hep-ph/yymmnnn
FTUV/99-25
IFIC/99-27
FSU-HEP-990712
Version of August 3, 1999

Two-Body Decays of the Lightest Stop in Supergravity with and without R-Parity

Marco A. Díaz[1], D. A. Restrepo[2], and J. W. F. Valle[2]

[1]Department of Physics, Florida State University,
Tallahassee, Florida 32306, USA

[2]Departamento de Física Teórica, IFIC-CSIC, Universidad de Valencia
Burjassot, Valencia 46100, Spain

Abstract

We study the decays of the lightest stop quark in supergravity models with and without R-parity. Using the simplest model with an effective explicit bilinear breaking of R-parity and radiative electroweak symmetry breaking we show that, below the threshold for decays into charginos, the lightest stop decays mainly into third generation fermions, $\tilde{t}_1 \to b\,\tau$ instead of the R-parity conserving mode $\tilde{t}_1 \to c\,\tilde{\chi}^0_1$, even for tiny tau neutrino mass values. Moreover we show that, even above the threshold for decays into charginos, the decay $\tilde{t}_1 \to b\,\tau$ may be dominant. We study the rôle played by the universality of the boundary conditions on the soft supersymmetry breaking terms. This new decay mode $\tilde{t}_1 \to b\,\tau$ as well as the cascades originated by the conventional $\tilde{t}_1 \to c\,\tilde{\chi}^0_1$ decay followed by the R-parity violating neutralino decays can provide new signatures for stop production at LEP and the Tevatron.

1 Introduction

The *Minimal Supersymmetric Standard Model (MSSM)* [1] or its minimal supergravity (SUGRA) version [2] are by far the most well studied realizations of supersymmetry. 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 [3] in which R-parity is broken [4, 5, 6, 7]. The violation of R-parity could arise explicitly [8] as a residual effect of some larger unified theory [5], or spontaneously, through nonzero vacuum expectation values (vev's) for scalar neutrinos [4, 6, 7]. 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 [9, 10, 11, 12] 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. It plays an important rôle in making these models fully consistent with astrophysics and cosmology [3]. If the majoron acquires a small mass due to explicit breaking effects at the Planck scale then it may decay into neutrinos and photons, on very large time scales of cosmological interest, playing a possible role as unstable dark matter [13]. In this case LSP decays would not be seen in laboratory experiments [13] and the conventional missing momentum signal would be kept. 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 [12]. 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

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

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

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

as well as

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

in the case the breaking of R-parity is spontaneous [9, 10]. This last decay conserves R-parity since the majoron has a large R-odd singlet sneutrino component.

Although most searches for supersymmetric particles have so far been performed within the framework of the MSSM it has become clear that one needs alternative studies. In either of the above cases the supersymmetric (SUSY) particles need not be produced only in pairs, and the lightest of them could decay. The effects of R-parity violation can be large enough to be experimentally observable.

In this paper we focus on the decay modes of the lightest top squark in supergravity models where supersymmetry is realized with R-parity violation. In such models the lightest stop could even be the lightest supersymmetric particle and be produced at LEP. Neither e^+e^- collider data [14] nor $p\bar{p}$ data from the Tevatron [15] preclude this possibility. This way we refine the work presented in ref. [16] first by improving the treatment of the R-parity conserving decays and also by discussing the R-parity violating ones both in the context of supergravity-type models. In contrast with ref. [16] we explicitly consider the case of light tau neutrino masses, motivated by the simplest oscillation interpretation of the Super-Kamiokande atmospheric neutrino data.

For definiteness and simplicity we focus on supersymmetric models where the breaking of R-parity is parametrized explicitly through a bilinear superpotential term of the type ℓH_u [17]. The stop can have new decay modes such as

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

due to mixing between charged leptons and charginos. We show that this decay may be dominant or at least comparable to the ordinary R-parity conserving mode

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

where $\tilde{\chi}_1^0$ denotes the lightest neutralino.

The paper is organized as follows. The model and an analytical analysis of the tree–level tau neutrino in terms of SUGRA parameters is described in section 2. The mass matrices are given in section 3 while in section 4 we present the top squark decay widths in the minimal supergravity model with universal soft SUSY breaking terms [18], MSUGRA, for short. The relevant Feynman rules and the squark decay widths and branching ratios are calculated in appendix A. They are studied numerically in section 5 and we present our conclusions in section 6.

2 The Model

The supersymmetric Lagrangian is specified by the superpotential W given by

$$W = \varepsilon_{ab}\left[h_U^{ij}\widehat{Q}_i^a\widehat{U}_j\widehat{H}_2^b + h_D^{ij}\widehat{Q}_i^b\widehat{D}_j\widehat{H}_1^a + h_E^{ij}\widehat{L}_i^b\widehat{R}_j\widehat{H}_1^a - \mu\widehat{H}_1^a\widehat{H}_2^b\right] + \varepsilon_{ab}\epsilon_i\widehat{L}_i^a\widehat{H}_2^b\,, \qquad (6)$$

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 $\frac{1}{3}$, -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

$$\begin{aligned}
V_{soft} \;=\;& M_Q^{ij2}\widetilde{Q}_i^{a*}\widetilde{Q}_j^a + M_U^{ij2}\widetilde{U}_i^*\widetilde{U}_j + M_D^{ij2}\widetilde{D}_i^*\widetilde{D}_j + M_L^{ij2}\widetilde{L}_i^{a*}\widetilde{L}_j^a + M_R^{ij2}\widetilde{R}_i^*\widetilde{R}_j \\
&+m_{H_1}^2 H_1^{a*}H_1^a + m_{H_2}^2 H_2^{a*}H_2^a \\
&-\left[\tfrac{1}{2}M_3\lambda_3\lambda_3 + \tfrac{1}{2}M\lambda_2\lambda_2 + \tfrac{1}{2}M'\lambda_1\lambda_1 + h.c.\right] \\
&+\varepsilon_{ab}\left[A_U^{ij}h_U^{ij}\widetilde{Q}_i^a\widetilde{U}_jH_2^b + A_D^{ij}h_D^{ij}\widetilde{Q}_i^b\widetilde{D}_jH_1^a + A_E^{ij}h_E^{ij}\widetilde{L}_i^b\widetilde{R}_jH_1^a\right. \\
&\left.-B\mu H_1^aH_2^b + B_i\epsilon_i\widetilde{L}_i^aH_2^b\right]\,, \qquad (7)
\end{aligned}$$

For definiteness and simplicity we assume only R-parity Violation (RPV) in the third generation, but we do allow for 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.

Note that the bilinear term ϵ_3 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.

In order to study the R–Parity violating decay mode $\tilde{t}_1 \to b\,\tau$ it is sufficient to consider the superpotential [19, 20, 21]

$$W = h_t \widehat{Q}_3 \widehat{U}_3 \widehat{H}_2 + h_b \widehat{Q}_3 \widehat{D}_3 \widehat{H}_1 + h_\tau \widehat{L}_3 \widehat{R}_3 \widehat{H}_1 - \mu \widehat{H}_1 \widehat{H}_2 + \epsilon_3 \widehat{L}_3 \widehat{H}_2 \tag{8}$$

This amounts to neglecting the effects of RPV on the two first families.

Such model can be described in various equivalent bases, for example

1. one in which bilinear term and sneutrino vev are non-zero, $\epsilon_3^I \neq 0$ and $v_3^I \neq 0$ [3]

2. one in which trilinear and sneutrino vev are non-zero, $\lambda_3^{II} \neq 0$ and $v_3^{II} \neq 0$ [22]

3. the vev-less basis in which in which ϵ_3^{III} and λ_3^{III} are non-zero but $v_3^{III} = 0$ [22, 24]

Note that the R-parity violating parameters can be expressed in terms of dimensionless basis-independent alignment parameters $\sin\xi$, $\sin\xi'$ and $\sin\xi''$ [21, 25] defined by ($X = I, II$ or III)

$$\sin\xi = \frac{\epsilon_3^X v_d^X + \mu^X v_3^X}{\mu' v_d'} = \frac{v_3^{II}}{v_d'} = \frac{\epsilon_3^{III}}{\mu'} \tag{9}$$

$$\sin\xi' = \frac{\mu^X \lambda_3^X + \epsilon_3^X h_b^X}{\mu' h_b'} = \frac{\lambda_3^{II}}{h_b'} = \frac{\epsilon_3^I}{\mu'} \tag{10}$$

$$\sin\xi'' = \frac{-v_d^X \lambda_3^X + v_3^X h_b^X}{v_d' h_b'} = \frac{v_3^I}{v_d'} = -\frac{\lambda_3^{III}}{h_b'} \tag{11}$$

of which only two are independent because they satisfy

$$\sin\xi'' = \cos\xi' \sin\xi - \sin\xi' \cos\xi \tag{12}$$

and where

$$h_b' = \sqrt{h_b^{X\,2} + \lambda_3^{X\,2}} \qquad \mu' = \sqrt{\mu^{X\,2} + \epsilon_3^{X\,2}} \qquad v_d' = \sqrt{v_d^{X\,2} + v_3^{X\,2}}, \quad X = I, II, \text{ or } III \tag{13}$$

In the limit when the R-parity violating parameters vanish one recovers the MSSM. From now on we will work in the $\lambda_3^I = 0$–basis, unless otherwise stated. As a result we will omit the label I in all the parameters associated with this basis. We also will drop out the prime in h_b. One of the advantages in working in this basis is that the RGE's evolution does not induce the trilinear R-parity violating terms neither in the superpotential nor in the scalar potential if at the beginning we impose universality between the parameters B and B_3 [21].

It is convenient to introduce the following notation in spherical coordinates for the vacuum expectation values (vev):

$$\begin{aligned} v_d &= v \sin\theta \cos\beta \\ v_u &= v \sin\theta \sin\beta \\ v_3 &= v \cos\theta \end{aligned} \tag{14}$$

which preserves the MSSM definition $\tan\beta = v_u/v_d$. In the MSSM limit, where $\epsilon_3 = v_3 = 0$, the angle θ is equal to $\pi/2$. This make sense in the $\lambda_3^I = 0$–basis where the usual MSSM relation

$$m_b = \frac{1}{\sqrt{2}} h_b v_d \tag{15}$$

holds. In this model the presence of RPV induces a mass for the tau neutrino at the tree level [6, 7], as well as radiative masses to the the ν_e and ν_μ . As already mentioned it is sufficient for our present discussion of stop decays to keep only the tau neutrino. In order to study the ν_τ mass it is convenient to have an analytical expression for m_{ν_τ} in this limit. The tree level tau neutrino mass may be expressed as [8, 5, 21] – [30]

$$m_{\nu_\tau} \approx -\frac{(g^2M' + g'^2M)\mu'^2}{4MM'\mu'^2 - 2(g^2M' + g'^2M)\mu' v_u v_d' \cos\xi} {v_d'}^2 \sin^2\xi \tag{16}$$

in terms of basis-independent parameters μ', v_d' and $\sin\xi$ defined in Eqs (13) and (9). The second term in the denominator may be neglected if $M, \mu \gtrsim m_Z$, as often happens in minimal supergravity models with universal soft SUSY breaking terms [18]. Thus one may obtain an estimate of the neutrino mass by keeping only the first term in the denominator.

$$m_{\nu_\tau} \approx \frac{g^2}{2M} {v_d'}^2 \sin^2\xi \tag{17}$$

where we have used $M' = Mg'^2/g^2$. For $\sin\xi \approx 1$ one can easily check that m_{ν_τ} could be as large as the experimental upper bound of 18 Mev [31]. However in MSUGRA models one may obtain naturally small $\sin\xi$ values, calculable from the RGE evolution from the unification scale down to the weak scale, because from the minimization equations it can be written in terms $\Delta m^2 = m_{H_1}^2 - m_{L_3}^2$ and $\Delta B = B_3 - B$ [32] as

$$\sin\xi = -\cos\xi' \sin\xi' \left(\cos\xi \frac{\Delta m^2}{m'^2_{\tilde\nu^0_\tau}} + \frac{v_2}{v_d'} \frac{\mu' \Delta B}{m'^2_{\tilde\nu^0_\tau}} \right) \tag{18}$$

One may give a simplified approximate analytical expression for the tau neutrino mass in this model by solving the renormalization group equations for the soft mass parameters

basis–independent				Basis I: $\lambda_3^I=0$				Basis II: $\epsilon_3^{II}=0$		Basis III: $v_3^{III}=0$		
$\sin\xi$		$\sin\xi'$		ϵ_3(GeV)		v_3(GeV)		λ_3^{II}	v_3^{II}(GeV)	λ_3^{III}	ϵ_3^{III}(GeV)	
10^{-5}	10^{-4}	10^{-2}	10^{-4}	1	10^{-2}	1	10^{-3}	10^{-4}	10^{-3}	10^{-4}	10^{-3}	10^{-2}

Table 1: Estimated magnitude of R-parity violating parameters required for a tau neutrino mass in the eV range, without requiring cancellation in $\sin\xi$ in the three bases defined before.

$m^2_{H_1}$, $m^2_{L_3}$, B, and B_3 in one-step approximation. This gives the results [32]

$$\begin{aligned}\sin\xi\,|_{\Delta m^2} &\approx -\cos\xi'\sin\xi'\cos\xi h_b^2\left[\frac{m^2_{H_1}+M_Q^2+M_D^2+A_D^2}{m'^2_{\tilde\nu^0_\tau}}\right]\left(\frac{3}{8\pi^2}\ln\frac{M_U}{m_t}\right)\\ &\sim -\cos\xi'\sin\xi'\cos\xi h_b^2\left(\frac{3}{8\pi^2}\ln\frac{M_U}{m_t}\right)\end{aligned}\tag{19}$$

and

$$\sin\xi\,|_{\Delta B}\approx\cos\xi'\sin\xi'\tan\beta' h_b^2\left[\frac{\mu' A_D}{m'^2_{\tilde\nu^0_\tau}}\right]\left(\frac{3}{8\pi^2}\ln\frac{M_U}{m_t}\right)\tag{20}$$

where $\sin\xi|_{\Delta m^2}$ and $\sin\xi|_{\Delta B}$ are the two terms contributing to $\sin\xi$ in eq. (18). With them, we find

$$m_{\nu_\tau}\,|_{\min}\sim\frac{g^2m_b^2}{M}\left(\sin^2\xi' h_b^2\right)\left(\frac{3}{8\pi^2}\ln\frac{M_U}{m_t}\right)^2\tag{21}$$

The minimum value for $\sin\xi' h_b$ is determined by the value $\sin\xi'$ and that of $\tan\beta$. For $\sin\xi'\sim 1$ and relatively small $\tan\beta$ so that h_t is perturbative, one has

$$m_{\nu_\tau}\gtrsim 10\text{KeV}\tag{22}$$

for $M\sim 1$ TeV. In order to get smaller ν_τ masses one need to suppress $\sin^2\xi'$ additionally, for example to reach one electron-volt the required R-parity violating parameters are given in Table 1. These order of magnitude estimates are given in terms of the basis–independent angles ξ and ξ', and in the three basis defined before. Note that whenever the parameter has two values, the first correspond to $\tan\beta=2$ (the lower perturbativity limit) and the second to $\tan\beta=35$. In Table 1, $\sin\xi$ was estimated from Eq. (17) and $\sin\xi'$ from Eq. (21). Note that the induced RGE suppression depends basically in the h_b^2 factor in Eq. (19) that is 10^{-3} (1) for small (large) $\tan\beta$ and therefore the bigger the value of $\tan\beta$ the smaller $\sin\xi'$ will be. The RPV parameters in the several bases were estimated from Eqs. (9–11) and (12). As expected, the constraint on the trilinear RPV term is stronger than the obtained from the one-loop induced tau neutrino mass [33]

In eq. (21) we have neglect ΔB contribution with respect to the one coming from Δm^2. It is possible, however, that the ΔB term may be sizeable. In this case then it may cancel the Δm^2 contribution in $\sin\xi$, leading to an additionally suppressed neutrino mass. As we will see, however, in SUGRA models with universal soft terms at the unification scale (ϵSUGRA for short) we do not need any substantial cancellation in order to obtain calculable ν_τ masses below the electron-volt scale.

3 Squark Mass Matrices

The up and down-type squark mass matrices of our model have already been given previously in ref. [29]. Here we generalize those to the three-generation case. The mass matrix of the up squark sector follows from the quadratic terms in the scalar potential

$$V_{quadratic} = \left[\begin{array}{cc} \tilde{\mathbf{u}}_L^\dagger & \tilde{\mathbf{u}}_R^\dagger \end{array} \right] \tilde{M}_U^2 \left[\begin{array}{c} \tilde{\mathbf{u}}_L \\ \tilde{\mathbf{u}}_R \end{array} \right] + \cdots \tag{23}$$

given by

$$\tilde{M}_U^2 = \left[\begin{array}{cc} M_Q^2 + \frac{1}{2}v_u^2 h_U {h_U}^\dagger + \Delta_{UL} & \frac{1}{\sqrt{2}}v_u A_U^h - \frac{1}{\sqrt{2}}(\mu v_d - \epsilon_3 v_3) h_U \\ \frac{1}{\sqrt{2}}v_u {A_U^h}^\dagger - \frac{1}{\sqrt{2}}(\mu v_d - \epsilon_3 v_3) {h_U}^\dagger & M_U^2 + \frac{1}{2}v_u^2 {h_U}^\dagger h_U + \Delta_{UR} \end{array} \right] \tag{24}$$

where $\Delta_{UL} = \frac{1}{8}(g^2 - \frac{1}{3}{g'}^2)(v_d^2 - v_u^2 + v_3^2)\mathbf{1}$ and $\Delta_{UR} = \frac{1}{6}{g'}^2(v_d^2 - v_u^2 + v_3^2)\mathbf{1}$ are the splitting in the squark mass spectrum produced by electro-weak symmetry breaking, and $A^h_{Uij} \equiv A_U^{ij} h_U^{ij}$. The eigenvalues of $\tilde{M}_U^2$ are

$$\mathrm{diag}\left\{m_{\tilde{u}_1}, m_{\tilde{u}_2}, \ldots, m_{\tilde{u}_6}\right\} = \left[\begin{array}{cc} \Gamma_{UL} & \Gamma_{UR} \end{array} \right] \tilde{M}_U^2 \left[\begin{array}{c} \Gamma_{UL}^\dagger \\ \Gamma_{UR}^\dagger \end{array} \right] \tag{25}$$

This way the six weak-eigenstate fields $\tilde{u}_{iL}$ and $\tilde{u}_{iR}$ ($i = 1, 2, 3$) combine into six up-type mass eigenstate squarks $\tilde{u}_k$ as follows: $\tilde{u}_{iL} = \Gamma_{UL}^{\dagger ik}\tilde{u}_k = \Gamma_{UL}^{*ki}\tilde{u}_k$, $\tilde{u}_{iR} = \Gamma_{UR}^{\dagger ik}\tilde{u}_k = \Gamma_{UR}^{*ki}\tilde{u}_k$.

For completeness, we also give the mass matrix of the down squark sector. The quadratic scalar potential includes

$$V_{quadratic} = \left[\begin{array}{cc} \tilde{\mathbf{d}}_L^* & \tilde{\mathbf{d}}_R^* \end{array} \right] \tilde{M}_D^2 \left[\begin{array}{c} \tilde{\mathbf{d}}_L \\ \tilde{\mathbf{d}}_R \end{array} \right] + \cdots \tag{26}$$

given by

$$\tilde{M}_D^2 = \left[\begin{array}{cc} M_Q^2 + \frac{1}{2}v_d^2 h_D {h_D}^\dagger + \Delta_{DL} & \frac{1}{\sqrt{2}}v_d A_D^h - \frac{1}{\sqrt{2}}\mu v_u h_D \\ \frac{1}{\sqrt{2}}v_d {A_D^h}^\dagger - \frac{1}{\sqrt{2}}\mu v_u {h_D}^\dagger & M_D^2 + \frac{1}{2}v_d^2 {h_D}^\dagger h_D + \Delta_{DR} \end{array} \right] \tag{27}$$

where $\Delta_{DL} = -\frac{1}{8}(g^2+\frac{1}{3}g'^2)(v_d^2-v_u^2+v_3^2)\mathbf{1}$, $\Delta_{DR} = -\frac{1}{12}g'^2(v_d^2-v_u^2+v_3^2)\mathbf{1}$, and $A^h_{Dij} \equiv A^{ij}_D h^{ij}_D$. The eigenvalues of $\tilde{M}_D^2$ are

$$\mathrm{diag}\,\{m_{\tilde{d}_1}, m_{\tilde{d}_2}, \ldots, m_{d_6}\} = \left[\begin{array}{cc} \Gamma_{DL} & \Gamma_{DR} \end{array}\right] \tilde{M}_D^2 \left[\begin{array}{c} \Gamma^\dagger_{DL} \\ \Gamma^\dagger_{DR} \end{array}\right] \tag{28}$$

One is left with six mass-eigenstate down squarks fields $\tilde{d}_k$ related to $\tilde{d}_{iL}$ and $\tilde{d}_{iR}$ fields as follows: $\tilde{d}_{iL} = \Gamma^{\dagger ik}_{DL}\tilde{d}_k = \Gamma^{*ki}_{DL}\tilde{d}_k$, $\tilde{d}_{iR} = \Gamma^{\dagger ik}_{DR}\tilde{d}_k = \Gamma^{*ki}_{UR}\tilde{d}_k$.

For the Higgs-slepton part of the quadratic scalar potential in the one generation case of the Bilinear R-parity Violating (BRpV) model, see refs. [34] and [17].

Of particular interest to us is the chargino/tau mass matrix. For our present purposes it is sufficient to have the form of this matrix for one generation, which is given by

$$M_C = \left[\begin{array}{ccc} M & \frac{1}{\sqrt{2}}gv_2 & 0 \\ \frac{1}{\sqrt{2}}gv_d & \mu & -\frac{1}{\sqrt{2}}h_\tau v_3 \\ \frac{1}{\sqrt{2}}gv_3 & -\epsilon_3 & \frac{1}{\sqrt{2}}h_\tau v_d \end{array}\right] \tag{29}$$

This form is common to all models with spontaneous breaking of R-parity, as well as in the simplest truncation of these models provided by the BRpV model. We note that the chargino sector decouples from the tau sector in the limit $\epsilon_3 = v_3 = 0$. As in the MSSM, the chargino mass matrix is diagonalized by two rotation matrices U and V

$$U^* M_C V^{-1} = \left[\begin{array}{ccc} m_{\tilde{\chi}_1^\pm} & 0 & 0 \\ 0 & m_{\tilde{\chi}_2^\pm} & 0 \\ 0 & 0 & m_\tau \end{array}\right]. \tag{30}$$

The lightest eigenstate of this mass matrix must be the tau lepton ($\tau^\pm$) and so the mass is constrained to be $1.77705^{+0.29}_{-0.26}$ 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 [34]. The composition of the tau is given by

$$\tau_R^+ = V_{3j}\psi_j^+, \qquad \tau_L^- = U_{3j}\psi_j^- \tag{31}$$

where $\psi^{+T} = (-i\lambda^+, \widetilde{H}_2^1, \tau_R^{0+})$ and $\psi^{-T} = (-i\lambda^-, \widetilde{H}_1^2, \tau_L^{0-})$. The two component Weyl spinors τ_R^{0-} and τ_L^{0+} are weak eigenstates and, similarly, the two component Weyl spinors τ_R^+ and τ_L^- are mass eigenstates. It follows easily from eq. (30) that the matrix $M_C M_C^T$ is diagonalized by U and the matrix $M_C^T M_C$ is diagonalized by V.

The soft SUSY breaking parameters at the electroweak scale needed for the evaluation of the mass matrices and couplings are calculated by solving the renormalization group equations (RGE's) of the MSSM and imposing the radiative electroweak symmetry breaking condition. Taking the quark masses, the CKM matrix and $\tan\beta$ as inputs, we first solve one-loop RGE's for the gauge and Yukawa coupling constants to calculate the values at the unification scale, where we assume universal soft supersymmetry breaking boundary conditions. Then we solve the RGE's for all MSSM parameters downward with universal initial conditions for each set of soft SUSY breaking parameters. We include all generation mixing in the RGE's for Yukawa coupling constants as well as soft SUSY breaking parameters. Next, we evaluate the Higgs potential at the m_t scale including the one-loop corrections induced by the Yukawa coupling constants of the third generation. The requirement of the radiative electroweak symmetry breaking fixes the magnitude of the SUSY Higgs mass parameter μ and the soft SUSY breaking parameters B and B_3. At this stage we also choose ϵ_3 and v_3 in order to get a sufficiently light tau neutrino. At this point, all RPV parameters at the electroweak scale are determined as functions of the input parameters ($\tan\beta$, m_0, $A_0, m_{1/2}, \mathrm{sign}(\mu), \epsilon_3$). Notice that due to the third minimization condition one can solve for v_3 as a function of ϵ_3. Iteration is required because μ and ϵ_3 are inputs to evaluate the loop corrected minimum. Having determined all parameters at the electroweak scale, we obtain the masses and the mixings of all the SUSY particles by diagonalizing the corresponding mass matrices.

We scan the soft SUSY breaking parameter space in the range

$$\begin{array}{rcl} & m_0 & \leq 700\,\mathrm{GeV} \\ 100 < & m_{1/2} & \leq 400\,\mathrm{GeV} \\ & |A_0| & \leq 1000\,\mathrm{GeV}, \\ & |\epsilon_3| & < 200\,\mathrm{GeV} \\ 1.8 < & \tan\beta & < 60 \end{array}$$

the previous range on $\tan\beta$ guarantee that both h_t and h_b will be perturbative. For the CKM matrix, we use the convention of the Particle Data Group, taking $V_{us} = 0.2205$, $V_{cb} = 0.041$, $|V_{ub}/V_{cb}| = 0.08$ and $\delta = 0$.

The resulting region of lightest stop and chargino masses is displayed in Fig. 1. Neglecting the three body decays, we find that in Region I of the $m_{\tilde{t}_1}$–$m_{\tilde{\chi}_1^+}$ plane, $BR(\tilde{t}_1 \to c\tilde{\chi}_1^0) + BR(\tilde{t}_1 \to b\tau) \approx 1$. In Region II $BR(\tilde{t}_1 \to b\tau) + BR(\tilde{t}_1 \to b\tilde{\chi}_i^+) \approx 1$ (i=1,2). In

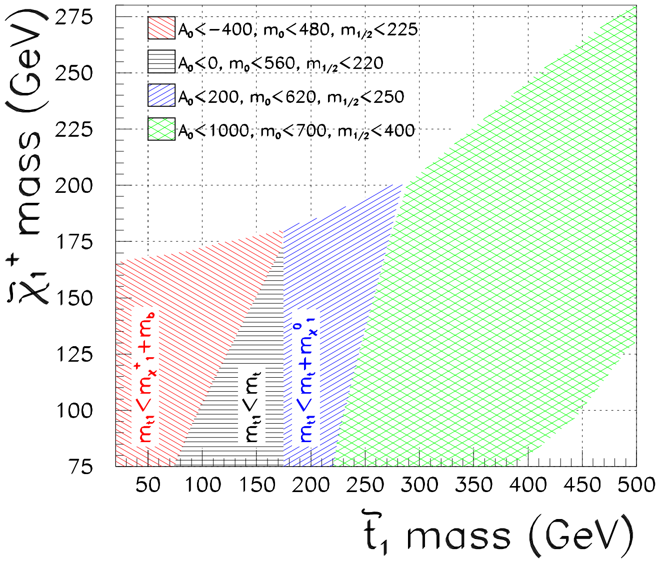


Figure 1: Kinematical regions in the $m_{\tilde{t}_1}$–$m_{\tilde{\chi}_1^+}$ plane. From left to right: Region I $m_{\tilde{t}_1} < m_{\tilde{\chi}_1^+} + m_b$; Region II $m_{\tilde{\chi}_1^+} + m_b < m_{\tilde{t}_1} < m_t$; Region III $m_t < m_{\tilde{t}_1} < m_{\tilde{\chi}_1^0} + m_t$; and region IV $m_{\tilde{t}_1} > m_{\tilde{\chi}_1^0} + m_t$

Region III $BR(\tilde{t}_1 \to b\,\tau) + BR(\tilde{t}_1 \to b + \tilde{\chi}_i^+) + BR(\tilde{t}_1 \to t\,\nu_\tau) \approx 1$ (i=1,2), while in region IV $BR(\tilde{t}_1 \to b\,\tau) + BR(\tilde{t}_1 \to b\,\tilde{\chi}_i^+) + BR(\tilde{t}_1 \to t\,\nu_\tau) + BR(\tilde{t}_1 \to t\,\tilde{\chi}_j^0) \approx 1$ ($j = 1, \ldots, 4$). Note that in each region the exact equality to 1 is reached when the FCNC processes are fully included.

In Appendix A we give the Feynman rules for all vertices involving squarks, quarks and charginos or neutralinos, as well as the two–body squark decay widths, for squarks of all three generations. These equations reduce to the expressions found in ref. [35] provided one identifies $\Gamma_{UL}^{33} = \cos\theta_{\tilde{t}}$ and $\Gamma_{UR}^{33} = \sin\theta_{\tilde{t}}$. They also generalize the results for the BRpV model to the three generation case.

4 Two-Body Decays of the Lightest Stop in MSUGRA

In the region I of Fig. 1 the main $\tilde{t}_1$ decay channel is expected to be the loop–induced and flavour–changing $\tilde{t} \to c\,\tilde{\chi}_1^0$ [36, 37, 38]. The FCNC processes in the MSSM in general involve a very large number of input parameters. Following common practice, we prefer to perform the phenomenological study of flavour changing processes in the framework of a supergravity theory with universal supersymmetry breaking. The simplest description of FCNC processes in R-parity conserving minimal SUGRA models uses the so-called one-step approximation. Here we start by reproducing the standard calculation for $\tilde{t} \to c\,\tilde{\chi}^0_1$ as in [36]. To do this consider only the effect of the third generation Yukawa coupling. From our general Eq. (A.9) we have for $\tilde{t}_1 = \tilde{u}_l$

$$\Gamma(\tilde{t}_1 \to c\,\tilde{\chi}^0_1) \approx \frac{g^2}{8\pi}\left(\Gamma_{UL13}\right)^2 \left[\tfrac{2}{3}\sin\theta_W N'_{11} + \left(\tfrac{1}{2} - \tfrac{2}{3}\sin^2\theta_W\right)\frac{N'_{12}}{\cos\theta_W}\right]^2 m_{\tilde{t}_1}\left(1 - \frac{m^2_{\tilde{\chi}^0_1}}{m^2_{\tilde{t}_1}}\right)^2 \tag{32}$$

with

$$\Gamma_{UL13} = \frac{\Delta_L \cos\theta_{\tilde{t}} - \Delta_R \sin\theta_{\tilde{t}}}{m^2_{\tilde{c}_L} - m^2_{\tilde{t}_L}} \tag{33}$$

and by using the one–step approximation[1]

$$\Delta_L = (\tilde{M}^2_U)_{23} \approx (M^2_Q)_{23} \approx -\frac{t_U}{16\pi^2}K_{cb}K_{tb}h_b^2(M^2_Q + M^2_D + m^2_{H_1} + A^2_b) \tag{34}$$

$$\Delta_R = (\tilde{M}^2_U)_{26} \approx (A^h_U)_{23} \approx -\frac{t_U}{16\pi^2}K_{cb}K_{tb}h_b^2 m_t(A_b + \frac{1}{2}A_t) \tag{35}$$

with $t_U = \ln(M_G/m_t)$. So, in the one–step approximation we have

$$\begin{aligned}\Gamma(\tilde{t}_1 \to c\,\tilde{\chi}^0_1) \approx\ & \frac{g^2}{16\pi}\left(\frac{t_U}{16\pi^2}K_{cb}K_{tb}\right)^2 h_b^4(\delta_{m_0^2}\cos\theta_{\tilde{t}} - \delta_A \sin\theta_{\tilde{t}})^2 \\ & \times\left[\frac{\sqrt{2}}{6}(\tan\theta_W\, N_{11} + 3N_{12})\right]^2 m_{\tilde{t}_1}\left(1 - \frac{m^2_{\tilde{\chi}^0_1}}{m^2_{\tilde{t}_1}}\right)^2\end{aligned} \tag{36}$$

where the prefactor is $\sim 6 \times 10^{-7}$ and the parameter $\delta_{m_0^2}$ is given by

$$\delta_{m_0^2} = \frac{M^2_Q + M^2_D + m^2_{H_1} + A^2_b}{m^2_{\tilde{c}_L} - m^2_{\tilde{t}_L}} \sim 1 \tag{37}$$

is basically independent of the initial conditions due to the m_0 dependence both in the numerator as in the denominator and

$$\delta_A = \frac{m_t(A_b + \frac{1}{2}A_t)}{m^2_{\tilde{c}_L} - m^2_{\tilde{t}_L}} \tag{38}$$

Note however, that the one-step approximation includes only the third generation Yukawa couplings and neglects the running of the soft breaking terms [36, 37, 38, 46]. Such an approximation is rather poor for our purposes, since we will be interested in comparing with other R-parity violating decay modes (see the next section). In order to have an accurate calculation of the respective branching ratios we need to go beyond the one-step approximation. We therefore use an improved calculation for the FCNC process

[1]In [36] the factor $\frac{1}{2}A_t$ is missing.

$\tilde{t} \to c\,\tilde{\chi}_1^0$ in which the running of the Yukawa couplings and soft breaking terms is taken into account. First we note that we have checked that indeed the effect of the Yukawas from the two first generations is negligible. However the same is not true for the running of the soft breaking terms. As can be seen from Figure 2 the range of variation that we obtain from the numerical solution is

$$\Gamma(\tilde{t}_1 \to c\,\tilde{\chi}_1^0) \sim (10^{-16} - 10^{-6})\mathrm{GeV} \tag{39}$$

depending on the assumed value of $m_{1/2}$ and $\tan\beta$. In this figure we have compared the decay width obtained from Eq. (A.9) with the approximate formulae in Eq. (32) for two fixed values of $m_{1/2}$, $\tan\beta$ and taking $A_0 = 0$. The approximate formulae only reproduce well the numerical result for the academic case of no SUSY breaking gaugino mass $m_{1/2} = 0$. For the more realistic case $m_{1/2} > 100\,\mathrm{GeV}$, the exact solution is usually one decade smaller than the approximate one. In the one-step approximation $\Gamma(\tilde{t}_1 \to c\,\tilde{\chi}_1^0)$ can be arbitrarily small if the two terms $\delta_{m_0^2}\cos\theta_{\tilde{t}}$ and $\delta_A \sin\theta_{\tilde{t}}$ in Eq. (36) cancel. This behaviour can be illustrated in Figure 2 by the dashed line labelled 358, which corresponds to $m_0 = 358$ GeV. One sees clearly that while the approximate solution goes to zero, while the numerical one reaches a minimum value around 10^{-11} GeV. The wrong behaviour of the approximate solution indicates that the δ_A depends strongly on the scale. For example, the RGE for A_b is very sensitive on $m_{1/2}$ and $\tan\beta$ and in the one-step approximation there is no explicit dependence on $m_{1/2}$, which is crucial. Both solutions increase with $\tan\beta$, as expected by the bottom Yukawa dependence explicit in Eq. (36) and remain practically constant for large $m_{1/2}$ values.

5 Two-Body Decays of the Lightest Stop: the R-parity violating case

With this at hand we can easily compute the R-parity violating stop decay width $\tilde{t}_1 \to b\,\tau$, where $\tau = F_3^+$ from eq. (A.10):

$$\begin{aligned}\Gamma(\tilde{t}_1 \to b\,\tau) &= \frac{g^2\lambda^{1/2}(m_{\tilde{t}_1}^2, m_b^2, m_\tau^2)}{16\pi m_{\tilde{t}_1}^3}\{-4U_{32}^*\hat{h}_b c_{\theta_{\tilde{t}}}(V_{32}^*\hat{h}_t s_{\theta_{\tilde{t}}} - V_{31}^* c_{\theta_{\tilde{t}}})m_b m_\tau \\ &\quad +[(V_{32}^*\hat{h}_t s_{\theta_{\tilde{t}}} - V_{31}^* c_{\theta_{\tilde{t}}})^2 + U_{32}^{*2}\hat{h}_b^2 c_{\theta_{\tilde{t}}}^2](m_{\tilde{t}_1}^2 - m_b^2 - m_\tau^2)\}\end{aligned} \tag{40}$$

which coincides with the result in ref. [38]. In [39] it was shown that, except for U_{32} which determines the SU(2)-conserving mixing of the Higgsino with the left-handed τ, all V_{3i}

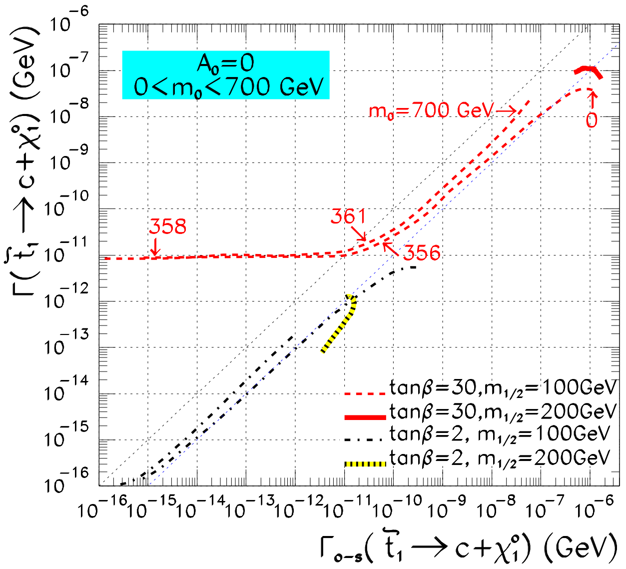


Figure 2: Comparison between the exact numerical calculation (ordinate) and the one–step approximation (abscissa) for the $\tilde{t}_1 \to c\,\tilde{\chi}_1^0$ decay width for various values of $\tan\beta$ and $m_{1/2}$ with $A_0 = 0$ and m_0 varying in the indicated range. The solid left diagonal line would signify the equality between the estimates, while the right diagonal line would indicate one order of magnitude difference. Results of both estimates indicated in lower right legend. More details in text.

and U_{3i} are proportional to v_3^{II} and therefore to the tau neutrino mass. Neglecting these terms we have from Eq. (40)

$$\Gamma(\tilde{t}_1 \to b\,\tau) \approx \frac{g^2 \lambda^{1/2}(m_{\tilde{t}_1}^2, m_b^2, m_\tau^2)}{16\pi m_{\tilde{t}_1}^3} \sin^2\xi' \hat{h}_b^2 c_{\theta_{\tilde{t}}}^2 (m_{\tilde{t}_1}^2 - m_b^2 - m_\tau^2) \tag{41}$$

noting that, to a good approximation,

$$|U_{32}| \approx \left|\frac{\epsilon_3}{\mu'}\right| = |\sin\xi'| \tag{42}$$

The dependence on the factor $\sin^2\xi' h_b^2$ may be more easily seen in basis II, where $\epsilon_3^{II} = 0$. In this case v_3^{II} is proportional to the tau neutrino mass and therefore in this basis all the elements U_{3i} and V_{3i} are small [21]. Neglecting these terms, $\Gamma(\tilde{t}_1 \to b\,\tau)$ may be written directly from the interaction term $\tilde{t}_L b_R \tau_L$. It is induced by the trilinear term in the $\epsilon_3^{II} = 0$–basis given in Eq. (10) as

$$\lambda_3^{II} = (\epsilon_3/\mu')\, h_b = h_b \sin\xi' \tag{43}$$

which reproduces Eq. (41). In our numerical calculation to be described in the next section we have used for $\Gamma(\tilde{t}_1 \to b\,\tau)$ the full expression given in Eq. (A.9) of the appendix.

In the next section we will determine the conditions under which the R-parity violating decay width $\Gamma(\tilde{t}_1 \to b\,\tau)$ can be dominant over the R-parity conserving ones, $\Gamma(\tilde{t}_1 \to c\,\tilde{\chi}^0_1)$ and $\Gamma(\tilde{t}_1 \to b\,\tilde{\chi}^+_1)$.

5.1 Region I

Using the one-step approximation for $\Gamma(\tilde{t}_1 \to c\,\tilde{\chi}^0_1)$ one finds

$$\Gamma(\tilde{t}_1 \to c\,\tilde{\chi}^0_1) \sim 10^{-6} h_b^4 m_{\tilde{t}_1} \tag{44}$$

Using the Eq. (41) and neglecting charm, tau and bottom masses we get

$$\frac{\Gamma(\tilde{t}_1 \to c\,\tilde{\chi}^0_1)}{\Gamma(\tilde{t}_1 \to b\,\tau)} \sim 10^{-5} \frac{h_b^2}{\sin^2 \xi'} \tag{45}$$

Therefore $\Gamma(\tilde{t}_1 \to c\,\tilde{\chi}^0_1)$ will start to compete with $\Gamma(\tilde{t}_1 \to b\,\tau)$ from $\sin\xi' \lesssim 5 \times 10^{-3}$ (10^{-4}) for $\tan\beta$ large (small). In fig 3 we compare $BR(\tilde{t}_1 \to c\,\tilde{\chi}^0_1)$ with $BR(\tilde{t}_1 \to b\,\tau)$ within the restricted region of the $m_{\tilde{t}_1}$–$m_{\tilde{\chi}^0_1}$ plane where only those two decay modes are open. We consider different m_{ν_τ} values (these correspond to relatively small values of the R-parity parameters $|\epsilon_3|, |v_3| \lesssim 1$ GeV. We vary the MSSM parameters randomly obeying the condition $m_{\tilde{t}_1} < m_{\tilde{\chi}^\pm_1} + m_b$ and depict the corresponding region in dark or light grey (yellow or green, respectively, in the coloured version). The upper–left triangular region corresponds to $m_{\tilde{t}_1} < m_{\tilde{\chi}^0_1} + m_c$ so that $BR(\tilde{t}_1 \to b\,\tau) = 100\,\%$. The lower–right triangular region corresponds to $m_{\tilde{t}_1} > m_{\tilde{\chi}^0_1} + m_b$. One notices from fig 3 that in the central region the dominant stop decay mode is $\tilde{t}_1 \to b\,\tau$ with branching ratio $BR(\tilde{t}_1 \to b\,\tau) > 0.9$. The dotted lines in the light grey region indicate maximum ν_τ mass values obtained in the scan. Therefore if the lightest stop only decays into the two modes considered here, the processes $\tilde{t}_1 \to b\,\tau$, will be important for the light tau neutrino mass range indicated by the simplest oscillation interpretation of the Super-Kamiokande atmospheric neutrino data. This may also be seen in Fig. 4 where the rôle played by $\tan\beta$ is manifest. In this figure we have shown $BR(\tilde{t}_1 \to b\,\tau)$ as function of the lighter stop mass for tau neutrino mass in the sub–eV range. We have obtained this kind of tau neutrino masses numerically allowing only one decade of cancellation between the two terms that contribute to $\sin\xi$ in eq. (18). The degree of suppression for ϵ_3/μ obtained numerically agrees very well with the expectations from the approximate formula for the minimal tau neutrino mass in Eq. (21).

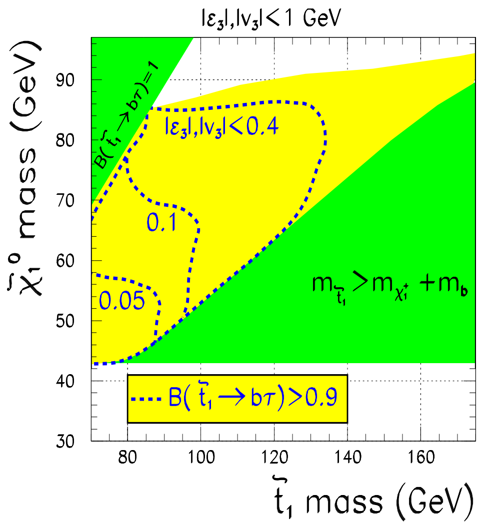


Figure 3: Regions where the $\tilde{t}_1 \to b\,\tau$ decay branching ratio exceeds 90% in the $m_{\tilde{t}_1}$–$m_{\tilde{\chi}_1^0}$ plane for different m_{ν_τ} values. The MSSM parameters are randomly varied as indicated in the text under the restriction $m_{\tilde{t}_1} < m_{\tilde{\chi}_1^\pm} + m_b$. The upper–left triangular region corresponds to $m_{\tilde{t}_1} < m_{\tilde{\chi}_1^0} + m_c$ so that only the $\tilde{t}_1 \to b\,\tau$ decay channel is open. The lower–right triangular region corresponds to $m_{\tilde{t}_1} > m_{\tilde{\chi}_1^+} + m_b$.

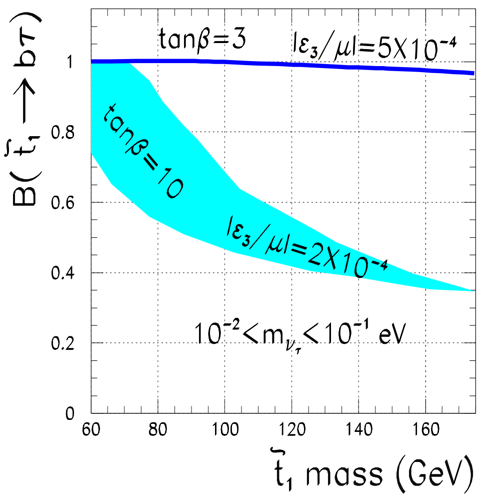


Figure 4: $BR(\tilde{t}_1 \to b\,\tau)$ as function of the lighter stop mass for tau neutrino mass in the sub–eV range and two different values of $\tan\beta$ and ϵ_3/μ. This prediction is natural in the sense that we have allowed only up to one order of magnitude of cancellation between the two terms that contribute to $\sin\xi$.

5.2 Region II

In region II the R–Parity conserving decay mode $\tilde{t}_1 \to b\tilde{\chi}_1^+$ is open and competes with the R–Parity violating mode $\tilde{t}_1 \to b\,\tau$. Replacing the subindex 3 by 1 on the diagonalization matrices U and V in eq. (40) we get the corresponding expression for $\Gamma(\tilde{t}_1 \to b\,\tilde{\chi}_1^+)$. In order to get an approximate expression for the ratio of the two main decay rates in this region, consider that in MSUGRA the lightest chargino is usually gaugino-like implying that $V_{11}^2 \sim 1$. In addition, the lightest stop is usually right, hence $\sin^2\theta_{\tilde{t}} \gtrsim \cos^2\theta_{\tilde{t}}$. This way we find

$$\frac{\Gamma(\tau)}{\Gamma(\tilde{\chi}_1^+)} \equiv \frac{\Gamma(\tilde{t}_1 \to b\,\tau)}{\Gamma(\tilde{t}_1 \to b\,\tilde{\chi}_1^+)} \approx \frac{\sin^2\xi' \hat{h}_b^2 \cos^2\theta_{\tilde{t}}}{\left[(V_{11}^* \cos\theta_{\tilde{t}} - V_{12}^* \hat{h}_t \sin\theta_{\tilde{t}})^2 + U_{12}^{*2} \hat{h}_b^2 \cos^2\theta_{\tilde{t}}\right]} K \tag{46}$$

where K is a kinematical factor that tends to increase the ratio, and $\hat{h}_{t,b} \equiv h_{t,b}/g$. The presence of the bottom quark Yukawa coupling indicates that large values of $\tan\beta$ are necessary to have large R–Parity violating branching ratios. In fact, we have numerically checked with the exact expressions that in Region II (RII) $\Gamma(\tau)/\Gamma(\tilde{\chi}_1^+) \gtrsim 1$ only for large $\tan\beta$ as we will see in the next figures. Note that if the stops have a small mixing ($\cos\theta_{\tilde{t}} \approx 0$), then $\Gamma(\tau)/\Gamma(\tilde{\chi}_1^+) \ll 1$ in RII.

In Fig. 5 we plot in the $m_{\chi_1^\pm} - m_{\tilde{t}_1}$ plane the regions where $BR(\tilde{t}_1 \to b\,\tau)$ dominates over $BR(\tilde{t}_1 \to b\tilde{\chi}_1^+)$. In the upper–left region the decay mode $\tilde{t}_1 \to b\tilde{\chi}_1^+$ is not allowed and corresponds to Region I. Below and to the right of this forbidden zone, and above and to the left of the inclined lines, lies the region where $\Gamma(\tau)/\Gamma(\tilde{\chi}_1^+) > 1$ of RII. Three lines have been drawn corresponding to $|\epsilon_3| < 80$ GeV (dashed), $|\epsilon_3| < 60$ GeV (dotted), and $|\epsilon_3| < 40$ GeV (dot–dashed). The proximity to the forbidden zone indicates that the the RPV decay dominates for large values of the RPV parameters only close to the threshold where there is a high kinematical suppression through the factor K.

An even more simpler expression for the ratio of decay rates in eq. (46) is obtained if we take $V_{11} \approx 1$ and assume no kinematical suppression in eq. (46) through the factor K:

$$\frac{\Gamma(\tau)}{\Gamma(\tilde{\chi}_1^+)} \sim \sin^2\xi' \hat{h}_b^2 . \tag{47}$$

Note that the presence of the parameter $\sin\xi' = \epsilon_3/\mu'$ indicates that the R–Parity violating decay mode is not proportional to the neutrino mass, and rather proportional to the BRpV parameter ϵ_3^2.

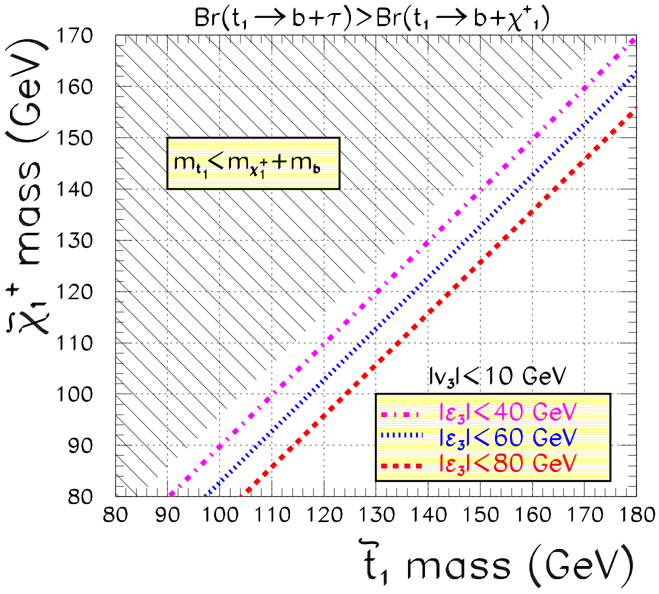


Figure 5: Contours of $BR(\tilde{t}_1 \to b\,\tau) > BR(\tilde{t}_1 \to b\tilde{\chi}_1^+)$ in the $m_{\tilde{t}_1}$–$m_{\tilde{\chi}_1^+}$ plane for $|v_3| < 10$ GeV. Three different maximum values for $|\epsilon_3|$ are considered: $|\epsilon_3| < 40$ GeV (dot-dash), $|\epsilon_3| < 60$ GeV (dots), and $|\epsilon_3| < 80$ GeV (dashes). The region where $m_{\tilde{t}_1} < m_b + m_{\tilde{\chi}_1^+}$ corresponds to the previously studied Region I.

In Fig. 6 we plot $\Gamma(\tau)/\Gamma(\tilde{\chi}_1^+)$ in RII as a function of the tau neutrino mass. Both decay rates have been calculated numerically from the exact formulas. In this figure we have imposed $m_{H_1}^2 = m_{L_3}^2$ at the GUT scale, but we have not imposed universality between B_3 and B. If we accept cancellation within one decade between the ΔB and Δm^2 terms in the neutrino mass formula of eq. (18), then the allowed region is at the right and down the corresponding dashed inclined line. If larger cancellation are accepted, the left boundary of the allowed region moves to the left as indicated in the figure. We have set the right boundary of the allowed region at the collider experimental limit of the tau neutrino mass, and have chosen fixed values of $\epsilon_3/\mu = 1$, 0.1, and 0.01. The allowed region for $\epsilon_3/\mu = 1$ is above the dashed line. In the case of $\epsilon_3/\mu = 0.1$ (0.01) the allowed region lies between the solid (dotted) lines. The effect of $\tan\beta$ is to increase the ratio $\Gamma(\tau)/\Gamma(\tilde{\chi}_1^+)$: the minimum value of the ratio is obtained for $\tan\beta \approx 2$ and the maximum for $\tan\beta \approx 60$. The extreme values of $\tan\beta$ are dictated by perturbativity.

A number of statistically less significant points appear outside the drawn regions in Fig. 6 and are not included. They correspond to points with $m_{\tilde{t}_1} - m_b - m_{\tilde{\chi}_1^\pm} < 10\,\text{GeV}$ which appear above the diagonal line, and points with $\cos\theta_{\tilde{t}} < 0.1$ which appear below the horizontal line corresponding to the lowest values of $\tan\beta$. In the last case, our

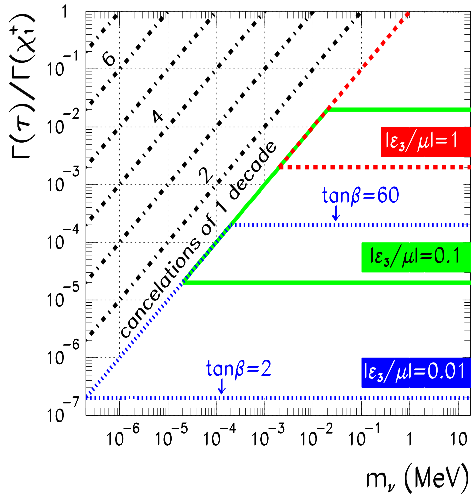


Figure 6: Zones for $\Gamma(\tilde{t}_1 \to b\,\tau)/\Gamma(\tilde{t}_1 \to b\,\tilde{\chi}_1^+)$ as a function of the tau neutrino mass for different amounts of cancellation between the two terms that contribute to the neutrino mass. We impose the universality condition $m_{H_1}^2 = m_{L_3}^2$ at the unification scale, but B and B_3 are not universal. We take $|\epsilon_3/\mu| = 1$ (inside the dashed lines), $|\epsilon_3/\mu| = 0.1$ (solid lines), and $|\epsilon_3/\mu| = 0.01$ (dotted lines).

approximation in eq. (47) does not work any more. On the other hand, If $\cos\theta_{\tilde{t}} > 0.1$, then eq. (47) predicts very well the behavior of $\Gamma(\tau)/\Gamma(\tilde{\chi}_1^+)$. For example considering $\epsilon_3/\mu = 1$, which is equivalent to $\sin\xi' = 1/\sqrt{2}$, we expect from eq. (47) a maximum value of order 1 for large $\tan\beta$ ($h_b \approx 1$) and a minimum value of order 10^{-3} for small $\tan\beta$ ($h_b^2 \approx 10^{-3}$), and this is confirmed by Fig. 6. In addition, if we accept only cancellation within a decade between the two terms that contribute to the tau neutrino mass, then our approximate formula in eq. (21) which predict the minimum tau neutrino mass, works very well.

In Fig. 7 we plot the ratio $\Gamma(\tau)/\Gamma(\tilde{\chi}_1^+)$ in RII as a function of the tau–neutrino mass, but this time imposing universality $B = B_3$ at the GUT scale within 0.1%. Cancellation between the Δm^2 and ΔB terms in the neutrino mass are accepted only within 1 decade. As a reference we have drawn the line corresponding to $\Delta B = 0$ and $\Delta m^2 = \Delta m^2_{min}$ (Δm^2 is negative and its magnitude is bounded from below by Δm^2_{min}) at the weak scale, which gives an idea of the value of the neutrino mass when there is no cancellation between the ΔB and Δm^2 terms. The region of high neutrino mass and small R–Parity violating branching ratio is disallowed, but this region is not theoretically interesting anyway. High values of the R–parity violating branching ratio for large values of ϵ_3 are highly restricted for large $\tan\beta$. This can be understood as follows. In the case of $\epsilon_3/\mu = 1$ and $\tan\beta = 60$

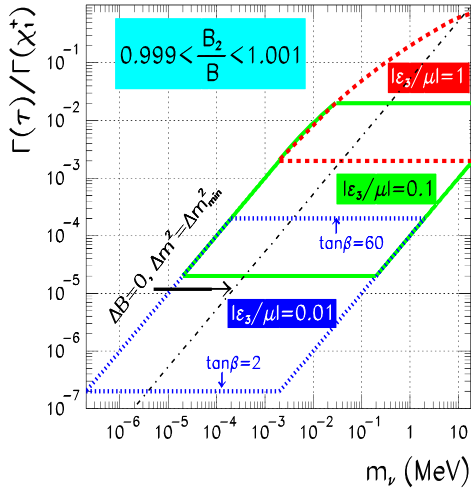


Figure 7: Zones for $\Gamma(\tilde{t}_1 \to b\,\tau)/\Gamma(\tilde{t}_1 \to b\,\tilde{\chi}_1^+)$ as a function of the tau neutrino mass, but now the universality condition $B = B_3$ at the unification scale condition is imposed at 0.1% as indicated. Its effect is to alter the maximum attainable tau neutrino mass. The dot-dash line corresponds to the case where $\Delta B = 0$ at the weak scale.

acceptable neutrino masses are obtained only if $\sin\xi \sim 1$. On the other hand, in this regime from eq. (18) we find that the ΔB term is large because of the high value of $\tan\beta$, and that the Δm^2 term is large because $m^2_{H_1}$ becomes negative and $\Delta m^2 = m^2_{H_1} - m^2_{L_3}$ grows in magnitude. In this way, acceptable neutrino masses are achieved only with cancellation within more than one decade. We think that Fig. 7 is very conservative considering that in MSSM–SUGRA with unification of top-bottom-tau Yukawa couplings, the large value of $\tan\beta$ implies that a cancellation of four decades among vev's is needed.

5.3 Effects of non–Universality

We now study the effect of possible non-universality of soft-breaking SUSY parameters on our previous results. SUGRA spectra are typically found for given values of $m_{1/2}$, m_0, A_0, $\tan\beta$ and Sgn(μ). In our case we have in addition $\sin\xi'$ (or equivalently, ϵ_3). The value of v_3 is determined by the previous parameters, and a relation between A_0 and the ratio B_3/B at the GUT scale emerges, which indicates the degree of universality. This relation can be seen in Fig. 8 for $\epsilon_3/\mu = 1$ and the values $\tan\beta = 3$, 40, and 60. The relation becomes more and more restrictive as $\tan\beta$ is increased, starting from $-1000 < A_0 < 1000$ GeV allowed for $\tan\beta = 3$, until practically only a single value of A_0 being compatible with unification if $\tan\beta = 60$.

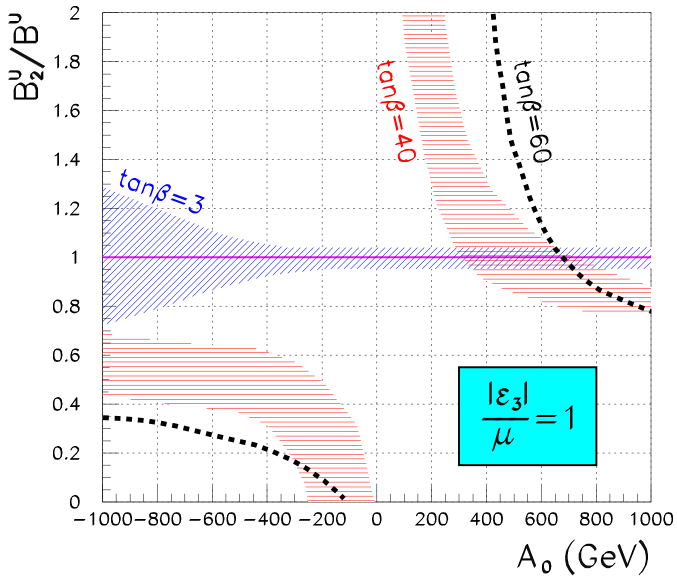


Figure 8: Universality condition $B = B_3$ at the unification scale as a function of A_0. As $\tan\beta$ increase, the allowed values of A_0 are more constrained in order to get universality at the GUT scale.

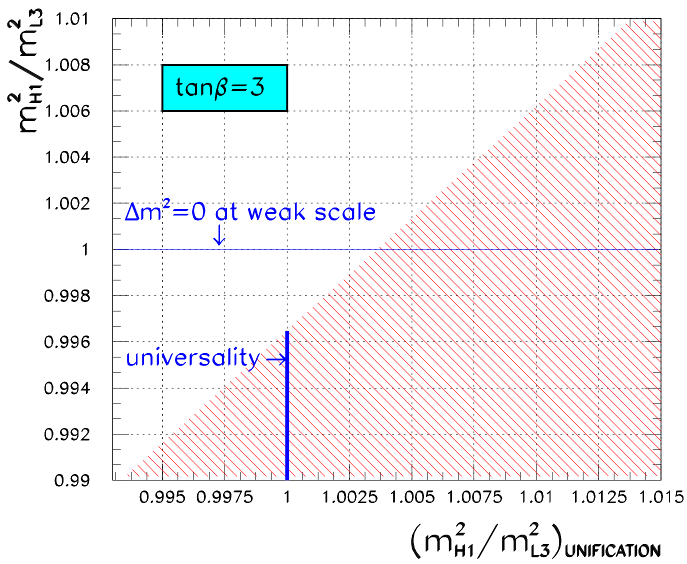


Figure 9: Comparison between the ratio $m^2_{H_1}/m^2_{L_3}$ at the weak and the unification scales for $\tan\beta = 3$. The universality at the unification scale, $m^2_{H_1}/m^2_{L_3} = 1$, implies a minimum values for this rate at the weak scale.

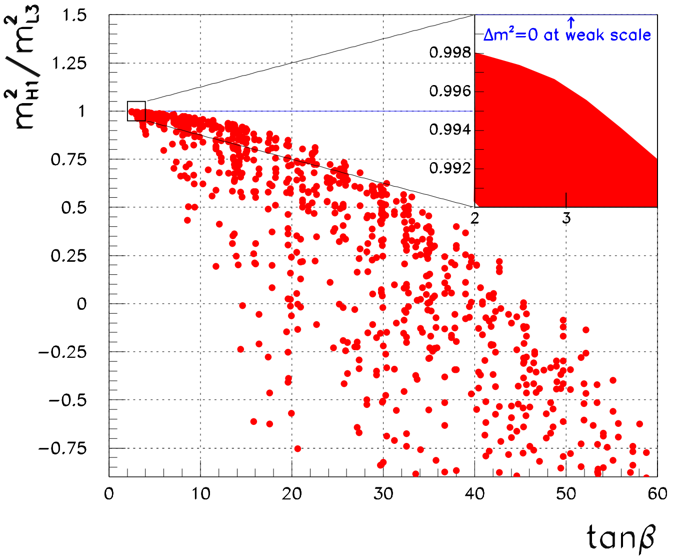


Figure 10: Misalignment in Δm^2 at the weak scale as a function of $\tan\beta$. It is always different from zero, negative and increase with $\tan\beta$.

In Fig. 9 we illustrate the effect of relaxing the universality condition $m^2_{H_1} = m^2_{L_3}$ at the GUT scale. For $\tan\beta = 3$ we plot the ratio $m^2_{H_1}/m^2_{L_3}$ at the weak scale as a function of the same ratio at the unification scale M_{GUT}. The shaded region is allowed implying an upper bound on the ratio at the weak scale for a given value of the ratio at the GUT scale. In particular, if we consider the universality $(m^2_{H_1}/m^2_{L_3})_{GUT} = 1$ case, we find an upper bound for $(m^2_{H_1}/m^2_{L_3})_{weak}$, implying that Δm^2 is negative and its magnitude has a minimum value $|\Delta m^2|_{min}$ which appears in Fig. 7. From this latest figure we see that for $\tan\beta = 3$ and $|\epsilon_3/\mu| = 1$ the minimum value of the neutrino mass is 2×10^{-3} MeV, and if no cancellation is allowed ($\Delta B = 0$ and $\Delta m^2 = \Delta m^2_{min}$), the neutrino mass is 4×10^{-2} MeV. The one–step approximation for Δm^2 give us a good estimation of this latest value. Smaller neutrino masses are obtained either with smaller values of ϵ_3 or accepting a larger cancellation. We also see from Fig. 9 that a relaxation of universality of 0.5% or more is enough to make $(m^2_{H_1}/m^2_{L_3})_{weak} = 0$ available, meaning that smaller neutrino masses are attainable without having to rely on a cancellation between the Δm^2 and ΔB terms or small values of ϵ_3.

However as we increase $\tan\beta$ the upper bound on $(\Delta m^2)_{weak}$ decreases, and thus, the required non universality between m_{H_1} and m_{L_3} at unification scale grows drastically. In Fig. 10 we show the ratio $m^2_{H_1}/m^2_{L_3}$ at the weak scale as a function of $\tan\beta$. We

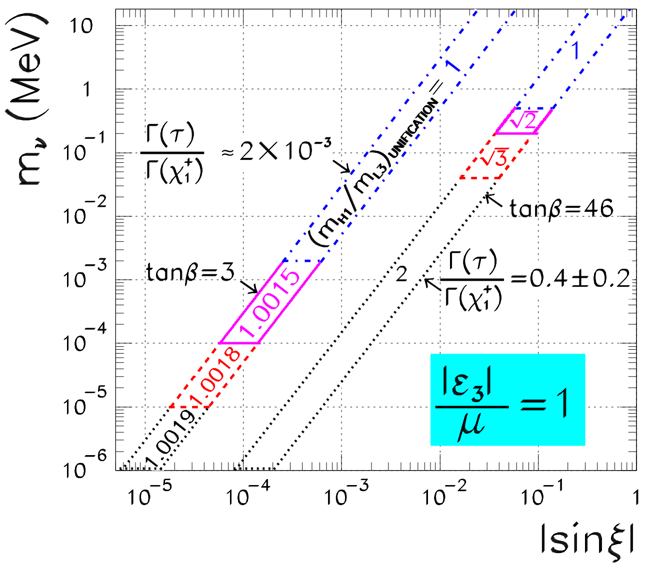


Figure 11: Minimum value of the tau neutrino mass as a function of $\sin\xi$ for different values of m_{H_1}/m_{L_3} at the GUT scale and two values of $\tan\beta$. The ratio ϵ_3/μ is fixed to the shown value, leading to a nearly constant value for $\Gamma(\tilde{t}_1 \to b\,\tau)/\Gamma(\tilde{t}_1 \to b\,\tilde{\chi}_1^+)$. Only a cancellation within one decade between the two term contributing to the tau neutrino mass is allowed.

appreciate clearly the growing of $|\Delta m^2|_{min}$ with $\tan\beta$. We remind the reader that this kind of non–universality in the soft terms is not uncommon in string models [43], or GUT models based on $SU(5)$ [44] or $SO(10)$ [45] for example.

The effect of non–universality it is also explored in Fig. 11 where it is shown the relation between the neutrino mass and the parameter $\sin\xi$ for $\epsilon_3/\mu = 1$. Two different bands are shown: one for $\tan\beta = 3$ and $\Gamma(\tau)/\Gamma(\tilde{\chi}_1^+) = 2\times10^{-3}$, and a second one for $\tan\beta = 46$ and $\Gamma(\tau)/\Gamma(\tilde{\chi}_1^+) = 0.4 \pm 0.2$. Inside the bands it is indicated the degree of universality at the GUT scale that is necessary. For example, in order to have neutrino masses of the order of eV for $\tan\beta = 3$, $m_{H_1}^2$ needs to be at least 0.2% larger than $m_{L_3}^2$. Similarly, for $\tan\beta = 46$ we need a $m_{H_1}^2$ twice as large as $m_{L_3}^2$ at the GUT scale in order to have neutrino masses of 1 eV. We stress the fact that for Fig. 11 we have accepted cancellation within 1 decade only, and we consider this a very conservative choice for reasons explained before.

6 Conclusions

We have studied the decays of the lightest stop quark in SUGRA models with and without R-parity. We have improved the calculation for the decay $\tilde{t}_1 \to c\,\tilde{\chi}^0$ by numerically solving the renormalization group equations (RGE's) of the MSSM including all generation mixing in the RGE's for Yukawa couplings as well as soft SUSY breaking parameters. The width decay is in general one order of magnitude smaller than the one obtained in the usual one–step approximation. This result will therefore enlarge the regions of the parameter space where the four–body decays of lightest stop dominate over the decay into a charm quark and the lightest neutralino. It will also affect the present experimental lower bound on the $\tilde{t}_1$ mass even in the R-parity conserving case [46]. In the RPV case, the lightest stop can be the LSP decaying with 100% rate into bottom and tau. We have shown that the decay mode $\tilde{t}_1 \to b\,\tau$ dominate over $\tilde{t}_1 \to c\,\tilde{\chi}^0$ even for neutrino masses in the range suggested by the simplest oscillation interpretation of the Super-Kamiokande atmospheric neutrino data. This result has a strong impact on the searches for the stop at LEP [47] and TEVATRON [48] where it is assumed that the decay mode $\tilde{t}_1 \to c\,\tilde{\chi}^0$ is the main decay channel. In addition to the signal of two jets and two taus present when the two produced stops decays through the R-parity violating channel, one expects a plethora of exotic high–multiplicity fermion events since in RPV the neutralino will decay inside the detector even for neutrino masses in the range suggested by the ν_μ to ν_τ oscillation interpretation of the atmospheric neutrino anomaly [49]. We have also compared the decay $\tilde{t}_1 \to b\,\tau$ with the R-parity and flavor conserving mode $\tilde{t}_1 \to b\,\tilde{\chi}^+$ and shown that the rate of the former can be comparable or even bigger than the second one if the tau neutrino mass is in the MeV range and $\tan\beta$ is large. However it is possible to have a sizeable branching of $\tilde{t}_1 \to b\,\tau$ in the case of suppressed tree level neutrino mass if larger cancelations between the two terms that contributes to $\sin\xi$ are accepted, or in some regions of the parameter space of non-univesal SUGRA models with $(m_{H_1}^2/m_{L_3}^2)_{GUT} \neq 1$. A detailed analysis of the detectability prospects of the related signatures at present and future accelerators lies outside of the scope of the present paper and it will be shown elsewhere [50]

Acknowledgments

We thank J. Ferrandis, O. J. P. Eboli and W. Porod for useful discussions. This work was supported by DGICYT under grants PB95-1077 and by the TMR network grant ERBFMRXCT960090 of the European Union. M.A.D. was supported by a DGICYT postdoctoral grant and by the U.S. Department of Energy under contract number DE-FG02-97ER41022. D.R. was supported by Colciencias fellowship.

A Feynman Rules

In this Appendix we derive the Feynman rules $F_j^0 q_i \tilde{q}_k$ (involving a neutralino/tau-neutrino, a quark, and a squark) and $F_j^{\pm} q_i \tilde{q}'_k$ (involving a chargino/tau, a quark, and a squark of different electric charge) in the case of three generations and RPV in the third generation. This is a generalization of the Feynman rules contained in [40], which are done for the R–Parity conserving MSSM and for one generation of quarks and squarks.

Following [41] we work in a quark interaction basis where $d_{L,R} = d^0_{L,R}$, $u_L = K u^0_L$, and $u_R = u^0_R$ (we denote q and q^0 the mass and current eigenstates respectively), as opposed to ref. [42] where a more general basis is used. In addition, we implement the notation $\tilde{q}_{L,R} \equiv \tilde{q}^0_{L,R}$ for the interaction basis.

The starting point is the following piece of the Lagrangian

$$
\begin{aligned}
\mathcal{L}_{u\tilde{u}F^0} \quad = \quad & -g\bar{u}^0_i \Bigg\{ \sqrt{2}\left[\sin\theta_W e_U N'_{J1} + \frac{1}{\cos\theta_W}(\tfrac{1}{2} - e_U \sin^2\theta_W) N'_{J2}\right] \tilde{u}_{iL} \\
& + \frac{m^{0U}_{ij}}{\sqrt{2} m_W \sin\beta \sin\theta} N'_{J4} \tilde{u}_{jR} \Bigg\} P_R F^0_J \\
& + g\bar{u}^0_i \Bigg\{ \sqrt{2}\left[\sin\theta_W e_U N'^*_{J1} + \frac{1}{\cos\theta_W}(-e_U \sin^2\theta_W) N'^*_{J2}\right] \tilde{u}_{iR} \\
& - \frac{m^{0U\dagger}_{ij}}{\sqrt{2} m_W \sin\beta \sin\theta} N'^*_{J4} \tilde{u}_{jL} \Bigg\} P_L F^0_J + \text{h.c.}
\end{aligned}
\tag{A.1}
$$

written in the quark interaction basis. The 5×5 matrix N' diagonalizes the neutralino/neutrino mass matrix in the $(\tilde{\gamma}, \tilde{Z}, \tilde{H}^0_1, \tilde{H}^0_2, \nu_\tau)$ basis as defined in [34], with the index $J = 1...5$. The 3×3 up–type quark mass matrix m^{0U} is not diagonal, with the indexes $i, j = 1, 2, 3$.

In order to write the above Lagrangian with mass eigenstates we use the basic relations mentioned before, in particular, $u^0_{iL} = (K^\dagger)^{ij} u_{jL}$, which implies that $\bar{u}^0_{iL} = \bar{u}_{jL} K^{ji}$. We need the following relations:

$$
\begin{aligned}
\bar{u}^0_{iL}\tilde{u}_{iL} &= \bar{u}_{iL}\left(\Gamma^*_{UL}(K^\dagger)^*\right)^{ki}\tilde{u}_k \\
\bar{u}^0_{iL} m^{0U}_{ij}\tilde{u}_{jR} &= \bar{u}_{iL}(\Gamma^*_{UR} m^U)^{ki}\tilde{u}_k \\
\bar{u}_{iR}\tilde{u}_{iR} &= \bar{u}_{iR}\Gamma^{*ki}_{UR}\tilde{u}_k
\end{aligned}
\tag{A.2}
$$

$$\bar{u}_{iR} m^{0U}_{ij} \tilde{u}_{jL} \quad = \quad \bar{u}_{iR}(\Gamma^*_{UL} K^* m^U)^{ki} \tilde{u}_k$$

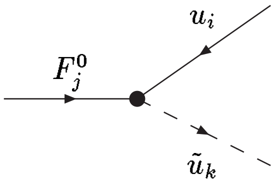
where $i, j = 1, 2, 3$ label the quark flavours, $k = 1 \ldots 6$ labels the squarks, and $m^U \equiv \mathrm{diag}\,\{m_u, m_c, m_t\}$ is the diagonal up–type quark mass matrix. In this way, the Lagrangian in Eq. (A.1) can be written as

$$\mathcal{L}_{u\tilde{u}F^0} = -g\bar{u}_i[(\sqrt{2}G^{*jki}_{0UL} + H^{*jki}_{0UR})P_R - (\sqrt{2}G^{*jki}_{0UR} - H^{*jki}_{0UL})P_L]F^0_j \tilde{u}_k + \text{h.c.} \qquad \text{(A.3)}$$

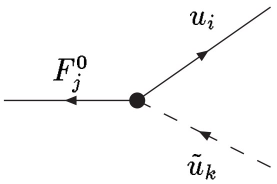
where the different couplings are

$$\begin{aligned}
G^{jki}_{0UL} &= \left[\sin\theta_W e_U N'^*_{j1} + \frac{1}{\cos\theta_W}(\tfrac{1}{2} - e_U \sin^2\theta_W)N'^*_{j2}\right]\left(\Gamma_{UL}K^\dagger\right)^{ki} \\
G^{jki}_{0UR} &= \left[\sin\theta_W e_U N'_{j1} + \frac{1}{\cos\theta_W}(-e_U \sin^2\theta_W)N'_{j2}\right]\Gamma^{ki}_{UR} \\
H^{jki}_{0UL} &= N'_{j4}(\Gamma_{UL}K^\dagger \hat{h}_U)^{ki} \\
H^{jki}_{0UR} &= N'^*_{j4}(\Gamma_{UR}\hat{h}_U)^{ki}
\end{aligned} \qquad \text{(A.4)}$$

and $\hat{h}_U \equiv \mathrm{diag}\,(m_u, m_c, m_t)/(\sqrt{2}m_W \sin\beta \sin\theta)$. Graphically, the $F^0_j u_i \tilde{u}_k$ Feynman rules are given by



$$-ig[(\sqrt{2}G^{jki}_{0UL} + H^{jki}_{0UR})P_L - (\sqrt{2}G^{jki}_{0UR} - H^{jki}_{0UL})P_R]$$

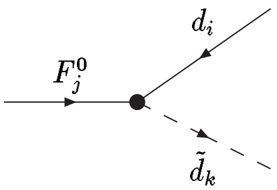


$$-ig[(\sqrt{2}G^{*jki}_{0UL} + H^{*jki}_{0UR})P_R - (\sqrt{2}G^{*jki}_{0UR} - H^{*jki}_{0UL})P_L]$$

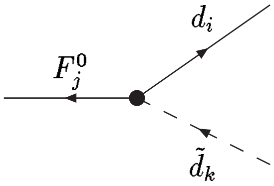
The analogous Feynman rules in the MSSM are obtained by replacing $F^0_i \to \tilde{\chi}^0_i$, by interpreting the matrix N as the usual 4×4 neutralino mixing matrix, and by setting $\theta = \pi/2$ in the formula for the Yukawa couplings.

Similarly, replacing all $u(\tilde{u})$ by $d(\tilde{d})$ in Eq. (A.1) and starting from Eq.(5.3) of [40] we can obtain the complete Feynman rules for the neutralino/tau-neutrino and chargino/tau with quarks and squarks. The results, that complements the obtained in [41], are

Neutralino–(d)quark–(d)squark



$$-ig[(\sqrt{2}G^{jki}_{0DL} + H^{jki}_{0DR})P_L - (\sqrt{2}G^{jki}_{0DR} - H^{jki}_{0DL})P_R]$$

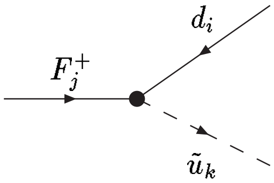


$$-ig[(\sqrt{2}G_{0DL}^{*jki} + H_{0DR}^{*jki})P_R - (\sqrt{2}G_{0DR}^{*jki} - H_{0DL}^{*jki})P_L]$$

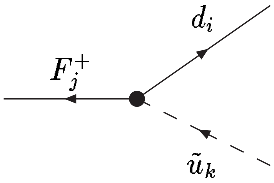
The mixing matrices G_{0D} and H_{0D} are defined as

$$\begin{aligned}
G_{0DL}^{jki} &= \left[\sin\theta_W e_D N'^{*}_{j1} + \frac{1}{\cos\theta_W}(T_{3D} - e_D \sin^2\theta_W) N'^{*}_{j2}\right]\Gamma_{DL}^{ki} \\
G_{0DR}^{jki} &= \left[\sin\theta_W e_D N'_{j1} + \frac{1}{\cos\theta_W}(-e_D \sin^2\theta_W) N'_{j2}\right]\Gamma_{DR}^{ki} \\
H_{0DL}^{jki} &= N'_{j3}(\Gamma_{DL}\hat{h}_D)^{ki} \\
H_{0DR}^{*jki} &= N'^{*}_{j3}(\Gamma_{DR}\hat{h}_D)^{ki} \qquad \text{(A.5)}
\end{aligned}$$

Chargino/tau–(d)quark–(u)squark



$$-ig(-C^{-1})[(G_{UL}^{jki} - H_{UR}^{jki})P_L - H_{UL}^{jki}P_R]$$

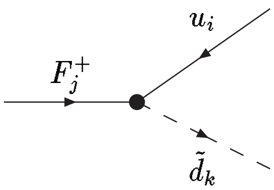


$$-ig[(G_{UL}^{*jki} - H_{UR}^{*jki})P_R - H_{UL}^{*jki}P_L]C$$

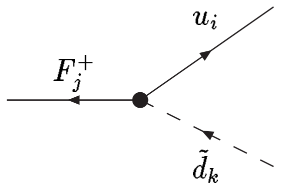
where C is the charge conjugation matrix (in spinor space) and the mixing matrices G_U and H_U are defined as

$$\begin{aligned}
G_{UL}^{jki} &\equiv V_{j1}^{*}\Gamma_{UL}^{ki}, \qquad H_{UL}^{jki} \equiv U_{j2}^{*}(\Gamma_{UL}\hat{h}_D)^{ki}, \\
H_{UR}^{jki} &\equiv V_{j2}^{*}(\Gamma_{UR}\hat{h}_U K)^{ki}, \qquad \text{(A.6)}
\end{aligned}$$

Chargino/tau–(u)quark–(d)squark



$$-ig(-C^{-1})[(G_{DL}^{jki} - H_{DR}^{jki})P_L - H_{DL}^{jki}P_R]$$



$$-ig[(G_{DL}^{*jki} - H_{DR}^{*jki})P_R - H_{DL}^{*jki} P_L]C$$

where the mixing matrices G_D and H_D are defined as

$$\begin{aligned} G_{DL}^{jki} &\equiv U_{j1}^*(\Gamma_{DL}K^\dagger)^{ki}, \qquad H_{DL}^{jki} \equiv V_{j2}^*(\Gamma_{DL}K^\dagger\hat{h}_U)^{ki}, \\ H_{DR}^{jki} &\equiv U_{j2}^*(\Gamma_{DR}\hat{h}_D K^\dagger)^{ki}, \end{aligned} \tag{A.7}$$

In order to derive the decays widths we write, for example Eq. (A.3) as

$$\mathcal{L}_{u\tilde{u}F^0} = g\bar{u}_i^0(f_U^{*jki}P_R + h_U^{*jki}P_L)F_j^0\tilde{u}_k + \text{h.c} \tag{A.8}$$

The result is

$$\begin{aligned} \Gamma(\tilde{q}_k \to q_i + F_j^0) &= \frac{g^2\lambda^{1/2}(m_{\tilde{q}_k}^2, m_{q_i}^2, m_{F_j^0}^2)}{16\pi m_{\tilde{q}_k}^3}\Big[-4h_Q^{jki}f_Q^{jki}m_{q_i}m_{F_j^0} \\ &\quad + \Big((h_Q^{jki})^2 + (f_Q^{jki})^2\Big)\Big(m_{\tilde{q}_k}^2 - m_{q_i}^2 - m_{F_j^0}^2\Big)\Big] \end{aligned} \tag{A.9}$$

$$\begin{aligned} \Gamma(\tilde{q}_k \to q_i' + F_j^\pm) &= \frac{g^2\lambda^{1/2}(m_{\tilde{q}_k}^2, m_{q_i'}^2, m_{F_j^\pm}^2)}{16\pi m_{\tilde{q}_k}^3}\Big[-4l_Q^{jki}H_{QL}^{jki}m_{q_i'}m_{F_j^\pm} \\ &\quad + \Big((l_Q^{jki})^2 + (H_{QL}^{jki})^2\Big)\Big(m_{\tilde{q}_k}^2 - m_{q_i'}^2 - m_{F_j^\pm}^2\Big)\Big] \end{aligned} \tag{A.10}$$

where $Q = U, D$ refers to $\tilde{q}$ and

$$f_Q^{jki} = -(\sqrt{2}G_{0QL}^{jki} + H_{0QR}^{jki}) \tag{A.11}$$

$$h_Q^{jki} = \sqrt{2}G_{0QR}^{jki} - H_{0QL}^{jki} \tag{A.12}$$

$$l_Q^{jki} = H_{QR}^{jki} - G_{QL}^{jki} \tag{A.13}$$

with the G and H couplings defined earlier in this appendix.

References

[1] H. Haber and G. Kane, Phys. Rep. **117** (1985) 75.

[2] H. P. Nilles, Phys. Rep. **110** (1984) 1.

[3] For a recent review see J.W.F. Valle, Supergravity Unification with Bilinear R Parity Violation, Proceedings of PASCOS98, ed. P. Nath, W. Scientific, [hep-ph 9808292]; J. W. F. Valle, in *Physics Beyond the Standard Model*, lectures given at the *VIII Jorge Andre Swieca Summer School* (Rio de Janeiro, February 1995) and at *V Taller Latinoamericano de Fenomenologia de las Interacciones Fundamentales* (Puebla, Mexico, October 1995); hep-ph/9603307.

[4] C. S. Aulakh, R.N. Mohapatra, Phys. Lett. **B119** (1982) 136.

[5] L. Hall, M. Suzuki, Nucl. Phys. **B231** (1984) 419.

[6] G. G. Ross, J. W. F. Valle, Phys. Lett. **151B** (1985) 375; J. Ellis, G. Gelmini, C. Jarlskog, G. G. Ross, J. W. F. Valle, Phys. Lett. **150B** (1985) 142.

[7] A. Santamaria, J.W.F. Valle, Phys. Lett. **195B** (1987) 423; Phys. Rev. **D39** (1989) 1780; Phys. Rev. Lett. **60** (1988) 397.

[8] S. Dimopoulos, L.J. Hall, Phys. Lett. **207B** (1988) 210; E. Ma, D. Ng, Phys. Rev. **D41** (1990) 1005; V. Barger, G. F. Giudice, T. Han, Phys. Rev. **D40** (1989) 2987; T. Banks, Y. Grossman, E. Nardi, Y. Nir, Phys. Rev. **D52** (1995) 5319; M. Nowakowski, A. Pilaftsis, Nucl. Phys. **B461** (1996) 19; G. Bhattacharyya, D. Choudhury, K. Sridhar, Phys. Lett. **B349** (1995) 118; B. de Carlos, P. L. White, Phys. Rev. **D54** (1996) 3427.

[9] A Masiero, J. W. F. Valle, Phys. Lett. **B251** (1990) 273; J. C. Romão, C. A. Santos, and J. W. F. Valle, Phys. Lett. **B288** (1992) 311 J.C. Romao, A. Ioannisian and J.W. F. Valle, Phys. Rev. **D55**, 427 (1997), hep-ph/9607401.

[10] G. Giudice, A. Masiero, M. Pietroni, A. Riotto, Nucl. Phys. **B396** (1993) 243; M. Shiraishi, I. Umemura, K. Yamamoto, Phys. Lett. **B313** (1993) 89; see also I. Umemura, K. Yamamoto, Nucl. Phys. **B423** (1994) 405.

[11] P. Nogueira, J. C. Romão, J. W. F. Valle, Phys. Lett. **B251** (1990) 142; R. Barbieri, L. Hall, Phys. Lett. **B238** (1990) 86; M. C. Gonzalez-García, J W F Valle, Nucl. Phys. **B355** (1991) 330; J. Romão, J. Rosiek and J. W. F. Valle, Phys. Lett. **B351** (1995) 497; J. C. Romão, N. Rius, J. W. F. Valle, Nucl. Phys. **B363** (1991) 369;

J. C. Romão and J. W. F. Valle Phys. Lett. **B272** (1991) 436; Nucl. Phys. **B381** (1992) 87.

[12] J. W. F. Valle, Phys. Lett. **B196** (1987) 157.

[13] V. Berezinskii and J.W.F. Valle, Phys. Lett. **B318** (1993) 360; A. D. Dolgov, S. Pastor, and J.W.F. Valle astro-ph/9506011.

[14] P. Abreu *et al.* [DELPHI Collaboration], CERN-EP-99-049; J.-F. Grivaz, Rapporteur Talk, International Europhysics Conference on High Energy Physics, Brussels, 1995; ALEPH collaboration, Phys. Lett. **B373** (1996) 246-260; H. Nowak and A. Sopczak, L3 Note 1887, Jan. 1996 ; S. Asai and S. Komamiya, OPAL Physics Note PN-205, Feb. 1996

[15] B. Abbott *et al.* [D0 Collaboration], hep-ex/9902013; D0 collaboration, Phys. Rev. Lett. **75** (1995) 618; and references therein, CDF collaboration, Phys. Rev. Lett. **69** (1992) 3439; UA2 collaboration, Phys. Lett. **B235** (1990) 363.

[16] A. Bartl, W. Porod, M. A. Garcia-Jareno, M. B. Magro, J. W. F. Valle, W. Majerotto, Phys. Lett. **B384** (1996) 151.

[17] F. de Campos, M. A. García-Jareño, A. Joshipura, J. Rosiek, J. W. F. Valle, Nucl. Phys. **B451** (1995) 3-15; H. Hempfling, Nucl. Phys. **B478** (1996) 3.

[18] See for example: Stephen P. Martin. Published in Perspectives on supersymmetry, edited G.L. Kane, World Scientific, 1998, 88pp. hep-ph/9709356, A. Bartl, *et al*, Published in the Proceedings of the 1996 DPF/DPB Summer Study on New Directions For High-Energy Physics, Edited by D.G. Cassel, L. Trindle Gennari, R.H. Siemann, 1997, p. 693-705. See also http://www.slac.stanford.edu/pubs/snowmass96/.

[19] M.A. Díaz, J. Ferrandis, J.C. Romão, and J.W.F. Valle, *Phys. Lett B* **453**, 263 (1999); M.A. Díaz, E. Torrente–Lujan, J.W.F. Valle, *Nucl. Phys. B* **551**, 78 (1999), L. Navarro, W. Porod and J.W. Valle, hep-ph/9903474.

[20] S. Roy and B. Mukhopadhyaya, Phys. Rev. **D55** (1997) 7020; ibid hep-ph/9903418; A. Datta, B. Mukhopadhyaya and S. Roy, hep-ph/9905549; T. Feng, hep-ph/980650; ibid hep-ph/9808379. C. Chang and T. Feng, hep-ph/9901260.

[21] J. Ferrandis, hep-ph/9810371, Accepted for publication in Physical Review D; M.A. Diaz, J. Ferrandis, J.C. Romao and J.W.F. Valle, hep-ph/9906343.

[22] F. Borzumati, J. Kneur and N. Polonsky, hep-ph/9905443.

[23] Y. Grossman and H.E. Haber, Phys. Rev. **D59** (1999) 093008; ibid Phys. Rev. Lett. **78** (1997) 3438.

[24] M. Bisset, O.C. Kong, C. Macesanu and L.H. Orr, Phys. Lett. **B430** (1998) 274; ibid hep-ph/9811498; ibid hep-ph/9907359.

[25] S. Davidson and J. Ellis, Phys. Lett. **B390** (1997) 210; S. Davidson and J. Ellis, Phys. Rev. **D56** (1997) 4182; S. Davidson, Phys. Lett. **B439** (1998) 63.

[26] E. Nardi, Phys. Rev. **D55** (1997) 5772.

[27] H.P. Nilles and N. Polonsky Nucl. Phys. **B484** (1997) 33.

[28] I.-H. Lee, Phys. Lett. **138B** (1984) 121; A.S. Joshipura and M. Nowakowski, Phys. Rev. **D51** (1995) 21; S. Roy and B. Mukhopadhyaya, Phys. Rev. **D55** (1997) 7020; F.M. Borzumati, Y. Grossman, E. Nardi and Y. Nir, Phys. Lett. **B384** (1996) 123.

[29] M.A. Díaz, J.C. Romao and J.W. Valle, Nucl. Phys. **B524**, 23 (1998).

[30] B. Mukhopadhyaya, S. Roy and F. Vissani, Phys. Lett. **B443** (1998) 191;V. Bednyakov, A. Faessler and S. Kovalenko, Phys. Lett. **B442** (1998) 203; O.C. Kong, Mod. Phys. Lett. **A14** (1999) 903; E.J. Chun, *et al*, Nucl. Phys. **B544** (1999) 89; S.Y. Choi, *et al*, hep-ph/9903465; A.S. Joshipura and S.K. Vempati, hep-ph/9903435; ibid hep-ph/9808232; R. Adhikari and G. Omanovic, hep-ph/9802390; D.E. Kaplan and A.E. Nelson, hep-ph/9901254; Y. Grossman and H.E. Haber, hep-ph/9906310.

[31] ALEPH Collaboration, Eur. Phys. J.**C2** (1998) 395 .

[32] M.A. Díaz, in *Beyond the Standard Model: From Theory to Experiment*, Proceedings of Valencia 97, ed. I. Antoniadis, L. E. Ibanez and J. W. F. Valle, W. Scientific, p 188 [hep-ph/9802407]

[33] S. Rakshit, G. Bhattacharyya and A. Raychaudhuri, Phys. Rev. **D59** (1999) 091701

[34] A. Akeroyd, M.A. Díaz, J. Ferrandis, M.A. Garcia–Jareño, and J.W.F. Valle, Nucl. Phys. **B529** (1998) 3.

[35] A. Bartl, W. Majerotto and W. Porod, Zeit. fur Physik **C64** (1994) 499.

[36] K. I. Hikasa and M. Kobayashi, Phys. Rev. **D36** (1987) 724

[37] H. Baer, M. Drees, R. Godbole, J.F. Gunion, X. Tata, Phys. Rev. **D44** (1991) ; T. Kon, T. Nonaka, Phys. Lett. **B319** (1993) 355; T. Kon, T. Kobayashi, S. Kitamura, K. Nakamura and S. Adachi, Zeit. fur Physik **C61** (1994) 239; T. Kon, T. Nonaka, ITP-SU-94-02, hep-ph/9404230; W. Porod and T. Wohrmann, Phys. Rev. **D55** (1997) 2907; W. Porod, Phys. Rev. **D59** (1999) 095009.

[38] W. Porod, Ph.D thesis, hep-ph/9804208 and references therein.

[39] A. G. Akeroyd, M.A. Díaz, J.W.F. Valle, Phys. Lett. **B441** (1998) 224.

[40] J.F. Gunion, and H.E. Haber, Nucl. Phys. **B272** (1986) 1.

[41] S. Bertolini, F. Borzumati, A. Masiero, and G. Ridolfi, Nucl. Phys. **B353** (1991) 591.

[42] J. Rosiek, Phys. Rev. **D41** (1990) 3464.

[43] A. Brignole, L.E. Ibañez, and C. Muñoz, hep-ph/9707209.

[44] N. Polonsky and A. Pomarol, *Phys. Rev. Lett.* **73**, 2292 (1994).

[45] H. Murayama, M. Olechowski, and S. Pokorski, *Phys. Lett. B* **371**, 57 (1996); H. Baer, M.A. Díaz, and J. Ferrandis, hep-ph/9907211.

[46] C. Boehm, A. Djouadi and Y. Mambrini, hep-ph/9907428.

[47] G. Abbiendi *et al.* [OPAL Collaboration], Phys. Lett. **B456** (1999) 95

[48] C. Holck [CDF Collaboration], Talk given at American Physical Society (APS) Meeting of the Division of Particles and Fields (DPF 99), Los Angeles, CA, 5-9 Jan 1999. hep-ex/9903060.

[49] J.C. Romao, M.A. Diaz, M. Hirsch, W. Porod and J.W. Valle, hep-ph/9907499.

[50] Fernando de Campos, *et al*, in preparation