br}(\tilde{\nu}\to t\bar{c})\sim 10^{-5}-10^{-4}$. For $\tan\beta\sim\mathcal{O}(1)$ the corresponding branching ratios for both top or sneutrino decays are too small to be measured at the CERN Large Hadron Collider (LHC). Therefore, the decays $t\to c\tilde{\nu}$ or $\tilde{\nu}\to t\bar{c}$ appear to be sensitive to $\tan\beta$ and may be detected at the LHC. The branching ratios of the corresponding decays to an up quark instead of a charm quark, e.g., $t\to u\tilde{\nu}$ or $\tilde{\nu}\to t\bar{u}$, may also be similar. [S0556-2821(99)04911-5]

PACS number(s): 14.65.Ha, 12.60.Jv, 14.80.Ly

I. INTRODUCTION

Flavor changing decays of heavy quarks which have been the subject of intense theoretical activity for a long time are especially significant as they provide important tests of the standard model (SM). The extraordinary large mass of the top quark renders the Glashow-Iliòpoulos-Maiani (GIM) mechanism very effective in the SM, so that the flavor-changing top decays are highly suppressed [1]. Experimental searches for the flavor-changing decays of the top are therefore very good probes of new physics. It is expected that the CERN Large Hadron Collider (LHC) will produce about 10^7-10^8 $t\bar{t}$ pairs, therefore rare top decays with branching ratios around 10^{-6} should be accessible at the LHC.

In the SM, the decays $t\to cV$ ($V=\gamma$, Z, or g) and $t\to cH^0$, have branching ratios (Br) much smaller than 10^{-6} [1,2] and should therefore be very useful in searching for new physics at the LHC. In some extensions of the SM such as multi-Higgs doublet models [1,3] and the minimal supersymmetric standard model (MSSM) [4], $\mathrm{Br}(t\to cV)$ can be several orders of magnitude larger than their SM value. However, although in the MSSM $\mathrm{Br}(t\to cg)\sim 10^{-5}$ may be possible, for $t\to c\gamma$ and $t\to cZ$, typically, $\mathrm{Br}(t\to c\gamma,cZ)\lesssim 10^{-6}$ [1,3,4].

The most optimistic results for $\mathrm{Br}(t\to cV)$ were reported in [5], where it was shown that, in the MSSM without R parity, these branching ratios can be as large as $\mathrm{Br}(t\to cg)\sim 10^{-3}$, $\mathrm{Br}(t\to cZ)\sim 10^{-4}$, and $\mathrm{Br}(t\to c\gamma)\sim 10^{-5}$ if the squarks have masses not much larger than 100 GeV.

For the decay $t\to cH^0$ the theoretical predictions for the branching ratio span several orders of magnitude as one considers beyond the SM scenarios. While in the MSSM $\mathrm{Br}(t\to cH^0)$ can reach $\sim 10^{-5}$ at the most [6], in a class of multi-Higgs doublets models in which the neutral Higgs boson can have a nonvanishing tree-level coupling to a $t\bar{c}$ pair, $\mathrm{Br}(t\to cH^0)$ around 10^{-2} is not ruled out [7].

In this paper we explore the new flavor-changing top decays $t\to c\tilde{\nu}$, $t\to c\overline{\tilde{\nu}}$ ($\tilde{\nu}=$sneutrino, $\bar{\tilde{\nu}}=$anti-sneutrino). These decays require the lepton number to be violated and, therefore, cannot occur within supersymmetry in its minimal version. We, therefore, consider $t\to c\tilde{\nu}$, $c\bar{\tilde{\nu}}$ or $\tilde{\nu}$, $\bar{\tilde{\nu}}\to t\bar{c}$, $\bar{t}c$ (depending on whether $m_t>m_{\tilde{\nu}}$ or $m_t<m_{\tilde{\nu}}$, respectively) in the MSSM with R-parity violation ($\not{R}_P$) [8].

The $\not{R}_P$ interaction of interest to us is the λ' type—lepton number violating—operator [8],

$$W_{\not{L}}=\lambda'_{ijk}\hat{L}_i\hat{Q}_j\hat{D}^c_k, \tag{1}$$

and the corresponding trilinear soft breaking operator:

$$\mathcal{L}^{\mathrm{soft}}_{\not{L}}=\lambda'_{ijk}A'\tilde{l}_i\tilde{q}_j\tilde{d}^c_k, \tag{2}$$

where in Eq. (1), $\hat{L}$ and $\hat{Q}$ are the SU(2)-doublet lepton and quark superfields, respectively and $\hat{D}^c$ are the quark singlet superfields. In Eq. (2) $\tilde{l}$, $\tilde{q}$, and $\tilde{d}$ are the sleptons, left-handed squarks, and right-handed down squarks corresponding to the superfields $\hat{L}$, $\hat{Q}$, and $\hat{D}$, respectively. Also, for simplicity, the trilinear soft breaking couplings in Eq. (2) have been defined to include the Yukawa-type λ' couplings. Notice that A' has a mass dimension and should naturally attain values of the order of the typical squark mass or equivalently the typical supersymmetry (SUSY) mass scale (see, e.g., [9], and references therein).

It is important to note that, on a purely phenomenological level, since there is no good theoretical reason which forbids $\not{R}_P$ SUSY models, the λ' couplings in Eqs. (1) and (2) are *a priori* expected to be of $\mathcal{O}(1)$ (see [10]). Existing bounds on the λ' couplings from low-energy processes suggest that, typically, $\lambda'_{ijk}<0.5$ for squark and slepton masses of 100 GeV. We wish to emphasize that these bounds are usually based on the assumption that an $\not R_P$ coupling, e.g., λ'_{ijk} is nonzero only for one combination of the indices *ijk* [8]. Indeed, if this assumption is relaxed, then in most cases $\lambda'_{ijk} \sim \mathcal{O}(1)$ cannot be ruled out.

Another useful observation is that in the presence of $\not R_P$ interactions, existing limits on the SUSY spectrum which are obtained from high-energy collider experiments become less restrictive, since in this model the lightest supersymmetric particle (LSP) is no longer stable. That is, the distinct missing energy signal associated with the LSP in the MSSM with *R*-parity conservation is lost. Thus, for example, squark masses of the order of $\sim$100 GeV are still allowed by present data [8].

With that in mind, in this work we show that if $m_t > m_{\tilde\nu}$, then $\mathrm{Br}(t\to c\tilde\nu)>10^{-5}$ is attainable within a relatively wide range of the relevant SUSY parameter space and may reach $10^{-4}-10^{-3}$ for $\tan\beta\sim m_t/m_b\sim 40$ (recall that $\tan\beta$ is the ratio between the vacuum expectation values of the two Higgs fields in the model) and with the typical squark mass around 100 GeV. Similarly, in the case $m_t < m_{\tilde\nu}$, we find that $\mathrm{Br}(\tilde\nu\to t\bar c)>10^{-6}$ is possible and can reach $10^{-5}-10^{-4}$ if, again, $M_s\sim 100$ GeV (where M_s denotes the squark mass) and $\tan\beta\sim m_t/m_b\sim 40$. In fact, if indeed some of the λ' couplings are saturated to be of order one, then the LHC can effectively serve as a sneutrino factory having about 10^8-10^9 sneutrinos with a mass of 200–300 GeV produced inclusively [11]. Therefore, if $m_t<m_{\tilde\nu}$, the above flavor-changing sneutrino decay, $\tilde\nu\to t\bar c$, appears to be as useful as the top decay $t\to c\tilde\nu$, which is of course relevant only for the corresponding mass range $m_t>m_{\tilde\nu}$.

We also note that, for an up quark instead of a charm quark, the branching ratios remain practically the same, i.e., $\mathrm{Br}(t\to u\tilde\nu)\simeq\mathrm{Br}(t\to c\tilde\nu)$ or $\mathrm{Br}(\tilde\nu\to t\bar u)\simeq\mathrm{Br}(\tilde\nu\to t\bar c)$ if the $\not R_P$ couplings relevant for the two cases are the same.

The paper is organized as follows: in Sec. II we present an analytical derivation and the numerical results for the branching ratio of $t\to c\tilde\nu$, $c\bar{\tilde\nu}$ in the MSSM with $\not R_P$. In Sec. III we evaluate the sneutrino and anti-sneutrino decays $\tilde\nu$, $\bar{\tilde\nu}\to t\bar c$, $\bar t c$ in the case $m_t<m_{\tilde\nu}$, and in Sec. IV we summarize.

II. THE CASE $m_t>m_{\tilde\nu}$ AND THE TOP DECAYS $t\to c\tilde\nu$, $t\to c\bar{\tilde\nu}$

The relevant Feynman diagrams responsible for the decay $t\to c\tilde\nu$ at the one-loop level and with insertion of only one $\not R_P$ coupling are depicted in Fig. 1. Those penguin-like diagrams give rise to an effective amplitude for the decay $t\to c\tilde\nu$ of the form

$$-i\mathcal{M}^{\tilde\nu}=\mathcal{C}^{\tilde\nu}\overline{u_c}\left[\left(\frac{m_c}{m_W}\right)\sum_\alpha \mathcal{A}_\alpha^{\tilde\nu}L+\left(\frac{m_t}{m_W}\right)\sum_\alpha \mathcal{B}_\alpha^{\tilde\nu}R\right]u_t, \quad (3)$$

where $L(R)=(1-(+)\gamma_5)/2$ and

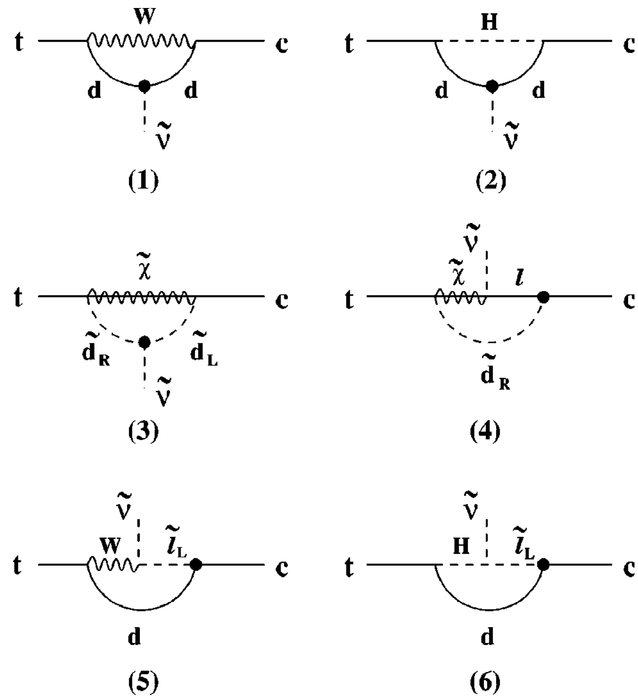

FIG. 1. The one-loop Feynman diagrams for the decay $t\to c\tilde\nu$, induced by an insertion of one lepton number violating λ' coupling. The heavy dot denotes the lepton number violating vertex. See also [12].

$$\mathcal{C}^{\tilde\nu}\equiv\frac{\alpha}{8\pi s_W^2}\frac{m_{d_k}}{m_W}\lambda'^{*}_{ijk}V_{tk}V^*_{cj}. \quad (4)$$

Also, s_W is the sine of the weak mixing angle and m_{d_k} is the mass of the down quark of the kth generation ($k=1$, 2 or 3). The form factors $A_\alpha^{\tilde\nu}$ and $B_\alpha^{\tilde\nu}$ corresponding to diagrams $\alpha=1-6$ in Fig. 1 are given by (keeping all the masses)

$$\mathcal{A}_1^{\tilde\nu}=m_t^2(C_{21}^1+C_{22}^1-2C_{23}^1)-2m_W^2(C_0^1+C_{12}^1)\\ -m_c^2C_{22}^1+2C_{24}^1-\tilde C_{12}^1, \quad (5)$$

$$\mathcal{A}_2^{\tilde\nu}=m_t^2(1+\cot^2\beta)(C_{12}^2-C_{11}^2)-m_c^2C_{12}^2\\ -m_{d_k}^2\tan^2\beta C_{12}^2-\tilde C_0^2, \quad (6)$$

$$\mathcal{A}_3^{\tilde\nu}=-A'\sum_{m=1,2}\frac{Z_{2m}^-}{\cos\beta}\left(\frac{Z_{2m}^+}{\sin\beta}m_{\tilde\chi m}C_0^3+\sqrt{2}Z_{1m}^{-*}m_WC_{12}^3\right), \quad (7)$$

$$\mathcal{A}_4^{\tilde\nu}=\sum_{m=1,2}\frac{Z_{2m}^-}{\cos\beta}\left(m_{l_i}^2\frac{Z_{2m}^{-*}}{\cos\beta}C_{12}^4-\sqrt{2}Z_{1m}^+m_Wm_{\tilde\chi m}(C_0^4\\ +C_{12}^4)\right), \quad (8)$$

$$\mathcal{A}_5^{\tilde{\nu}}=(m_t^2-m_c^2)(C_{22}^5-C_{23}^5)-m_{\tilde{\nu}_i}^2(C_{22}^5+C_{23}^5)+m_W^2(C_{12}^5+2C_0^5)-2C_{24}^5-\tilde{C}_{12}^5, \tag{9}$$

$$\mathcal{A}_6^{\tilde{\nu}}=(m_W^2\sin 2\beta-m_{l_i}^2\tan\beta)\tan\beta C_{12}^6, \tag{10}$$

and:

$$\mathcal{B}_1^{\tilde{\nu}}=(m_t^2-m_{\tilde{\nu}_i}^2)(2C_{23}^1-C_{21}^1-C_{22}^1)-2m_W^2(C_{11}^1-C_{12}^1)+m_c^2(C_{22}^1-C_{21}^1)+\tilde{C}_{12}^1-\tilde{C}_{11}^1, \tag{11}$$

$$\mathcal{B}_2^{\tilde{\nu}}=m_{d_k}^2\tan^2\beta(C_{12}^2-C_{11}^2-C_0^2)-m_{d_k}^2C_0^2-m_c^2(C_0^2+C_{11}^2)-m_c^2\cot^2\beta(C_0^2+C_{12}^2), \tag{12}$$

$$\mathcal{B}_3^{\tilde{\nu}}=\sqrt{2}A'm_W\sum_{m=1,2}\frac{Z_{2m}^-Z_{1m}^-}{\cos\beta}(C_{12}^3-C_{11}^3), \tag{13}$$

$$\mathcal{B}_4^{\tilde{\nu}}=\sum_{m=1,2}\frac{Z_{2m}^-}{\cos\beta}\left(\sqrt{2}Z_{1m}^+m_Wm_{\tilde{\chi}_m}(C_0^4+C_{11}^4)-m_{l_i}^2\frac{Z_{2m}^{-*}}{\cos\beta}C_{11}^4\right), \tag{14}$$

$$\mathcal{B}_5^{\tilde{\nu}}=(m_t^2-m_c^2)(C_{21}^5-C_{23}^5)+m_{\tilde{\nu}_i}^2(C_{21}^5+C_{23}^5)-m_W^2(C_{11}^5+2C_0^5)+2C_{24}^5+\tilde{C}_{11}^5, \tag{15}$$

$$\mathcal{B}_6^{\tilde{\nu}}=-(m_W^2\sin 2\beta-m_{l_i}^2\tan\beta)(\tan\beta(C_{11}^6+C_0^6)+\cot\beta C_0^6), \tag{16}$$

where the lepton flavor (l_i) as well as the sneutrino flavor ($\tilde{\nu}_i$), $i=1$, 2, or 3, are set by the choice of the first index in the $R\!\!\!/_P$ coupling λ'_{ijk}. Also, $m_{\tilde{\chi}m}$, $m=1,2$, are the chargino masses and $Z^{\pm}$ are the 2×2 matrices that diagonalize the charginos 2×2 mass matrix (see, e.g., Ref. [13]). The C_0^α, $\tilde{C}_0^\alpha$, C_{pq}^α, and $\tilde{C}_{pq}^\alpha$, $p=1,2$ and $q=1-4$, in Eqs. (5)–(16) are the three-point loop from factors associated with the loop integrals of diagrams $\alpha=1$–6 in Fig. 1. Using the notation in [13], they are given by

$$C^1\equiv C^1(m_W^2,m_{d_j}^2,m_{d_k}^2,m_t^2,m_{\tilde{\nu}_i}^2,m_c^2), \tag{17}$$

$$C^2\equiv C^2(m_{H^+}^2,m_{d_j}^2,m_{d_k}^2,m_t^2,m_{\tilde{\nu}_i}^2,m_c^2), \tag{18}$$

$$C^3\equiv C^3(m_{\tilde{\chi}_m}^2,m_{\tilde{d}_{Rj}}^2,m_{\tilde{d}_{Lk}}^2,m_t^2,m_{\tilde{\nu}_i}^2,m_c^2), \tag{19}$$

$$C^4\equiv C^4(m_{\tilde{\chi}_m}^2,m_{\tilde{d}_{Rj}}^2,m_{l_i}^2,m_t^2,m_c^2,m_{\tilde{\nu}_i}^2), \tag{20}$$

$$C^5\equiv C^5(m_W^2,m_{d_j}^2,m_{\tilde{l}_{Li}}^2,m_t^2,m_c^2,m_{\tilde{\nu}_i}^2), \tag{21}$$

$$C^6\equiv C^6(m_{H^+}^2,m_{d_j}^2,m_{\tilde{l}_{Li}}^2,m_t^2,m_c^2,m_{\tilde{\nu}_i}^2). \tag{22}$$

Note that since the arguments are the same for all the C^α's belonging to diagram α, we have omitted their subscripts and tildes. $m_{\tilde{d}_{Rj}}$ and $m_{\tilde{d}_{Lk}}$ are the masses of the right-handed and left-handed squarks of the j and k generation (set by the second and third index in λ'), respectively. Also, $m_{\tilde{l}_{Li}}$ is the mass of the left-handed slepton of flavor i corresponding to the first index in λ', and m_{H^+} is the charged Higgs boson mass.

Notice that $\tilde{C}_0$, $\tilde{C}_{11}$, $\tilde{C}_{12}$, and C_{24} contain ultraviolet divergences which, in dimensional regularization, appear as $1/\epsilon$ where $\epsilon=4-D$ in D dimensions. Of course, $\Sigma_{\alpha=1}^6\mathcal{A}_\alpha$ and $\Sigma_{\alpha=1}^6\mathcal{B}_\alpha$ should be free of those infinities; indeed, as can be checked analytically from Eqs. (5)–(16), those infinities cancel out.

The width for $t\to c\tilde{\nu}_i$ is then given by

$$\Gamma(t\to c\tilde{\nu}_i)=\frac{|\mathcal{C}^{\tilde{\nu}}|^2}{32\pi}\frac{\omega(m_t^2,m_c^2,m_{\tilde{\nu}_i}^2)}{m_t^3}\frac{m_t^2m_c^2}{m_W^4}\left\{(m_t^2+m_c^2-m_{\tilde{\nu}_i}^2)\times\left(\frac{m_W^2}{m_t^2}|\overline{\mathcal{A}^{\tilde{\nu}}}|^2+\frac{m_W^2}{m_c^2}|\overline{\mathcal{B}^{\tilde{\nu}}}|^2\right)+4m_W^2\,\mathrm{Re}[\overline{\mathcal{A}^{\tilde{\nu}}}(\overline{\mathcal{B}^{\tilde{\nu}}})^*]\right\}, \tag{23}$$

where we have defined

$$\overline{\mathcal{A}^{\tilde{\nu}}}\equiv\sum_{\alpha=1}^{6}\mathcal{A}_\alpha^{\tilde{\nu}},\quad \overline{\mathcal{B}^{\tilde{\nu}}}\equiv\sum_{\alpha=1}^{6}\mathcal{B}_\alpha^{\tilde{\nu}}, \tag{24}$$

and

$$\omega(a^2,b^2,c^2)\equiv\sqrt{[a^2-(b+c)^2][a^2-(b-c)^2]}. \tag{25}$$

For the top decay to an antisneutrino, $t\to c\bar{\tilde{\nu}}$, the amplitude can be similarly written as

$$-i\mathcal{M}^{\bar{\tilde{\nu}}}=\mathcal{C}^{\bar{\tilde{\nu}}}\overline{u_c}\left[\left(\frac{m_c}{m_W}\right)\sum_\alpha\mathcal{A}_\alpha^{\bar{\tilde{\nu}}}L+\left(\frac{m_t}{m_W}\right)\sum_\alpha\mathcal{B}_\alpha^{\bar{\tilde{\nu}}}R\right]u_t, \tag{26}$$

where now

$$\mathcal{C}^{\bar{\tilde{\nu}}}\equiv\frac{\alpha}{8\pi s_W^2}\frac{m_{d_k}}{m_W}\lambda'_{ijk}V_{tj}V_{ck}^*, \tag{27}$$

and the form factors $A_\alpha^{\bar{\tilde{\nu}}},B_\alpha^{\bar{\tilde{\nu}}}$, for $\alpha=1$–6, are related to $A_\alpha^{\tilde{\nu}},B_\alpha^{\tilde{\nu}}$, as follows

$$A_\alpha^{\bar{\tilde{\nu}}}=[B_\alpha^{\tilde{\nu}}(m_t\leftrightarrow m_c)]^*, \tag{28}$$

$$B_\alpha^{\bar{\tilde{\nu}}}=[A_\alpha^{\tilde{\nu}}(m_t\leftrightarrow m_c)]^*. \tag{29}$$

The width $\Gamma(t\to c\bar{\tilde{\nu}}_i)$ is then given by Eq. (23) with the replacements $\mathcal{C}^{\tilde{\nu}}\to\mathcal{C}^{\bar{\tilde{\nu}}}$, $\overline{\mathcal{A}^{\tilde{\nu}}}\to\overline{\mathcal{A}^{\bar{\tilde{\nu}}}}$ and $\overline{\mathcal{B}^{\tilde{\nu}}}\to\overline{\mathcal{B}^{\bar{\tilde{\nu}}}}$, where

$$\overline{\mathcal{A}^{\bar{\tilde{\nu}}}}\equiv\sum_{\alpha=1}^{6}\mathcal{A}_\alpha^{\bar{\tilde{\nu}}},\quad \overline{\mathcal{B}^{\bar{\tilde{\nu}}}}\equiv\sum_{\alpha=1}^{6}\mathcal{B}_\alpha^{\bar{\tilde{\nu}}}, \tag{30}$$

and $\mathcal{A}_\alpha^{\bar{\tilde{\nu}}}$ and $\mathcal{B}_\alpha^{\bar{\tilde{\nu}}}$ are given in Eqs. (28) and (29), respectively.

We are now interested in calculating the branching ratio for a top quark to decay to $c\tilde{\nu}_i$ and to $c\bar{\tilde{\nu}}_i$. In what follows we will assume that $m_{H^+}>m_t$, therefore, the branching ratio in question is given by

$$\mathrm{Br}_i \equiv \frac{\Gamma(t\to c\tilde{\nu}_i)+\Gamma(t\to c\bar{\tilde{\nu}}_i)}{\Gamma(t\to bW^+)}. \tag{31}$$

It is clear from Eq. (4) that $\Gamma(t\to c\tilde{\nu}_i)$ is largest for $j=2$ and $k=3$, i.e., for λ'_{i23}, in which case it is proportional to the diagonal Cabibbo-Kobayashi-Maskawa (CKM) elements $V_{tb}V_{cs}^*\sim 1$. Similarly, from Eq. (27) we see that $\Gamma(t\to c\bar{\tilde{\nu}}_i)$ is largest for $j=3$ and $k=2$, being suppressed by the small off-diagonal CKM elements if $\lambda'_{ijk}\neq\lambda'_{i32}$. Therefore, taking $\lambda'_{i23}\neq 0$ for $i=1$, 2 or 3, one finds that $\Gamma(t\to c\tilde{\nu}_i)\gg\Gamma(t\to c\bar{\tilde{\nu}}_i)$, leading to $\mathrm{Br}_i\approx\Gamma(t\to c\tilde{\nu}_i)/\Gamma(t\to bW^+)$. In fact, for a given sneutrino flavor i, even if one allows $|\lambda'_{i23}|=|\lambda'_{i23}|\neq 0$, in which case the leading contribution to both $\Gamma(t\to c\tilde{\nu}_i)$ and $\Gamma(t\to c\bar{\tilde{\nu}}_i)$ is proportional to $V_{tb}V_{cs}^*$, one still finds that $\Gamma(t\to c\tilde{\nu}_i)/\Gamma(t\to c\bar{\tilde{\nu}}_i)\approx(\mathcal{C}^{\tilde{\nu}}/\mathcal{C}^{\bar{\tilde{\nu}}})^2\simeq m_b^2/m_s^2\gg 1$.

In the one-coupling scheme, i.e., only one λ' is nonzero, and if the relative mixing in the quark sector is solely due to absolute mixing in the up-quark sector, the existing 1σ bounds are $\lambda'_{i23}\lesssim 0.2$ for sleptons and squarks masses of 100 GeV [8]. These bounds are obtained from data on $D^0-\bar{D}^0$ mixing [14] and D decays [15]. However, if the one-coupling scheme assumption is relaxed, then these bounds no longer apply due to possible cancelations between the λ''s with different indices (for details, see [14]).

In fact, the branching ratio stays practically the same if, for example, some other λ'_{ijk} with $j\neq 2$ and $k\neq 3$ are also nonzero, since their contribution to $\Gamma(t\to c\tilde{\nu}_i)$ is suppressed by off-diagonal CKM elements. Therefore, for our purpose it is sufficient to assume that there are only three nonzero couplings: $\lambda'_{i23}\neq 0$ for $i=1,2,3$. The decays to an e-sneutrino, μ-sneutrino, and to a τ-sneutrino occur with the same rate if the three sneutrinos are degenerate with a mass $<m_t$ and if $|\lambda'_{123}|=|\lambda'_{223}|=|\lambda'_{323}|$. In that case, $\mathrm{Br}\equiv\mathrm{Br}_1=\mathrm{Br}_2=\mathrm{Br}_3$ and the branching ratio for a top quark to decay to all sneutrino flavors is simply the sum $\Sigma_{i=1}^3\ \mathrm{Br}_i=3\mathrm{Br}$. For definiteness, and without loss of generality, throughout our analysis we present results for a given sneutrino flavor, say the τ-sneutrino $(i=3)$, and we drop the sneutrino index i. We take $|\lambda'_{323}|\equiv\lambda'=1$ (which, as mentioned above, is not ruled out once the one coupling scheme approach is not realized), thus scaling out the $\not{R}_P$ coupling from the branching ratio and presenting results for

$$B^t\equiv \mathrm{Br}/|\lambda'|^2, \tag{32}$$

where Br is defined in Eq. (31).

Before we continue we note that, to one-loop order, there is one additional diagram contributing to $t\to c\tilde{\nu}$ and similarly to $t\to c\bar{\tilde{\nu}}$. This diagram is shown in Fig. 2 and is purely $\not{R}_P$ since it involves insertion of three $\not{R}_P$ couplings. The amplitude corresponding to this diagram has the same structure as in Eq. (3), with the replacement

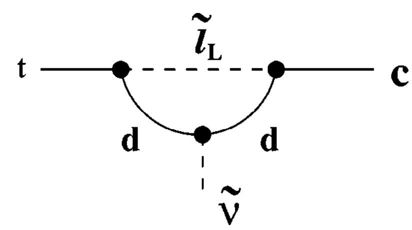

FIG. 2. The one-loop Feynman diagram for the decay $t\to c\tilde{\nu}$, induced by an insertion of three lepton number violating λ' couplings. The heavy dots denote the lepton number violating vertices.

$$\mathcal{C}^{\tilde{\nu}}\to-\frac{\lambda'^*_{lnj}\lambda'_{lpk}}{16\pi^2}\lambda'^*_{ijk}\frac{m_{d_j}}{m_W}V_{tp}V_{cn}^*, \tag{33}$$

and with the form factors

$$\mathcal{A}^{\tilde{\nu}}=m_W^2 C_{12}, \tag{34}$$

$$\mathcal{B}^{\tilde{\nu}}=m_W^2(C_0+C_{11}-C_{12}), \tag{35}$$

where

$$C\equiv C(m_{\tilde{l}_{L_l}}^2,m_{d_k}^2,m_{d_j}^2,m_t^2,m_{\tilde{\nu}_i}^2,m_c^2). \tag{36}$$

For $t\to c\bar{\tilde{\nu}}$ we get

$$\mathcal{C}^{\bar{\tilde{\nu}}}\to-\frac{\lambda'^*_{lnk}\lambda'_{lpj}}{16\pi^2}\lambda'_{ijk}\frac{m_{d_j}}{m_W}V_{tp}V_{cn}^*, \tag{37}$$

and, again, the corresponding form factors $\mathcal{A}^{\bar{\tilde{\nu}}}$ and $\mathcal{B}^{\bar{\tilde{\nu}}}$ are extracted from $\mathcal{A}^{\tilde{\nu}}$ and $\mathcal{B}^{\tilde{\nu}}$ using the relations in Eqs. (28) and (29).

From Eqs. (33) and (37) we see that the amplitude of the diagram in Fig. 2 is largest when $j=3$ $(m_{d_j}=m_b)$ and for $p=3$, $n=2$ $(V_{tp}V_{cn}^*=V_{tb}V_{cs}^*)$. This will require both couplings of the type λ'_{l2k} and λ'_{l3k} to be nonzero. However, even within such a coupling scenario, we find that the contribution of the pure $\not{R}_P$ diagram, in the best cases and for high $\tan\beta$ values where—as will be demonstrated below—B^t can become as much as $\sim 10^{-4}$, is typically one-order of magnitude smaller than that of the diagrams in Fig. 1. We will, therefore, assume for simplicity that $\lambda'_{l3k}=0$, thus neglecting the contribution from the pure $\not{R}_P$ diagram for both $\Gamma(t\to c\tilde{\nu})$ and $\Gamma(t\to c\bar{\tilde{\nu}})$.

The free parameters of the low-energy SUSY relevant for the decays in question are $m_{\tilde{\nu}}$, $m_{\tilde{d}_R}$, $m_{\tilde{d}_L}$, $m_{\tilde{l}_L}$, m_{H^+}, $m_{\tilde{\chi}_m}$, A', and $\tan\beta$. To simplify our analysis below we wish to reduce the number of free parameters by making some plausible simplifying assumptions on the low-energy SUSY spectrum. In particular, we find that some of the above parameters have very little effect on B^t. We fix the values of these parameters and vary the rest.

(1) We find that $\Gamma(t\to c\tilde{\nu})$ is practically insensitive to the slepton and charged Higgs boson masses. We, therefore, set $m_{H^+}=200\,\mathrm{GeV}$ and $m_{\tilde{l}_L}=m_{\tilde{\nu}}$.

(2) Since a possible mass splitting between the left- and right-handed down squarks has no effect on our scaled

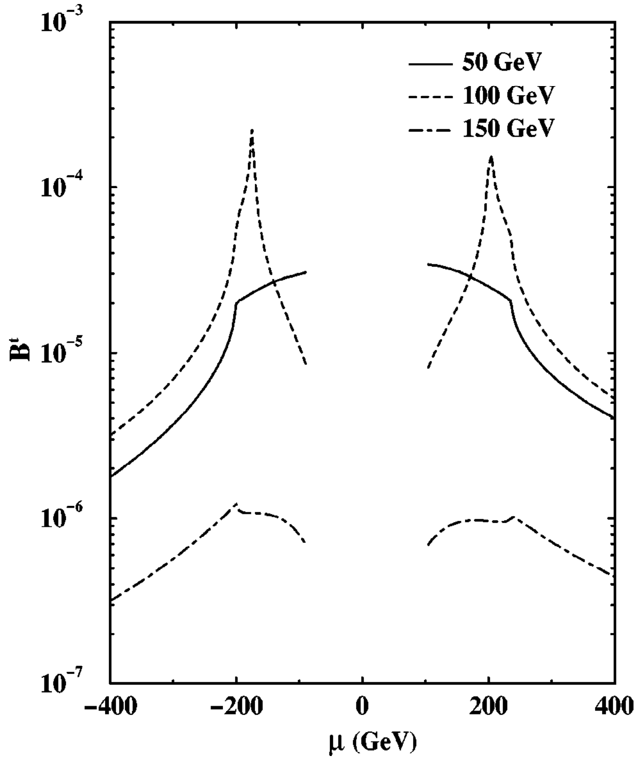

FIG. 3. The scaled branching ratio B^t, defined in Eq. (32), for a single neutrino flavor, as a function of μ for $m_{\tilde{\nu}}=50$ GeV (solid line), $m_{\tilde{\nu}}=100$ GeV (dashed line), and $m_{\tilde{\nu}}=150$ GeV (dashed-dotted line). Also, $\tan\beta=35$, $M_s=100$ GeV, $m_{H^+}=200$ GeV, $m_{\tilde{l}_L}=m_{\tilde{\nu}}$ and $A'=M_s$ are used. See also [16] and [17].

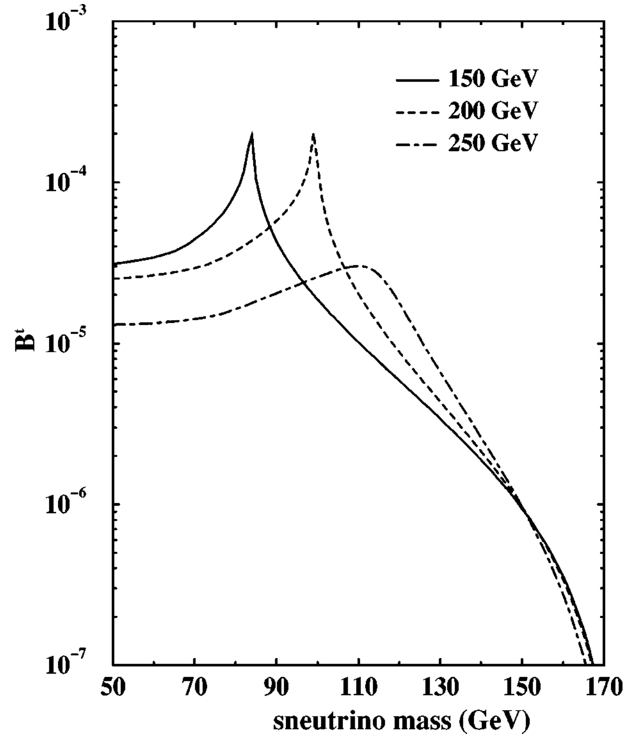

FIG. 4. The scaled branching ratio B^t as a function of the sneutrino mass $m_{\tilde{\nu}}$, for $\mu=150$ GeV (solid line), $\mu=200$ GeV (dashed line), and $\mu=250$ GeV (dashed-dotted line). See also caption to Fig. 3.

branching ratio B^t, we set $m_{\tilde{d}_R}=m_{\tilde{d}_L}=M_s$ for all squark flavors. Therefore, since M_s is our typical SUSY mass scale, it is only natural to set A'—the R_P trilinear soft breaking term in Eq. (2)—to be $A'=M_s$.

(3) The two physical chargino masses $m_{\tilde{\chi}_m}$ $(m=1,2)$ and the mixing matrices $Z^{\pm}$ are extracted by diagonalizing the chargino mass matrix which depends on the low-energy values of the Higgs mass parameter μ, the gaugino mass $\tilde{m}_2$, and $\tan\beta$ (for more details, see [13]). It is, however, sufficient to vary only one of the two mass parameters μ and $\tilde{m}_2$ in order to investigate the dependence of B^t on the chargino masses. We, therefore, set $\tilde{m}_2=85$ GeV. In the traditional grand unified theory (GUT) assumption, i.e., that there is an underlying grand unification, $\tilde{m}_2=85$ GeV corresponds to a gluino mass of ~ 300 GeV since, in that case, the gaugino masses are unified at the GUT scale leading to the relation, $\tilde{m}_2/m_{\text{gluino}}=\alpha/s_W^2\alpha_s$, at the electroweak scale (see, e.g., [13]).

(4) Since we are not interested here in *CP* violation, we take μ, $\tilde{m}_2$, and λ' to be real.

As it turns out, with a low $\tan\beta$, i.e., $\tan\beta \lesssim 10$, the branching ratio B^t is typically $\lesssim 10^{-6}$. Therefore, in Figs. 3 and 4 we focus on the high value $\tan\beta=35$. The dependence of B^t on $\tan\beta$ is shown in Fig. 5. Also, we find that $\Gamma(t \to c\tilde{\nu})$ drops as the squark mass scale M_s is increased, and in what follows, we present results for $M_s=100$ GeV. However, it is important to note that $B^t > 10^{-6}$ is still possible for squark masses $\lesssim 190$ GeV.

In Fig. 3 we plot the scaled branching ratio B^t, as a function of μ, for $\tan\beta=35$, $M_s=100$ GeV, and for three values of the sneutrino mass, $m_{\tilde{\nu}}=50,100,150$ GeV. Also, as stated above, here and throughout the rest of the paper we set $m_{H^+}=200$ GeV and $m_{\tilde{l}_L}=m_{\tilde{\nu}}$. We vary μ in the range $-400\,\text{GeV}<\mu<400\,\text{GeV}$ subject to $m_{\tilde{\chi}_m}>50$ GeV for $m=1$ or 2 (see also [16]). Evidently, in the range $-400\,\text{GeV}<\mu<400\,\text{GeV}$, for $m_{\tilde{\nu}}=50(100)$ GeV there is a $\sim 290(360)$ GeV range of μ in which $B^t>10^{-5}$. Note that B^t is largest for $m_{\tilde{\nu}}=100$ GeV [17], reaching $B^t>10^{-4}$ in the rather narrow μ mass ranges: $-185\,\text{GeV}\lesssim\mu\lesssim-165\,\text{GeV}$ and $200\,\text{GeV}\lesssim\mu\lesssim 210\,\text{GeV}$.

Figure 4 shows the dependence of B^t on the sneutrino mass, for $\mu=150$, 200, and 250 GeV. Here also, $\tan\beta=35$ and $M_s=100$ GeV. We see that $B^t>10^{-5}$ for $m_{\tilde{\nu}} \lesssim 110-125$ GeV depending on the value of μ. Again, B^t can reach above 10^{-4} in some small sneutrino mass ranges around ~ 85 and ~ 100 GeV (see also [17]).

Finally, in Fig. 5 we show the dependence of B^t on $\tan\beta$, for $m_{\tilde{\nu}}=50$, 100, and 150 GeV and for $\mu=200$ GeV. The rest of the parameters are fixed to the same values as in Figs. 3 and 4. As mentioned before, $B^t<10^{-6}$ for small $\tan\beta$ values of $\mathcal{O}(1)$ and it increases with $\tan\beta$. It is interesting to note that, for $m_{\tilde{\nu}}=100$ GeV and with $\tan\beta \approx m_t/m_b \sim 40$, B^t is well above 10^{-4} reaching almost 10^{-3} (see also [17]).

To conclude this section, noting that at the LHC about 10^7-10^8 top quarks will be produced, the observability of this flavor changing top quark decay will require at least $B^t \gtrsim 10^{-6}$. Therefore, the decay $t \to c\tilde{\nu}$ may be a very good

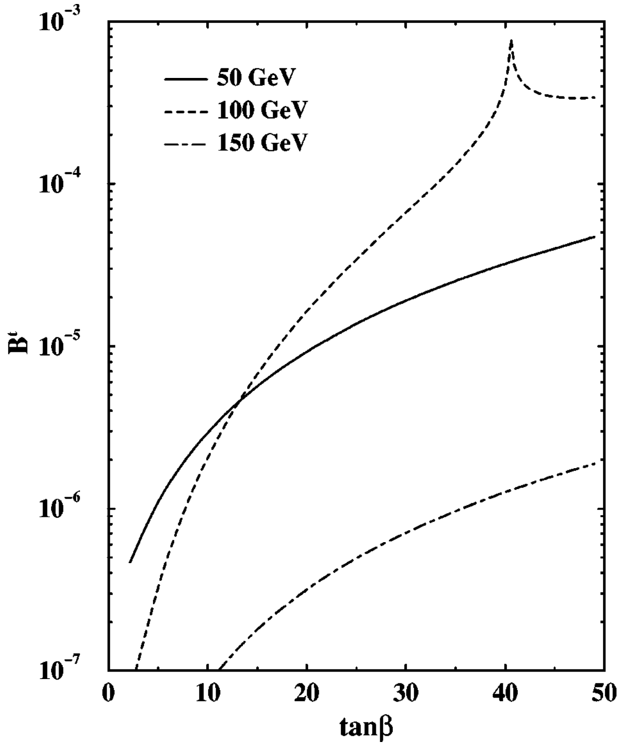

FIG. 5. The scaled branching ratio B^t as a function of $\tan\beta$, for $\mu=200$ GeV and for $m_{\tilde\nu}=50$ GeV (solid line), $m_{\tilde\nu}=100$ GeV (dashed line), and $m_{\tilde\nu}=150$ GeV (dashed-dotted line). The rest of the parameters are held fixed to their values in Figs. 3 and 4. See also the caption to Fig. 3.

venue to determine $\tan\beta$ in $R\!\!\!/_P$ SUSY models, since its branching ratio does not reach this limit if $\tan\beta$ is smaller than about 10. However, for a high $\tan\beta$ scenario, e.g., with $\tan\beta\approx m_t/m_b$, the decay $t\to c\tilde\nu$ may have a branching ratio well above 10^{-5} in some wide ranges of the SUSY parameters, provided that $m_{\tilde\nu}\lesssim 120$ GeV and that the squark masses are of $\mathcal{O}(100)$ GeV. Although not explicitly shown above, we find that $B^t\gtrsim 10^{-6}$ is possible as long as the squark masses are $\lesssim$190 GeV. Moreover, in some small ranges of the SUSY parameter space, $t\to c\tilde\nu$ can have a branching ratio above 10^{-4} reaching even 10^{-3}. It is interesting to note that, since the leading contribution to $\mathrm{Br}(t\to c\tilde\nu)$ is independent of m_c [see the term proportional to $|\overline{\mathcal{B}}^{\tilde\nu}|^2$ in Eq. (23)], $\mathrm{Br}(t\to u\tilde\nu)\simeq\mathrm{Br}(t\to c\tilde\nu)$ if $\lambda'_{i13}\sim\lambda'_{i23}\sim\mathcal{O}(1)$.

It is also useful to note that, with $\lambda'_{323}\sim\mathcal{O}(1)$, the τ-sneutrino will decay predominantly to a pair of $b\bar s$ with a branching ratio of $\mathcal{O}(1)$ (see Sec. III). Therefore, a good way for experimentally searching for this rare flavor changing top decay, i.e., $t\to c\tilde\nu$, may be to look for the three jets signature $t\to cb\bar s$, where the invariant $b\bar s$ mass reconstructs the sneutrino mass.

III. THE CASE $m_{\tilde\nu}>m_t$ AND THE SNEUTRINO DECAYS $\tilde\nu\to t\bar c,\bar t c$ AND $\bar{\tilde\nu}\to t\bar c,\bar t c$

In this section we discuss the opposite mass case, i.e., $m_{\tilde\nu}>m_t$, and thus the prospects for observing the ''reversed'' flavor-changing sneutrino decays

$$\tilde\nu_i\to t\bar c,\bar t c,\quad \bar{\tilde\nu}_i\to t\bar c,\bar t c, \tag{38}$$

at the LHC.

The calculation of the widths for the sneutrino decays in Eq. (38) is straightforward using our formulas for $t\to c\tilde\nu_i$ and $t\to c\bar{\tilde\nu}_i$ in the previous section. For $\tilde\nu_i\to\bar t c(\bar{\tilde\nu}_i\to\bar t c)$ the amplitude is the same as the amplitude for $t\to c\bar{\tilde\nu}_i(t\to c\tilde\nu_i)$ with the top spinor u_t replaced by the antitop spinor v_t. We, therefore, have

$$\Gamma(\tilde\nu_i\to\bar t c)=N_c\frac{|\mathcal{C}^{\bar{\tilde\nu}}|^2}{16\pi}\frac{\omega(m^2_{\tilde\nu_i},m_t^2,m_c^2)}{m^3_{\tilde\nu_i}}\frac{m_t^2m_c^2}{m_W^4}\times\Bigg\{(m^2_{\tilde\nu_i}-m_t^2-m_c^2)\left(\frac{m_W^2}{m_t^2}|\overline{\mathcal{A}}^{\bar{\tilde\nu}}|^2+\frac{m_W^2}{m_c^2}|\overline{\mathcal{B}}^{\bar{\tilde\nu}}|^2\right)-4m_W^2\,\mathrm{Re}[\overline{\mathcal{A}}^{\bar{\tilde\nu}}(\overline{\mathcal{B}}^{\bar{\tilde\nu}})^*]\Bigg\}, \tag{39}$$

where $N_c=3$ is the number of colors, $\omega(m^2_{\tilde\nu_i},m_t^2,m_c^2)$ is defined in Eq. (25), and $\overline{\mathcal{A}}^{\bar{\tilde\nu}}$, $\overline{\mathcal{B}}^{\bar{\tilde\nu}}$ are defined in Eq. (30).

For $\bar{\tilde\nu}_i\to\bar t c$ the width is given by Eq. (39) with the replacements: $\mathcal{C}^{\bar{\tilde\nu}}\to\mathcal{C}^{\tilde\nu}$ and $\overline{\mathcal{A}}^{\bar{\tilde\nu}}\to\overline{\mathcal{A}}^{\tilde\nu}$, $\overline{\mathcal{B}}^{\bar{\tilde\nu}}\to\overline{\mathcal{B}}^{\tilde\nu}$, where $\overline{\mathcal{A}}^{\tilde\nu}$ and $\overline{\mathcal{B}}^{\tilde\nu}$ are defined in Eq. (24).

In addition, if CP is a good symmetry (as is in our case since we have assumed that the λ''s, μ, and $\tilde m_2$ are all real), then

$$\Gamma(\tilde\nu_i\to t\bar c)=\Gamma(\bar{\tilde\nu}_i\to\bar t c), \tag{40}$$

$$\Gamma(\bar{\tilde\nu}_i\to t\bar c)=\Gamma(\tilde\nu_i\to\bar t c). \tag{41}$$

We are now interested in the branching ratio of a sneutrino, say again the τ-sneutrino $(i=3)$ and dropping the index i, to decay to a $t\bar c$ and $\bar t c$ pairs:

$$B^s\equiv\frac{\Gamma(\tilde\nu\to t\bar c)+\Gamma(\tilde\nu\to\bar t c)}{\Gamma^{\tilde\nu}_{\mathrm{tot}}}. \tag{42}$$

Due to CP invariance the corresponding branching ratio for $\bar{\tilde\nu}$ equals that of $\tilde\nu$, i.e., B^s defined above. It is also worth noting that, similar to the arguments given in the previous section, in our scenario with $\lambda'_{i32}=0$ and $\lambda'_{i23}\neq 0$, $\Gamma(\tilde\nu\to t\bar c)\gg\Gamma(\tilde\nu\to\bar t c)$ and $B^s\approx\Gamma(\tilde\nu\to t\bar c)/\Gamma^{\tilde\nu}_{\mathrm{tot}}$.

For our purpose we define the τ-sneutrino total width as

$$\Gamma^{\tilde\nu}_{\mathrm{tot}}=\Gamma(\tilde\nu\to b\bar s)+\Gamma(\tilde\nu\to\tilde\chi\tau)+\Gamma(\tilde\nu\to\tilde\chi^0\nu_\tau). \tag{43}$$

This should serve as an approximate expression for the $\tilde\nu$ total width since, as assumed in this analysis, with $m_{\tilde\nu_i}=m_{\tilde l_{Li}}$ and when the three sneutrino flavors are degenerate, the decays $\tilde\nu_i\to W^+\tilde l_{Li},H^+\tilde l_{Li},Z^0\tilde\nu_j,H^0\tilde\nu_j$ are kinematically forbidden.

The width of the $R\!\!\!/_P$ sneutrino decay to a $b\bar{s}$ pair is given by

$$\Gamma(\tilde{\nu}\to b\bar{s})=(\lambda'_{323})^2\frac{N_c}{16\pi}m_{\tilde{\nu}}, \tag{44}$$

and the widths of the R-parity conserving decays in Eq. (43) to a chargino ($\tilde{\chi}$) and to a neutralino ($\tilde{\chi}^0$) are [18]

$$\Gamma(\tilde{\nu}\to\tilde{\chi}\tau);\ \Gamma(\tilde{\nu}\to\tilde{\chi}^0\nu_\tau)=\mathcal{C}\frac{g^2}{16\pi}m_{\tilde{\nu}}\times[(1-m_{\tilde{\chi}}^2/m_{\tilde{\nu}}^2)^2;$$
$$(1-m_{\tilde{\chi}^0}^2/m_{\tilde{\nu}}^2)^2], \tag{45}$$

where $\mathcal{C}<1$, since it is proportional to the square of the charginos or neutralinos mixing matrices [18]. Therefore, $\Gamma(\tilde{\nu}\to\tilde{\chi}\tau);\Gamma(\tilde{\nu}\to\tilde{\chi}^0\tilde{\nu}_\tau)<10^{-2}m_{\tilde{\nu}}$ and, for simplicity, in our numerical results we conservatively ignore the additional phase-space factors in Eq. (45), by setting $\Gamma(\tilde{\nu}\to\tilde{\chi}\tau)=\Gamma(\tilde{\nu}\to\tilde{\chi}^0\tilde{\nu}_\tau)=10^{-2}m^{\tilde{\nu}}$. Evidently, with $\lambda'_{323}\sim\mathcal{O}(1)$, $\Gamma_{\rm tot}^{\tilde{\nu}}$ is dominated by the width of the $R\!\!\!/_P$ decay $\Gamma(\tilde{\nu}\to b\bar{s})\sim 6\times10^{-2}m_{\tilde{\nu}}$.

Let us now discuss our numerical results for the branching ratio B^s, defined in Eq. (42). As in the previous section, we fix $m_{H^+}=200$ GeV, $m_{\tilde{l}_L}=m_{\tilde{\nu}}$ (B^s is also found to be practically insensitive to these two parameters), $\tilde{m}_2=85$ GeV, and we study the dependence of B^s on the parameters $m_{\tilde{\nu}}$, M_s (again setting $A'=M_s$), μ, and $\tan\beta$.

In Fig. 6 we plot the branching ratio B^s as a function of μ for $\lambda'_{i23}=1$, $\tan\beta=35$, $M_s=100$ GeV, and for three values of the sneutrino mass, $m_{\tilde{\nu}}=200$, 250, and 300 GeV. We see that for $m_{\tilde{\nu}}=200$–300 GeV, $B^s>10^{-6}$ as long as $|\mu|\lesssim300$ GeV. Moreover, $B^s>10^{-5}$ in the ranges $-205\text{ GeV}\lesssim\mu\lesssim-160$ GeV and $185\text{ GeV}\lesssim\mu\lesssim240$ GeV.

Although, as in the case of the top decay $t\to c\tilde{\nu}$, we find that in general B^s decreases as the squark mass M_s increases, for $\tilde{\nu}\to t\bar{c},\bar{t}c$ the dependence is not as trivial as in $t\to c\tilde{\nu}$. We, therefore, plot in Fig. 7 B^s as a function of the squarks mass M_s, for $\tan\beta=35$, $\mu=200$ GeV, and $m_{\tilde{\nu}}=200$, 250 or 300 GeV. Evidently, while for $m_{\tilde{\nu}}=200$ GeV the branching ratio B^s is above 10^{-7} only if $M_s\lesssim130$ GeV, for $m_{\tilde{\nu}}=250$ GeV, and $m_{\tilde{\nu}}=300$ GeV, B^s can be $>10^{-7}$ for larger M_s values: $M_s\lesssim180$ GeV and $M_s\lesssim160$ GeV respectively.

In Fig. 8 we show the dependence of B^s, on $\tan\beta$, for the three sneutrino masses $m_{\tilde{\nu}}=200$, 250, and 300 GeV and with $\lambda'_{i23}=1$, $M_s=100$ GeV and $\mu=200$ GeV. The same behavior as in the case of $t\to c\tilde{\nu}$ is found. In particular, $B^s<10^{-7}$ for $\tan\beta\sim\mathcal{O}(1)$ and it increases with $\tan\beta$ such that $B^s>10^{-7}$ for $\tan\beta\gtrsim10$. For $m_{\tilde{\nu}}=200$ GeV and $\tan\beta\gtrsim30$, $B^s>10^{-5}$.

Note that, as in the case of the top decays and for the same reasons, the results above apply also to $\text{Br}(\tilde{\nu}\to t\bar{u})$ if instead $\lambda'_{i13}\sim\mathcal{O}(1)$.

III. SUMMARY AND CONCLUSIONS

We have calculated the branching ratios for the flavor-changing top quark decays $t\to c\tilde{\nu}$, $c\bar{\tilde{\nu}}$ when $m_{\tilde{\nu}}<m_t$, and for

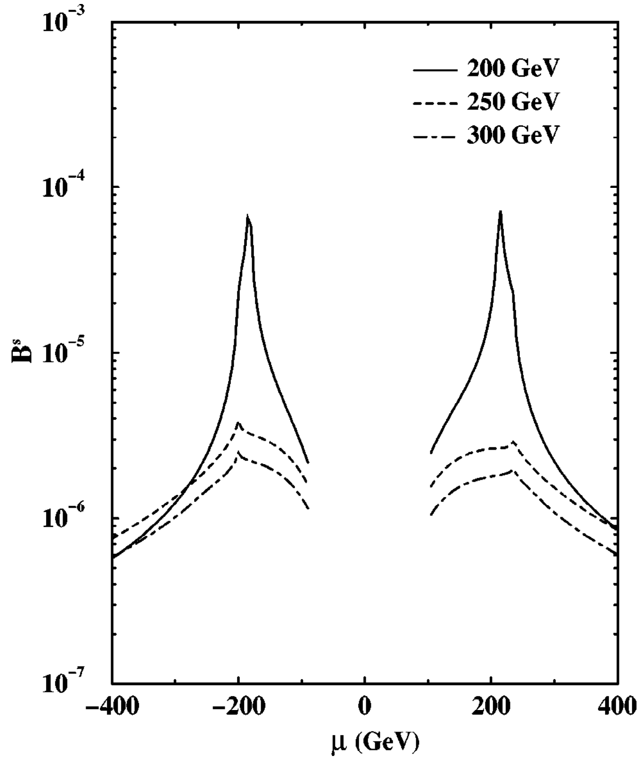

FIG. 6. The branching ratio B^s, defined in Eq. (42), as a function of μ, for: $m_{\tilde{\nu}}=200$ GeV (solid line), $m_{\tilde{\nu}}=250$ GeV (dashed line), and $m_{\tilde{\nu}}=300$ GeV (dashed-dotted line). Also, $\lambda'_{i23}=1$ for a given seutrino flavor i and the rest of the parameters are set to $\tan\beta=35$, $M_s=100$ GeV, $m_{H^+}=200$ GeV, $m_{\tilde{l}_L}=m_{\tilde{\nu}}$ and $A'=M_s$. See also [16].

the corresponding reversed sneutrino decays $\tilde{\nu}\to t\bar{c}$, $\bar{t}c$, and $\bar{\tilde{\nu}}\to t\bar{c}$, $\bar{t}c$ when $m_{\tilde{\nu}}>m_t$, in the MSSM with R-parity violation. In this model, these decays can occur at the one-loop level and they depend predominantly on the squark masses, the Higgs mass parameter μ, $\tan\beta$ and for a sneutrino flavor $i=e$, μ or τ, on the $R\!\!\!/_P$ couplings λ'_{i23}.

We have considered the values $\lambda'_{i23}\sim\mathcal{O}(1)$, which are not ruled out if one allows more than one $R\!\!\!/_P$ coupling to be nonzero, and showed that these rare decays are sensitive probes of the parameter $\tan\beta$, since their branching ratios become experimentally accessible (i.e., at the LHC), since typically, $\text{Br}(t\to c\tilde{\nu})\gtrsim10^{-6}$ and $\text{Br}(\tilde{\nu}\to t\bar{c})\gtrsim10^{-7}$ only for $\tan\beta\gtrsim10$.

For the top decays in the case $m_{\tilde{\nu}}<m_t$, we found that $\Gamma(t\to c\tilde{\nu})\gg\Gamma(t\to c\bar{\tilde{\nu}})$ and that $\text{Br}(t\to c\tilde{\nu})$ can be well within the reach of the LHC with 10^7-10^8 $t\bar{t}$ pairs produced. In particular, it was shown that $\text{Br}(t\to c\tilde{\nu})>10^{-5}$ for $\tan\beta\gtrsim30$ in a $\gtrsim300$ GeV range of the Higgs mass parameter μ, as long as the squark masses are not much larger than 100 GeV. In some cases, with $\tan\beta\sim m_t/m_b\sim40$, we found that $\text{Br}(t\to c\tilde{\nu})>10^{-4}$ reaching almost 10^{-3}.

For the reversed sneutrino decays in the case $m_{\tilde{\nu}}>m_t$, we found that similar to the above-mentioned t decays, $\Gamma(\tilde{\nu}\to t\bar{c})\gg\Gamma(\tilde{\nu}\to\bar{t}c)$. Furthermore, $\text{Br}(\tilde{\nu}\to t\bar{c})\gtrsim10^{-6}$ for $|\mu|\lesssim300$ GeV, $\tan\beta\gtrsim30$, again, as long as the squarks have masses around 100 GeV. Here also the branching ratio can

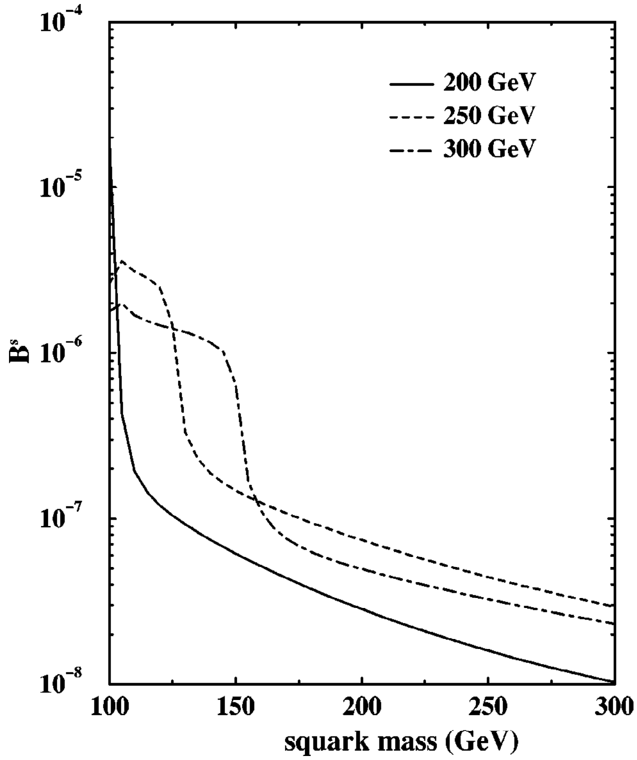

FIG. 7. The branching ratio B^s, as a function of the squark mass M_s, for $\mu=200$ GeV and for $m_{\tilde{\nu}}=200$ GeV (solid line), $m_{\tilde{\nu}}=250$ GeV (dashed line), and $m_{\tilde{\nu}}=300$ GeV (dashed-dotted line). See also caption to Fig. 6.

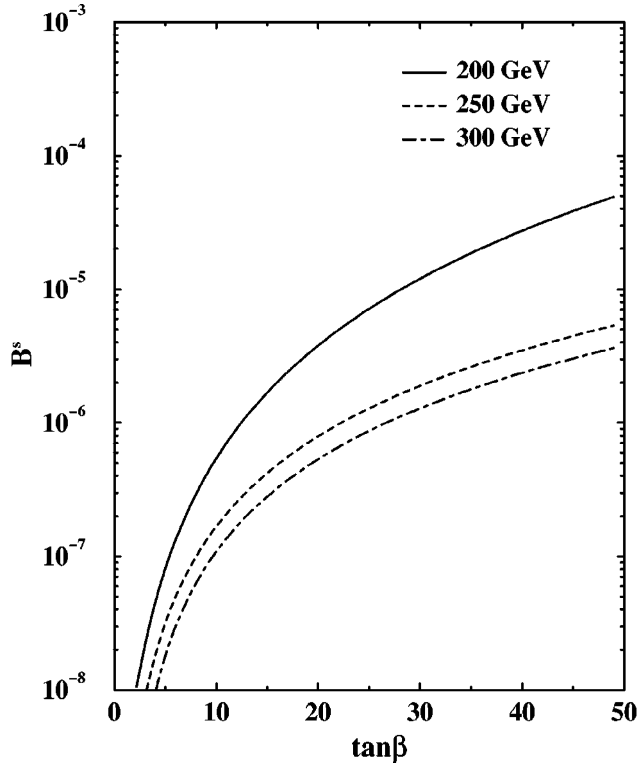

FIG. 8. The branching ratio B^s as a funciton of $\tan\beta$, for $\mu=200$ GeV and for $m_{\tilde{\nu}}=200$ GeV (solid line), $m_\nu=250$ GeV (dashed line), and $m_{\tilde{\nu}}=300$GeV (dashed-dotted line). See also caption to Fig. 6.

be more than an order of magnitude larger, e.g., for $\tan\beta\sim m_t/m_b\sim 40$ and squark masses around 100 GeV, $\mathrm{Br}(\tilde{\nu}\to t\bar{c})\sim 10^{-5}-10^{-4}$ is possible. As mentioned in the introduction, the LHC will be able to produce 10^8-10^9 sneutrinos with a mass of 200–300 GeV, if indeed some of the $R\!\!\!/_P$ couplings of the λ' type are saturated to be of $\mathcal{O}(1)$. Therefore, if $t\to c\tilde{\nu}$ is not detected at the LHC, then our results above indicate that it may still be useful to search for the reversed sneutrino decay $\tilde{\nu}\to t\bar{c}$ via the reaction $pp\to\tilde{\nu}+X\to t\bar{c}+X$.

Finally, changing the charm quark to an up quark in these decays has a negligible effect on the branching ratios since, by assumption, the couplings are flavor independent. Therefore, if in addition to $\lambda'_{i23}\sim\mathcal{O}(1)$ also $\lambda'_{i13}\sim\mathcal{O}(1)$, then the same results are obtained for the decays $t\to u\tilde{\nu}$, $u\bar{\tilde{\nu}}$ and $\tilde{\nu}\to t\bar{u},\bar{t}u$.

ACKNOWLEDGMENTS

We acknowledge partial support from the US Israel BSF (G.E. and A.S.) and from the US DOE Contracts No. DE-AC02-76CH00016(BNL) (A.S.), DE-FG03-94ER40837(UCR) (S.B.). G.E. also thanks the VPR Fund at the Technion for partial support.

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[17] We want to comment here on the sharp peaks in some of the Figs. 3–8. The loop integrals associated with the diagrams in Fig. 1 are very sensitive to the masses of the particles in the loop such that an enhancement occurs whenever an absorptive threshold is open (i.e., when two particles in the loop can go on their mass shell). Thus, for example, if $m_{\tilde{\nu}} \gtrsim m_{\tilde{\chi}} + m_l$, such a threshold is satisfied in diagram 4 in Fig. 1, in which case it causes the branching ratio B^t to be larger for $m_{\tilde{\nu}} = 100$ GeV compared to $m_{\tilde{\nu}} = 50$ GeV, in spite of phase-space suppression.

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