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Light Stop: MSSM versus R–parity violation

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Abstract

We discuss the phenomenology of the lightest stops in models where R-parity is broken by bilinear terms. In this class of models we consider scenarios where the R-parity breaking two-body decay $\tilde{t}_1 \to \tau^+ b$ competes with the leading three-body decays $\tilde{t}_1 \to W^+ b\,\tilde{\chi}^0_1$, $\tilde{t}_1 \to H^+ b\,\tilde{\chi}^0_1$, $\tilde{t}_1 \to b\,\tilde{l}^+_i\,\nu_l$, and $\tilde{t}_1 \to b\,\tilde{\nu}_l\,l^+$ ($l = e, \mu, \tau$). We demonstrate that the R–parity violating decay can be the dominant one. In particular we focus on the implications for a future e^+e^- Linear Collider.

1 Introduction

The search for supersymmetry (SUSY) [1, 2] plays an important rôle in the experimental program at the colliders LEP2 and Tevatron. It will be even more important at future colliders, e.g. an upgraded Tevatron, LHC, an e^+e^- linear collider. Therefore many phenomenological studies have been carried out in recent years (see e.g. [3, 4, 5, 6] and references therein). Most of them have been carried out in the context of the minimal supersymmetric standard model (MSSM) [2, 7]. However, neither gauge invariance nor supersymmetry requires the conservation of R-parity. Indeed, there is considerable theoretical and phenomenological interest in studying possible implications of alternative scenarios [8] in which R-parity is broken [9, 10, 11, 12]. The violation of R-parity could arise explicitly [13] as a residual effect of some larger unified theory [10], or spontaneously, through nonzero vacuum expectation values (vev's) for scalar neutrinos [9, 11, 12]. In realistic spontaneous R-parity breaking models there is an $SU(2) \otimes U(1)$ singlet sneutrino vev characterizing the scale of R-parity violation [14, 15, 16, 17] which is expected to be the same as the effective supersymmetry breaking scale.

There are two generic cases of spontaneous R-parity breaking models to consider. In the absence of any additional gauge symmetry, these models lead to the existence of a physical massless Nambu-Goldstone boson, called majoron (J) which is *the lightest SUSY* *particle*, massless and therefore stable. If lepton number is part of the gauge symmetry and R-parity is spontaneously broken then there is an additional gauge boson which gets mass via the Higgs mechanism, and there is no physical Goldstone boson [17]. As in the standard case in R-parity breaking models the lightest SUSY particle (LSP) is in general a neutralino. However, it now decays mostly into visible states, therefore diluting the missing momentum signal and bringing in increased multiplicity events which arise mainly from three-body decays such as

$$\tilde{\chi}^0_1 \to f \bar{f} \nu, \tag{1}$$

where f denotes a charged fermion. The neutralino also has the invisible decay mode

$$\tilde{\chi}^0_1 \to 3\nu. \tag{2}$$

as well as

$$\tilde{\chi}^0_1 \to \nu J, \tag{3}$$

in the case the breaking of R-parity is spontaneous [14, 15]. This last decay conserves R-parity since the majoron has a large R-odd singlet sneutrino component.

Owing to the large top Yukawa coupling the stops have a quite different phenomenology compared to those of the first two generations of up–type squarks (see e.g. [18] and references therein). The large Yukawa coupling implies a large mixing between $\tilde{t}_L$ and $\tilde{t}_R$ [19] and large couplings to the higgsino components of neutralinos and charginos. The large top quark mass also implies the existence of scenarios where all two-body decay modes of $\tilde{t}_1$ (e.g. $\tilde{t}_1 \to t\,\tilde{\chi}^0_i, b\,\tilde{\chi}^+_j, t\,\tilde{g}$) are kinematically forbidden. In these scenarios higher order decays of $\tilde{t}_1$ become relevant: [20, 21]: $\tilde{t}_1 \to c\,\tilde{\chi}^0_{1,2}$, $\tilde{t}_1 \to W^+ b\,\tilde{\chi}^0_1$, $\tilde{t}_1 \to H^+ b\,\tilde{\chi}^0_1$, $\tilde{t}_1 \to b\,\tilde{l}^+_i\,\nu_l$, $\tilde{t}_1 \to b\,\tilde{\nu}_l\,l^+$, where l denotes e, μ, τ. In [21] it has been shown that the three-body decay modes are in general much more important than the two body decay mode in the framework of the MSSM. Recently it has been demonstrated that not only LSP decays are sign of R-parity violation but that also the light stop is possible candidate for observing R-parity violation even if R-parity violation is small [6, 22, 23]. It has been demonstrated that there exists a large parameter region where the R-parity violating decay

$$\tilde{t}_1 \to b\,\tau \tag{4}$$

is much more important than

$$\tilde{t}_1 \to c\,\tilde{\chi}^0_{1,2} \tag{5}$$

in scenarios where only those decay modes are possible. It is therefore natural to ask if there exist scenarios where the decay $\tilde{t}_1 \to b\,\tau$ is as important as the three–body decays. Note that in the R-parity violating models under consideration the neutral (charged) Higgs–bosons mix with the neutral (charged) sleptons. These states are denoted by S^0_i, P^0_j, and $S^\pm_k$ for the neutral scalars, pseudoscalars and charged scalars, respectively. Therefore in the R-parity violating case one has the following three-body decay modes:

$$\begin{aligned}
\tilde{t}_1 &\to W^+ b\,\tilde{\chi}^0_1 && (6)\\
\tilde{t}_1 &\to S^+_k\, b\,\tilde{\chi}^0_1 && (7)\\
\tilde{t}_1 &\to S^+_k\, b\,\nu_l && (8)\\
\tilde{t}_1 &\to b\,S^0_i\,l^+\,, && (9)\\
\tilde{t}_1 &\to b\,P^0_j\,l^+\,. && (10)
\end{aligned}$$

We will demonstrate that $\tilde{t}_1 \to b\,\tau^+$ can indeed be the most important decay mode. In particular we will consider a mass range of $\tilde{t}_1$, where it is difficult for the LHC to discover the light stop within the MSSM due to the large top background [24]. The rest of paper is organized in the following way: in the next section we will introduce the model. In Sect. 3 numerical results for stop decays are presented and their implications for LC. In Sect. 4 we present our conclusions.

2 The model

The supersymmetric Lagrangian is specified by the superpotential W given by

$$W = \varepsilon_{ab}\left[h_U^{ij}\widehat{Q}_i^a\widehat{U}_j\widehat{H}_2^b + h_D^{ij}\widehat{Q}_i^b\widehat{D}_j\widehat{H}_1^a + h_E^{ij}\widehat{L}_i^b\widehat{R}_j\widehat{H}_1^a - \mu\widehat{H}_1^a\widehat{H}_2^b\right] + \varepsilon_{ab}\epsilon_i\widehat{L}_i^a\widehat{H}_2^b\,, \tag{11}$$

where $i,j = 1,2,3$ are generation indices, $a,b = 1,2$ are $SU(2)$ indices, and ε is a completely antisymmetric 2×2 matrix, with $\varepsilon_{12} = 1$. The symbol "hat" over each letter indicates a superfield, with $\widehat{Q}_i$, $\widehat{L}_i$, $\widehat{H}_1$, and $\widehat{H}_2$ being $SU(2)$ doublets with hypercharges $1/3$, -1, -1, and 1 respectively, and $\widehat{U}$, $\widehat{D}$, and $\widehat{R}$ being $SU(2)$ singlets with hypercharges $-\frac{4}{3}$, $\frac{2}{3}$, and 2 respectively. The couplings h_U, h_D and h_E are 3×3 Yukawa matrices, and μ and ϵ_i are parameters with units of mass.

Supersymmetry breaking is parametrized by the standard set of soft supersymmetry breaking terms

$$\begin{aligned}
V_{soft} \;=\;& M_Q^{ij2}\widetilde{Q}_i^{a*}\widetilde{Q}_j^a + M_U^{ij2}\widetilde{U}_i^*\widetilde{U}_j + M_D^{ij2}\widetilde{D}_i^*\widetilde{D}_j + M_L^{ij2}\widetilde{L}_i^{a*}\widetilde{L}_j^a + M_R^{ij2}\widetilde{R}_i^*\widetilde{R}_j \\
&+m_{H_1}^2 H_1^{a*}H_1^a + m_{H_2}^2 H_2^{a*}H_2^a \\
&-\left[\tfrac{1}{2}M_3\lambda_3\lambda_3 + \tfrac{1}{2}M\lambda_2\lambda_2 + \tfrac{1}{2}M'\lambda_1\lambda_1 + h.c.\right] \\
&+\varepsilon_{ab}\left[A_U^{ij}h_U^{ij}\widetilde{Q}_i^a\widetilde{U}_jH_2^b + A_D^{ij}h_D^{ij}\widetilde{Q}_i^b\widetilde{D}_jH_1^a + A_E^{ij}h_E^{ij}\widetilde{L}_i^b\widetilde{R}_jH_1^a\right. \\
&\left.-B\mu H_1^aH_2^b + B_i\epsilon_i\widetilde{L}_i^aH_2^b\right]\,,
\end{aligned} \tag{12}$$

Note that, in the presence of soft supersymmetry breaking terms the bilinear terms ϵ_i can not be rotated away, since the rotation that eliminates it reintroduces an R–Parity violating trilinear term, as well as a sneutrino vacuum expectation value [25].

For our discussion it suffices to assume R-parity Violation (RPV) only in the third generation. However we do allow for R-parity-conserving Flavour Changing Neutral Currents (FCNC) effects, such as the process $\tilde{t}_1 \to c\,\tilde{\chi}_1^0$ involving the three generations of quarks. In this case we will omit the labels i,j in the soft breaking terms. In order to study the R–Parity violating decay mode $\tilde{t}_1 \to b\,\tau$ it is sufficient to consider the superpotential [25, 26, 27, 28]

$$W = h_t\widehat{Q}_3\widehat{U}_3\widehat{H}_2 + h_b\widehat{Q}_3\widehat{D}_3\widehat{H}_1 + h_\tau\widehat{L}_3\widehat{R}_3\widehat{H}_1 - \mu\widehat{H}_1\widehat{H}_2 + \epsilon_3\widehat{L}_3\widehat{H}_2 \tag{13}$$

This amounts to neglecting the effects of RPV on the two first families. A short discussion on $\tilde{t}_1 \to b\,l^+$ in the three generation model will be given at the end of Sect. 3.

The bilinear term in Eq. (13) leads to a mixing between the charginos and the τ–lepton which in turn leads to the decay $\tilde{t}_1 \to b\,\tau$. The mass matrix is given by

$$\mathbf{M_C} = \begin{bmatrix} M & \frac{1}{\sqrt{2}}gv_2 & 0 \\ \frac{1}{\sqrt{2}}gv_d & \mu & -\frac{1}{\sqrt{2}}h_\tau v_3 \\ \frac{1}{\sqrt{2}}gv_3 & -\epsilon_3 & \frac{1}{\sqrt{2}}h_\tau v_d \end{bmatrix} \tag{14}$$

As in the MSSM, the chargino mass matrix is diagonalized by two rotation matrices $\mathbf{U}$ and $\mathbf{V}$

$$\mathbf{U}^*\mathbf{M_C}\mathbf{V}^{-1} = \begin{bmatrix} m_{\tilde\chi_1^\pm} & 0 & 0 \\ 0 & m_{\tilde\chi_2^\pm} & 0 \\ 0 & 0 & m_\tau \end{bmatrix} . \tag{15}$$

The lightest eigenstate of this mass matrix must be the tau lepton ($\tau^\pm$) and so the mass is constrained to be 1.7771 GeV. To obtain this the tau Yukawa coupling becomes a function of the parameters in the mass matrix, and the full expression is given in [29].

The stop mass matrix is given by

$$M^2_{\tilde t} = \begin{bmatrix} M_Q^2 + \frac{1}{2}v_u^2 {h_t}^2 + \Delta_{UL} & \frac{h_U}{\sqrt{2}}(v_u A_t - \mu v_d + \epsilon_3 v_3) \\ \frac{h_U}{\sqrt{2}}(v_u A_t - \mu v_d + \epsilon_3 v_3) & M_U^2 + \frac{1}{2}v_u^2 {h_t}^2 + \Delta_{UR} \end{bmatrix} \tag{16}$$

with $\Delta_{UL} = \frac{1}{8}(g^2 - \frac{1}{3}g'^2)(v_d^2 - v_u^2 + v_3^2)$ and $\Delta_{UR} = \frac{1}{6}g'^2(v_d^2 - v_u^2 + v_3^2)$. The sum of the v_i^2 is given by $m_W^2 = g^2(v_d^2 + v_u^2 + v_3^2)/2$. The mass eigenstates are given by $\widetilde t_1 = \widetilde t_L \cos\theta_{\tilde t} + \widetilde t_R \sin\theta_{\tilde t}$ and $\widetilde t_2 = \widetilde t_R \cos\theta_{\tilde t} - \widetilde t_L \sin\theta_{\tilde t}$. The sfermion mixing angle is given by

$$\cos\theta_{\tilde t} = \frac{-M^2_{\tilde t_{12}}}{\sqrt{(M^2_{\tilde t_{11}} - m^2_{\tilde t_1})^2 + (M^2_{\tilde t_{12}})^2}}, \qquad \sin\theta_{\tilde t} = \frac{M^2_{\tilde t_{11}} - m^2_{\tilde t_1}}{\sqrt{(M^2_{\tilde t_{11}} - m^2_{\tilde t_1})^2 + (M^2_{\tilde t_{12}})^2}} . \tag{17}$$

In addition the charged Higgs bosons mix with charged sleptons and the real (imaginary) parts of the sneutrino mix the scalar (pseudoscalar) Higgs bosons. The formulas can be found e.g. in [29, 30]. Their main decay modes for the mass range considered in this study are:

$$S_i^0 \;\to\; b\,\bar b\,,\; \tau^+\tau^-\,,\; \tilde\chi_1^0\nu_\tau \tag{18}$$

$$P_j^0 \;\to\; b\,\bar b\,,\; \tau^+\tau^-\,,\; \tilde\chi_1^0\nu_\tau \tag{19}$$

$$S_k^- \;\to\; s\,\bar c\,,\; \tau^-\nu_\tau\,,\; \tilde\chi_1^0\tau^- \tag{20}$$

3 Numerical results

In this section we present our numerical results for the branching ratios of the higher order decays of $\tilde t_1$. Here we consider scenarios where all two-body decays induced at tree-level are kinematically forbidden. Before going into detail it is useful to have some approximate formulas at hand [23]:

$$\Gamma(\tilde t_1 \to b\,\tau) \;\approx\; \frac{g^2|U_{32}|^2 h_b^2 \cos^2_{\theta_{\tilde t}} m_{\tilde t_1}}{16\pi} \tag{21}$$

$$\Gamma(\tilde t_1 \to c\,\tilde\chi_1^0) \;\approx\; 10^{-6} h_b^4 m_{\tilde t_1} \tag{22}$$

where $|U_{32}| \approx |\epsilon_3/\mu|$ if $|\epsilon_3| \ll |\mu|$ and $v_3 \ll m_W$. The complete formulas are given in [22, 23]. For the three–body decays the formulas given in [21] can be used as a good approximation if the mixings induced by R-parity violation are small. The complete formulas for the three–body decays in the R–parity violating case will be given elsewhere [30].

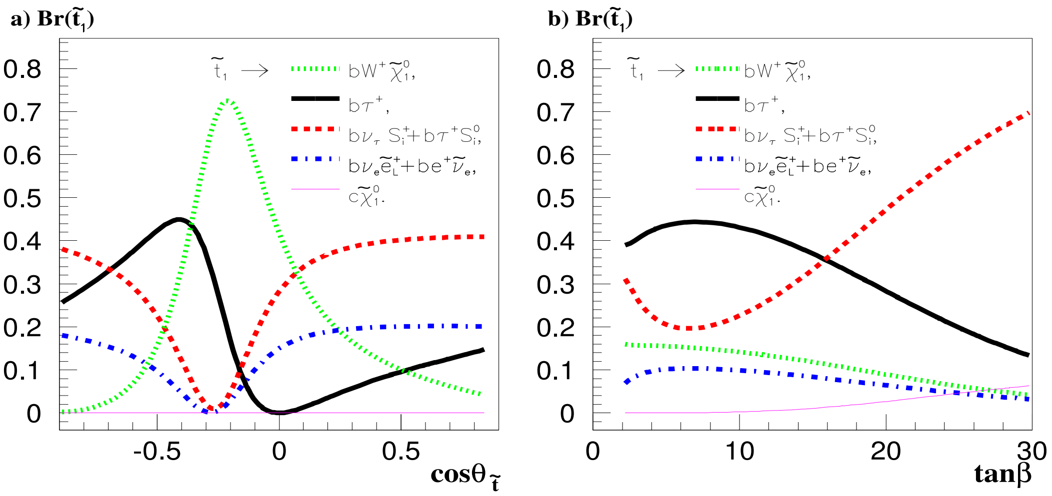

Figure 1: Branching ratios for $\tilde{t}_1$ decays for $m_{\tilde{t}_1} = 220$ GeV, $\mu = 500$ GeV, $M = 240$ GeV, and $m_\nu = 100$ eV. In a) the branching ratios are shown as a function of $\cos\theta_{\tilde{t}}$ for $\tan\beta = 4$, in b) as a function of $\tan\beta$ for $\cos\theta_{\tilde{t}} = 0.25$. Note, that the graph - - displays the sum $b\,\nu_\tau\, S_i^+ + b\,\tau^+\, S_i^0 + b\,\tau^+\, P_i^0$.

Input:	$\tan\beta = 4$	$\mu = 500$ GeV	$M = 240$ GeV
	$M_{\tilde{D}} = 370$ GeV	$M_{\tilde{Q}} = 340$ GeV	$A_b = 150$ GeV
	$M_{\tilde{E}} = 190$ GeV	$M_{\tilde{L}} = 190$ GeV	$A_\tau = 150$ GeV
	$m_{\tilde{t}_1} = 220$ GeV	$\cos\theta_{\tilde{t}} = 0.25$	$m_{P_2^0} = 110$ GeV
Calculated	$m_{\tilde{\chi}_1^0} = 120$	$m_{\tilde{\chi}_1^+} = 225$	$m_{\tilde{\chi}_2^+} = 520$
	$m_{\tilde{b}_1} = 340$ GeV	$m_{\tilde{b}_2} = 375$ GeV	$\cos\theta_{\tilde{b}} = 0.925$
	$m_{S_1^0} = 82$ GeV	$m_{S_2^0} = 128$ GeV	$m_{S_3^0} = 182$ GeV
	$m_{P_2^0} = 110$ GeV	$m_{P_3^0} = 182$ GeV	
	$m_{S_2^-} = 136$ GeV	$m_{S_3^-} = 187$ GeV	$m_{S_4^-} = 204$ GeV
	$m_{\tilde{e}_L} = 213$ GeV	$m_{\tilde{\nu}_e} = m_{\tilde{\nu}_\tau} = 204$ GeV	

Table 1: Input parameters and resulting quantities used in Fig. 1.

We have fixed the parameters as in [21] to avoid colour breaking minima: we have used $m_{\tilde{t}_1}$, $\cos\theta_{\tilde{t}}$, $\tan\beta$, and μ as input parameters in the top squark sector. For the sbottom (stau) sector we have fixed $M_{\tilde{Q}}$, $M_{\tilde{D}}$ and A_b ($M_{\tilde{E}}$, $M_{\tilde{L}}$, and A_τ) as input parameters. In addition we choose the R-parity violating parameters ϵ_3 and v_3 in such a way that the tau neutrino mass is fixed ([23] and references therein):

$$m_{\nu_\tau} \approx -\frac{(g^2 M' + g'^2 M)\mu'^2}{4MM'\mu'^2 - 2(g^2M' + g'^2M)\mu' v_u v_d' \cos\xi} {v_d'}^2 \sin^2\xi \qquad (23)$$

with

$$\sin\xi = \frac{\epsilon_3 v_d + \mu v_3}{\sqrt{\mu^2+\epsilon_3^2}\sqrt{v_d^2+v_3^2}} \qquad (24)$$

$$\mu' = \sqrt{\mu^2+\epsilon_3^2}\,, \qquad v_d' = \sqrt{v_d^2+v_3^2}\,. \qquad (25)$$

For simplicity, we assume that the soft SUSY breaking parameters are equal for all generations.

In Fig. 1(a) and (b) we show the branching ratios of $\tilde{t}_1$ as a function of $\cos\theta_{\tilde{t}}$. The parameters and physical quantities are given in Tab. 1. In Fig. 1(a) we show BR($\tilde{t}_1 \rightarrow$

$b\,\tau^+$), BR($\tilde{t}_1 \to c\,\tilde{\chi}^0_1$), BR($\tilde{t}_1 \to b\,W^+\,\tilde{\chi}^0_1$), BR($\tilde{t}_1 \to b\,e^+\,\tilde{\nu}_e$) + BR($\tilde{t}_1 \to b\,\nu_e\,\tilde{e}^+_L$). The branching ratios for decays into $\tilde{\mu}_L$ or $\tilde{\nu}_\mu$ are practically the same as those into $\tilde{e}_L$ or $\tilde{\nu}_e$. We have summed up those branching ratios for the decays into sleptons that give the same final state, for example:

$$\tilde{t}_1 \to b\,\nu_e\,\tilde{e}^+_L \;\to\; b\,e^+\,\nu_e\,\tilde{\chi}^0_1\;, \quad \tilde{t}_1 \to b\,e^+\,\tilde{\nu}_e \;\to\; b\,e^+\,\nu_e\,\tilde{\chi}^0_1 \tag{26}$$

Note that in Fig. 1 we have also summed the decay branching ratios BR($\tilde{t}_1 \to b\,S^+_k\,\nu_\tau$) + BR($\tilde{t}_1 \to b\,\tau^+\,S^0_i$) + BR($\tilde{t}_1 \to b\,\tau^+\,P^0_j$).

In the above cases the assumption $m_{\tilde{t}_1} - m_b < m_{\tilde{\chi}^+_1}$ implies $m_{\tilde{\chi}^+_1} > m_{\tilde{l}}$. Therefore, charginos can not arise as decay products of sleptons. The latter can only decay into the corresponding lepton plus $\tilde{\chi}^0_1$ except for a small parameter region where the decay into $\tilde{\chi}^0_2$ is possible. However, this decay is negligible due to kinematics in that region. In addition there exists the possibility of R-parity violating decays. However, these will be small because the neutrinos mix mainly with higgsinos implying that the partial decay widths are proportional to the squared product of an R-parity violating mixing parameter and small Yukawa coupling. For this set of parameters BR($\tilde{t}_1 \to c\,\tilde{\chi}^0_1$) is $O(10^{-4})$ independent of $\cos\theta_{\tilde{t}}$ and therefore negligible. Near $\cos\theta_{\tilde{t}} = -0.3$ one has $\tilde{t}_1 \to b\,W^+\,\tilde{\chi}^0_1$ as dominant decay channel, since the $\tilde{t}_1$-$\tilde{\chi}^+_1$-b coupling vanishes implying that the main contribution for the decays into the scalars vanishes. Moreover, the width for $\tilde{t}_1 \to b\,\tau^+$ is somewhat suppressed because of the $\cos^2\theta_{\tilde{t}}$ factor in Eq. (21). Note, from the figure that the branching ratios for the various decays into selectrons e-sneutrino is roughly a factor two smaller than the sum of the decays into $S^\pm_k$, S^0_i, P^0_j. The reason is that, for this choice of parameters the P^0_2 is mainly the pseudoscalar Higgs boson A^0 of the MSSM with mass 110 GeV. In this case the R-parity violating channel $\tilde{t}_1 \to b\,\tau^+\,P^0_2$ is comparable to the corresponding R-parity conserving decays. This state appears additional to the states which carry tau–lepton number in the MSSM limit giving rise to the observed difference. Note, that one has to expect additional jets from the states containing the scalars S^0_i, P^0_j, and $S^\pm_k$ because they have admixtures of the original Higgs boson. In case of negative $\cos\theta_{\tilde{t}}$ the decay into $\tilde{t}_1 \to b\,\tau^+$ is important and can be even the most important one. Therefore one has events with $\tau^+\,\tau^-\,b\,\bar{b}$ in the final state which can be used for a full mass reconstruction of the light stop.

In Fig. 1(b) the $\tan\beta$ dependence of branching ratios is shown. For this specific choice of $\cos\theta_{\tilde{t}}$ the decay $\tilde{t}_1 \to b\,\tau^+$ is the most important one for $\tan\beta \lesssim 15$. Above this value the final states which contain the scalars corresponding to the lighter MSSM stau are the most important ones. The growth of the decay branching ratios into these states with $\tan\beta$ is a feature independent of $\cos\theta_{\tilde{t}}$.

The assumption that no tree-level-induced two-body decays are kinematically allowed implies that $m_{\tilde{\chi}^+_1} > m_{\tilde{t}_1} - m_b$. Therefore, one expects an increase of BR($\tilde{t}_1 \to b\,W^+\,\tilde{\chi}^0_1$) if $m_{\tilde{t}_1}$ increases, because the decay into $b\,W^+\,\tilde{\chi}^0_1$ is dominated by the t exchange, whereas for the decays into scalars $\tilde{\chi}^+_j$ exchange dominates. This trend is indeed observed in Fig. 2, where we show the branching ratios for $m_{\tilde{t}_1} = 350$ GeV. Here we concentrate on the range of $\cos\theta_{\tilde{t}}$ where $A_t \le 1$ TeV to avoid possible minima in the scalar potential which break either color or electric charge. Notice that for the heavy stop case the decay $\tilde{t}_1 \to b\,W^+\,\tilde{\chi}^0_1$ is the most important one, independently of $\cos\theta_{\tilde{t}}$ and $\tan\beta$.

Note however that also in this case R-parity violation implies a distinct signature compared to what is expected in the MSSM due to the decays of $\tilde{\chi}^0_1$. One gets the

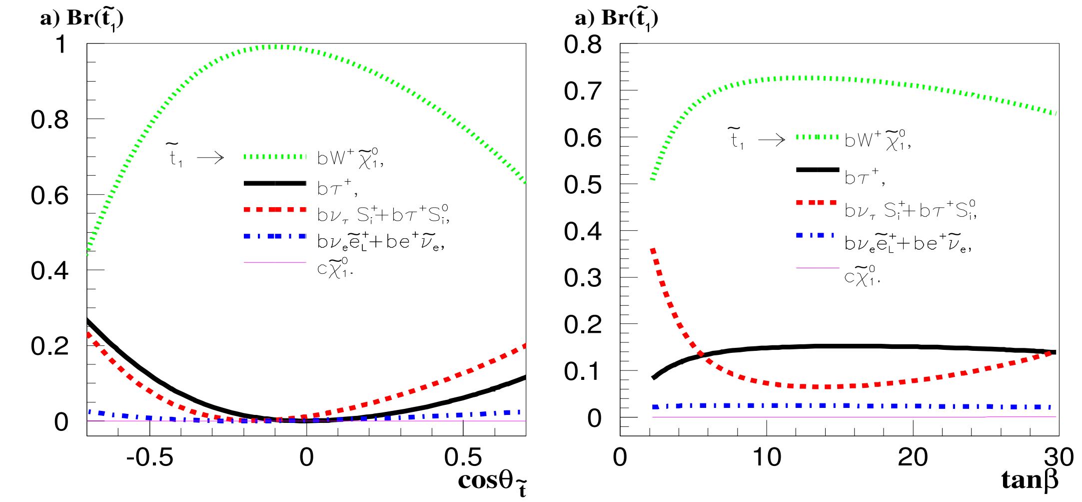

Figure 2: Branching ratios for $\tilde{t}_1$ decays for $m_{\tilde{t}_1} = 350$ GeV, $\mu = 750$ GeV, $M = 380$ GeV, and $m_\nu = 1$ keV. In a) the branching ratios are shown as a function of $\cos\theta_{\tilde{t}}$ for $\tan\beta = 4$, in b) as a function of $\tan\beta$ for $\cos\theta_{\tilde{t}} = 0.7$. Note, that the graph - - displays the sum $b\,\nu_\tau\, S_i^+ + b\,\tau^+\, S_i^0 + b\,\tau^+\, P_i^0$.

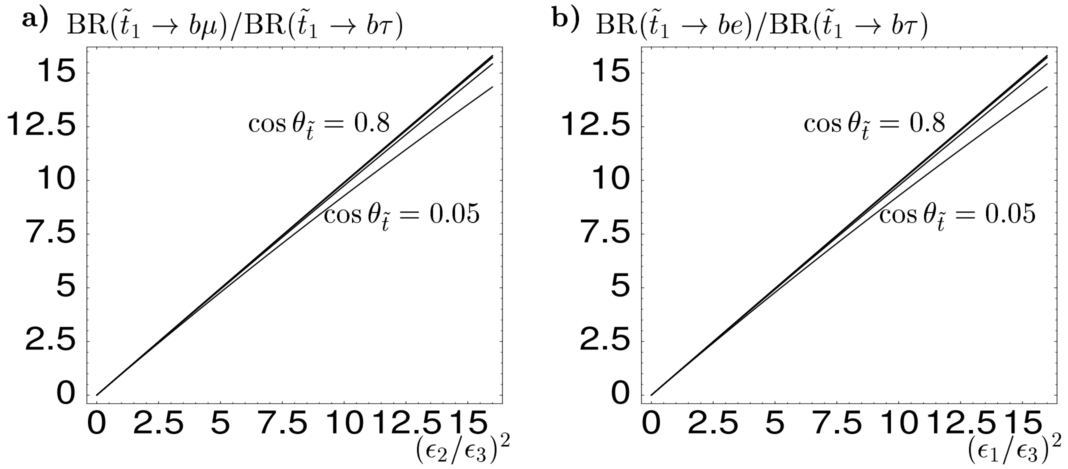

Figure 3: Ratio of branching ratios: a)BR($\tilde{t}_1 \to b\mu$)/BR($\tilde{t}_1 \to b\tau$) as a function of $(\epsilon_2/\epsilon_3)^2$ and b) BR($\tilde{t}_1 \to be$)/BR($\tilde{t}_1 \to b\tau$) as a function of $(\epsilon_1/\epsilon_3)^2$ for for $m_{\tilde{t}_1} = 220$ GeV, $\mu = 500$ GeV, $M = 240$ GeV, and $\epsilon_1^2 + \epsilon_2^2 + \epsilon_3^2 = 1$ GeV^2; $\cos\theta_{\tilde{t}} = 0.05, 0.1, 0.2, 0.40.8$, $m_{\nu_\tau} = 100$ eV.

following high-multiplicity final states:

$$\begin{aligned} \tilde{t}_1 &\to b\, W^+\, f\, \bar{f}\, \nu & (27)\\ &\to b\, W^+\, f\, \bar{f}'\, l^\pm & (28)\\ &\to b\, W^+\, \nu\, J & (29) \end{aligned}$$

where f denotes a Standard Model fermion. Here the decays of the W–boson will give additional leptons and jets. Therefore, one has in general additional jets and leptons compared to the MSSM case.

In the event that $\epsilon_{1,2}$ are of the same order of magnitude, as suggested by a solution to the present neutrino anomalies [31] one has in addition the decays into $b\,e^+$ and $b\,\mu^+$. If one passes from the 1-generation model to the 3-generation model the situation changes as follows. From Eq. (21) it follows that the sum of the modes $\Gamma(\tilde{t}_1 \to b\,l^+)$ in the 3–generation model is nearly equal to $\Gamma(\tilde{t}_1 \to b\,\tau^+)$ in the 1–generation model, if $(\epsilon_1')^2 +$

$(\epsilon_2')^2 + (\epsilon_2')^2 = \epsilon_3^2$ where the ϵ_i' (ϵ_3) are the parameters of the 3–generation (1–generation) model. In Fig. 3 we show the ratios of branching ratios for different $\tilde{t}_1 \to b\, l^+$ modes versus the ratios of different ϵ's squared and for different values of $\cos\theta_{\tilde{t}}$. In both cases we have fixed $\epsilon_1^2 + \epsilon_2^2 + \epsilon_3^2 = 1$ GeV^2 and in Fig. 3a $\epsilon_1 = \epsilon_3$ whereas in Fig. 3b $\epsilon_2 = \epsilon_3$. One can see that the dependence is nearly linear even for rather small $\cos\theta_{\tilde{t}}$. This result depends on $(\epsilon_1^2 + \epsilon_2^2 + \epsilon_3^2)/\mu^2$ and on the neutrino mass m_{ν_τ}, since both determine the mixings of the leptons with the charginos. The lines indicated in the figure come closer to the diagonal if $(\epsilon_1^2 + \epsilon_2^2 + \epsilon_3^2)/\mu^2$ increases and m_{ν_τ} decreases.

4 Conclusions

We have studied the phenomenology of the lightest stop in scenarios where the R-parity violating decay $\tilde{t}_1 \to b\,\tau^+$ competes with three–body decays. We have found that for $m_{\tilde{t}_1} \lesssim 250$ GeV there are regions of parameter where $\tilde{t}_1 \to b\,\tau^+$ is an important decay mode if not the most important one. This implies that there exists the possibility for full stop mass reconstruction from $\tau^+\tau^-\, b\,\bar{b}$ final states. Moreover, in this mass range the discovery of the lightest stop might not be possible at the LHC (certainly this is the case in the MSSM). This implies that one has to take into account the importance of this new decay mode when designing the stop search strategies at a future e^+e^- Linear Collider. Spontaneously and bilinearly broken R-parity violation imply additional leptons and/or jets in stop cascade decays. Looking at the three generation model the decays into $\tilde{t}_1 \to b\, l^+$ imply the possibility of measuring $\epsilon_e^2/\epsilon_\tau^2$ and $\epsilon_\mu^2/\epsilon_\tau^2$ and thereby probing the parameters associated with the present solar and atmospheric neutrino anomalies.

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