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Charged Higgs Boson and Stau Phenomenology in the Simplest R–Parity Breaking Model

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Abstract

We consider the charged scalar boson phenomenology in the simplest effective low-energy R-parity breaking model characterized by a bilinear violation of R-parity in the superpotential. This induces a mixing between staus and the charged Higgs boson. We show that the charged Higgs boson mass can be lower than expected in the MSSM, even before including radiative corrections. We also study the charged scalar boson decay branching ratios and show that the R-parity violating decay rates can be comparable or even bigger than the R-parity conserving ones. Moreover, if the stau is the LSP it will have only decays into standard model fermions. These features could have important implications for charged supersymmetric scalar boson searches at future accelerators.

1 Introduction

A lot of emphasis has been put into the phenomenological study of the supersymmetric Higgs boson sector [1]. However, so far most of these phenomenological studies have been made in the framework of the Minimal Supersymmetric Standard Model (MSSM) [2, 3] with conserved R-parity [4]. R-parity is a discrete symmetry assigned as $R_p = (-1)^{(3B+L+2S)}$, where L is the lepton number, B is the baryon number and S is the spin of the state. If R-parity is conserved all supersymmetric particles must always be pair-produced, while the lightest of them must be stable. In particular, supersymmetric Higgs scalar bosons must decay only to normal standard model particles or to pairs of lowest-lying supersymmetric particles, which are usually heavy. On the other hand staus decay only to supersymmetric states, like a neutralino and a tau lepton.

The study of alternative supersymmetric scenarios where the effective low energy theory violates R-parity [5] has recently received a lot of attention [6] both due to its phenomenological interest, as well as due to the intrinsic importance of investigating the issue of R-parity breaking at a deeper level.

It is well–known that the simplest supersymmetric extension of the Standard Model violates R-parity through a set of cubic superpotential terms involving a very large number of arbitrary Yukawa couplings. Although highly constrained by proton stability, many of such scenarios could still be viable. Nevertheless their systematic study at a phenomenological level is hardly possible, due to the enormous number of parameters present, in addition to those of the MSSM.

As with other fundamental symmetries, it could well be that R-parity is a symmetry at the Lagrangian level but is broken by the ground state. In order to comply with LEP precision measurements of the invisible Z decay width these models require the introduction of $SU(2) \otimes U(1)$ singlet superfields [7]. Such scenarios provide a very *systematic* way to include R parity violating effects, automatically consistent with low energy *baryon number conservation*. They have many added virtues, such as the possibility of providing a dynamical origin for the breaking of R-parity, through radiative corrections, similar to the electroweak symmetry [8]. The simplest truncated version of such a model, in which the violation of R-parity is effectively parametrized by a bilinear superpotential term $\epsilon_i \widehat{L}_i^a \widehat{H}_2^b$ has been widely discussed [9, 10]. It has also been shown recently [10] that this model is consistent with minimal N=1 supergravity unification with radiative breaking of the electroweak symmetry and universal scalar and gaugino masses. This one-parameter extension of the MSSM-SUGRA model therefore provides the simplest reference model for the breaking of R-parity and constitutes a consistent truncation of the complete dynamical models with spontaneous R-parity breaking proposed previously [7]. In this case there is no physical Goldstone boson, the Majoron, associated to the spontaneous breaking of R-parity, since in this effective truncated model the superfield content is exactly the standard one of the MSSM. Formulated as an effective theory at the weak scale, the model contains only two new parameters in addition to those of the MSSM. Therefore our model provides also the simplest parametrization of R-parity breaking effects. In contrast to models with tri-linear R-parity breaking couplings, it leads to a very restrictive and systematic pattern of R-parity violating interactions, which can be taken as a reference model.

In this paper we focus on the phenomenology of the charged scalar boson sector of the simplest R-parity breaking model. This complements a previous study of the electrically neutral sector [11]. We show that

1. the mass of the charged Higgs boson can be lower than expected in the MSSM, even before including radiative corrections,
2. if the stau is the LSP it will have only R-parity violating decay channels into standard model fermions,
3. the branching ratio for the R-parity violating charged Higgs boson decays can be comparable or even bigger than the R-parity conserving ones.

We illustrate how these features arising from the charged scalar boson sector of the simplest R-parity breaking model could play an important role in designing charged supersymmetric scalar boson searches at e^+e^- colliders such as LEP II. For example they can give rise to striking signatures consisting of high multiplicity events, such as di–tau + 4 jets + missing energy or 4 taus + 4 jets. Such processes, forbidden in the MSSM, are expected to have high rates and negligible background. As for hadron colliders, we also can have very high leptonic multiplicity events such as six leptons of which at least two are taus, plus missing momentum. This should be easy to see at the LHC, due again to the negligible standard model background.

2 The Model

The supersymmetric Lagrangian is specified by the superpotential W given by *

$$W = \varepsilon_{ab}\left[h_U^{ij}\widehat{Q}_i^a\widehat{U}_j\widehat{H}_2^b + h_D^{ij}\widehat{Q}_i^b\widehat{D}_j\widehat{H}_1^a + h_E^{ij}\widehat{L}_i^b\widehat{R}_j\widehat{H}_1^a - \mu\widehat{H}_1^a\widehat{H}_2^b + \epsilon_i\widehat{L}_i^a\widehat{H}_2^b\right] \tag{1}$$

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

*We are using here the notation of refs. [3] and [12].

μ and ϵ_i are parameters with units of mass. The last term in eq. (1) is the only R–parity violating term.

Supersymmetry breaking is parametrized with a set of soft supersymmetry breaking terms which do not introduce quadratic divergences to the unrenormalized theory [13]

$$\begin{aligned} V_{soft} &= M_Q^{ij2}\widetilde{Q}_i^{a*}\widetilde{Q}_j^a + M_U^{ij2}\widetilde{U}_i^*\widetilde{U}_j + M_D^{ij2}\widetilde{D}_i^*\widetilde{D}_j + M_L^{ij2}\widetilde{L}_i^{a*}\widetilde{L}_j^a + M_R^{ij2}\widetilde{R}_i^*\widetilde{R}_j \\ &+m_{H_1}^2 H_1^{a*}H_1^a + m_{H_2}^2 H_2^{a*}H_2^a - \left[\tfrac{1}{2}M_s\lambda_s\lambda_s + \tfrac{1}{2}M\lambda\lambda + \tfrac{1}{2}M'\lambda'\lambda' + h.c.\right] \\ &+\varepsilon_{ab}\left[A_U^{ij}h_U^{ij}\widetilde{Q}_i^a\widetilde{U}_jH_2^b + A_D^{ij}h_D^{ij}\widetilde{Q}_i^b\widetilde{D}_jH_1^a + A_E^{ij}h_E^{ij}\widetilde{L}_i^b\widetilde{R}_jH_1^a - B\mu H_1^aH_2^b + B_2\epsilon_i\widetilde{L}_i^aH_2^b\right], \end{aligned} \tag{2}$$

and again, the last term in eq. (2) is the only R–parity violating term.

Following previous discussions we will focus for simplicity on the case of one generation, namely the third [11, 14]. In contrast we will keep in our discussion the theory as defined at low energies by the most general set of soft-breaking masses, tri-linear and bilinear soft-breaking parameters, gaugino masses and the Higgs superfield mixing parameter μ.

The electroweak symmetry is broken when the two Higgs doublets H_1 and H_2, and the third component of the left slepton doublet $\widetilde{L}_3$ acquire vacuum expectation values. We introduce the notation:

$$H_1 = \begin{pmatrix} \frac{1}{\sqrt{2}}[\chi_1^0 + v_1 + i\varphi_1^0] \\ H_1^- \end{pmatrix}, \qquad H_2 = \begin{pmatrix} H_2^+ \\ \frac{1}{\sqrt{2}}[\chi_2^0 + v_2 + i\varphi_2^0] \end{pmatrix},$$
$$\widetilde{L}_3 = \begin{pmatrix} \frac{1}{\sqrt{2}}[\tilde{\nu}_\tau^R + v_3 + i\tilde{\nu}_\tau^I] \\ \tilde{\tau}^- \end{pmatrix}. \tag{3}$$

Note that the W boson acquires a mass $m_W^2 = \frac{1}{4}g^2v^2$, where $v^2 \equiv v_1^2 + v_2^2 + v_3^2 = (246\text{ GeV})^2$. We introduce the following notation in spherical coordinates for the vacuum expectation values (VEV):

$$\begin{aligned} v_1 &= v\sin\theta\cos\beta \\ v_2 &= v\sin\theta\sin\beta \\ v_3 &= v\cos\theta \end{aligned} \tag{4}$$

which preserves the MSSM definition $\tan\beta = v_2/v_1$. In the MSSM limit, where $\epsilon_3 = v_3 = 0$, the angle θ is equal to $\pi/2$.

In addition to the above MSSM parameters, our model contains three new parameters, ϵ_3, v_3 and B_2, of which only two are independent, and these may be chosen as ϵ_3 and v_3.

The full scalar potential at tree level is

$$V_{total} = \sum_i \left|\frac{\partial W}{\partial z_i}\right|^2 + V_D + V_{soft} \tag{5}$$

where z_i is any one of the scalar fields in the superpotential, V_D are the D-terms, and V_{soft} the SUSY soft breaking terms given in eq. (2).

The scalar potential contains linear terms

$$V_{linear} = t_1^0 \chi_1^0 + t_2^0 \chi_2^0 + t_3^0 \tilde{\nu}_\tau^R \,, \tag{6}$$

where

$$\begin{aligned} t_1^0 &= (m_{H_1}^2 + \mu^2) v_1 - B\mu v_2 - \mu \epsilon_3 v_3 + \tfrac{1}{8}(g^2 + g'^2) v_1 (v_1^2 - v_2^2 + v_3^2) \,, \\ t_2^0 &= (m_{H_2}^2 + \mu^2 + \epsilon_3^2) v_2 - B\mu v_1 + B_2 \epsilon_3 v_3 - \tfrac{1}{8}(g^2 + g'^2) v_2 (v_1^2 - v_2^2 + v_3^2) \,, \\ t_3^0 &= (m_{L_3}^2 + \epsilon_3^2) v_3 - \mu \epsilon_3 v_1 + B_2 \epsilon_3 v_2 + \tfrac{1}{8}(g^2 + g'^2) v_3 (v_1^2 - v_2^2 + v_3^2) \,. \end{aligned} \tag{7}$$

These $t_i^0, i = 1, 2, 3$ are the tree level tadpoles, and are equal to zero at the minimum of the potential.

Now a few theoretical comments on the model. The first refers to the choice of basis in the original Lagrangian in eq. (1). We could have used a rotated basis in which the bi-linear coupling disappears [15]. However, if we choose to do so, new R–parity violating terms appear not only in the tri-linear superpotential, in the form of a DQL term, but also in the scalar sector of the theory, i.e. the Higgs potential due to supersymmetry breaking. Authors who adopt this basis [16] have a tendency to neglect SUSY breaking in the Higgs potential, which is most crucial for our subsequent analysis. Therefore we prefer to use in our calculations the basis in which the bi-linear term is not rotated away. The final physics results are completely basis-independent [17].

Another important feature of this model is that lepton number is violated by the ϵ_3 term and by the presence of the sneutrino vacuum expectation value v_3. This induces a mass for the tau neutrino since ν_τ mixes with the neutralinos (see the Appendix). This mass turns out to be proportional to an *effective* neutralino mixing parameter $\xi \equiv (\mu v_3 + \epsilon_3 v_1)^2$ characterizing the violation of R–parity, either through v_3 or ϵ_3. One can show [17] that this parameter corresponds to the R–parity violating VEV in the rotated basis. If we stick to the simplest unified supergravity version of the model with universal boundary conditions for the soft breaking parameters, then ξ will be small since contributions arising from *gaugino* mixing will cancel, to a large extent, those from *Higgsino* mixing. This cancellation will happen automatically so that in this case m_{ν_τ} will be naturally small and radiatively calculable in terms of the bottom Yukawa coupling h_b, thus accounting naturally for the smallness of the ν_τ mass in this model. In the language of [18] one can say that universality of the soft breaking terms implies an approximate and radiatively calculable *alignment* and as a result a suppression in ξ and on m_{ν_τ}. This offers a *hybrid* scenario combining the see-saw and radiative schemes of ν mass generation. The rôle of the right-handed mass which appears in the see-saw model is played by the neutralino mass (which lies at the SUSY scale) while the rôle of the seesaw-scheme Dirac mass is played by the *effective* neutralino mixing ξ which is induced radiatively. The ν_τ mass

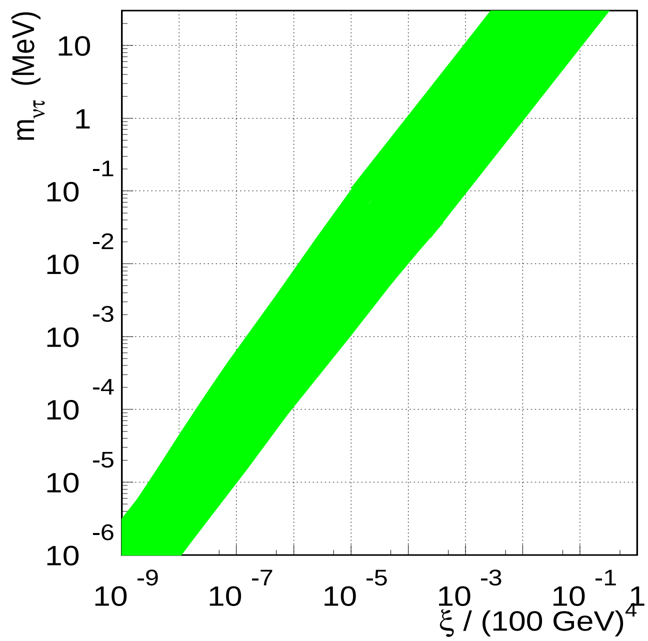

Figure 1: Tau neutrino mass versus de R–parity violating parameter ξ

induced this way is directly correlated with the magnitude of the effective parameter ξ so that R–parity violation acts as the origin for neutrino mass. In Fig. 1 we display the allowed values of m_{ν_τ}. One can see that m_{ν_τ} values can cover a very wide range, from MeV values comparable to the present LEP limit [19] down to values in the eV range, even though the individual values of v_3 and ϵ_3 can be rather large, e. g. $v_3, \epsilon_3 \sim 100$ GeV, because they are naturally "aligned", without any need for fine-tuning to get small m_{ν_τ}. The alignment follows from the universality of the soft breaking parameters at the unification scale [10]. However, for generality, in this paper we take the parameters at the weak scale independent (*i.e.* no universality assumption), and always impose $m_{\nu_\tau} \lesssim 20$ MeV which corresponds to the laboratory limit.

This brings us to a discussion about cosmology. Clearly our model leads to a tau–neutrino which can be much heavier than the limits that follow from the cosmological critical density as well as primordial nucleosynthesis would allow [20]. However, in this model the ν_τ is unstable and decays via neutral current into three lighter neutrinos [21]. In order for this mode to be efficient we estimate that m_{ν_τ} must exceed 100 keV or so. On the other hand, to avoid problems with primordial nucleosynthesis it is safer to consider masses below 1 MeV or so. In order to sharpen these estimates (which are not strict bounds) a detailed investigation is required [22].

One should bear in mind, however, that m_{ν_τ} can be as large as the present laboratory bound [19] in the more complete versions of the model in which R–parity is broken spontaneously due to sneutrino expectation values [7, 6]. This is so because such models contain a majoron, denoted as J, which opens new decay channels ν_τ into $\nu + J$ where

ν is a lighter neutrino [23] as well as new annihilation channels $\nu_\tau + \nu_\tau \to J + J$. It has been shown explicitly that the lifetimes that can be achieved in the spontaneous broken R–parity versions of the model can be sufficiently short to obey the critical density limit [24]. Moreover, it has been shown that the annihilation channel is efficient enough in order to comply the primordial nucleosynthesis bound [25], while decays may also play an important rôle [26].

Finally, a word about the magnitude of R–parity violation. It will depend to some extent on the process considered. Some R–parity violating observables turn out to be proportional to an *effective* neutralino mixing parameter $\xi \equiv (\mu v_3 + \epsilon_3 v_1)^2$ characterizing the violation of R–parity, either through v_3 or ϵ_3. An example is the mass of the tau–neutrino (see below). However, not all R–parity-violating processes are determined by ξ: some single production processes or R–parity violating decay branching ratios, such as the ones discussed in the present paper, depend separately on v_3 or ϵ_3 and can be rather large even for small ξ and m_{ν_τ}. An obvious and important example is the decay of the lightest neutralino, which is determined by ϵ_3 only since, in the rotated basis it is determined by the D Q L superpotential term only. A detailed discussion lies outside the scope of this paper [17].

3 Scalar Mass Matrices

The mass matrix of the charged scalar sector follows from the quadratic terms in the scalar potential

$$V_{quadratic} = [H_1^-, H_2^-, \tilde{\tau}_L^-, \tilde{\tau}_R^-] \boldsymbol{M^2_{S^\pm}} \begin{bmatrix} H_1^+ \\ H_2^+ \\ \tilde{\tau}_L^+ \\ \tilde{\tau}_R^+ \end{bmatrix} + \ldots \tag{8}$$

For convenience reasons we will divide this 4×4 matrix into 2×2 blocks in the following way:

$$\boldsymbol{M^2_{S^\pm}} = \begin{bmatrix} \boldsymbol{M}^2_{HH} & \boldsymbol{M}^{2T}_{H\tilde{\tau}} \\ \boldsymbol{M}^2_{H\tilde{\tau}} & \boldsymbol{M}^2_{\tilde{\tau}\tilde{\tau}} \end{bmatrix} \tag{9}$$

where the charged Higgs block is

$$\boldsymbol{M}^2_{HH} = \tag{10}$$

$$\begin{bmatrix} B\mu\frac{v_2}{v_1} + \frac{1}{4}g^2(v_2^2 - v_3^2) + \mu\epsilon_3\frac{v_3}{v_1} + \frac{1}{2}h_\tau^2 v_3^2 + \frac{t_1}{v_1} & B\mu + \frac{1}{4}g^2 v_1 v_2 \\ B\mu + \frac{1}{4}g^2 v_1 v_2 & B\mu\frac{v_1}{v_2} + \frac{1}{4}g^2(v_1^2 + v_3^2) - B_2\epsilon_3\frac{v_3}{v_2} + \frac{t_2}{v_2} \end{bmatrix}$$

and h_τ is the tau Yukawa coupling. This matrix reduces to the usual charged Higgs mass matrix in the MSSM when we set $v_3 = \epsilon_3 = 0$ and we call $m_{12}^2 = B\mu$. The stau block is given by

$$\boldsymbol{M}^2_{\tilde{\tau}\tilde{\tau}} = \tag{11}$$

$$\begin{bmatrix} \frac{1}{2}h_\tau^2 v_1^2 - \frac{1}{4}g^2(v_1^2 - v_2^2) + \mu\epsilon_3\frac{v_1}{v_3} - B_2\epsilon_3\frac{v_2}{v_3} + \frac{t_3}{v_3} & \frac{1}{\sqrt{2}}h_\tau(A_\tau v_1 - \mu v_2) \\ \frac{1}{\sqrt{2}}h_\tau(A_\tau v_1 - \mu v_2) & m_{R_3}^2 + \frac{1}{2}h_\tau^2(v_1^2 + v_3^2) - \frac{1}{4}g'^2(v_1^2 - v_2^2 + v_3^2) \end{bmatrix}$$

We recover the usual stau mass matrix again by replacing $v_3 = \epsilon_3 = 0$, nevertheless, we need to replace the expression of the third tadpole in eq. (7) before taking the limit. The mixing between the charged Higgs sector and the stau sector is given by the following 2×2 block:

$$\boldsymbol{M}_{H\tilde{\tau}}^2 = \begin{bmatrix} -\mu\epsilon_3 - \frac{1}{2}h_\tau^2 v_1 v_3 + \frac{1}{4}g^2 v_1 v_3 & -B_2\epsilon_3 + \frac{1}{4}g^2 v_2 v_3 \\ -\frac{1}{\sqrt{2}}h_\tau(\epsilon_3 v_2 + A_\tau v_3) & -\frac{1}{\sqrt{2}}h_\tau(\mu v_3 + \epsilon_3 v_1) \end{bmatrix} \tag{12}$$

and as expected, this mixing vanishes in the limit $v_3 = \epsilon_3 = 0$. The charged scalar mass matrix in eq. (9), after setting $t_1 = t_2 = t_3 = 0$, has determinant equal to zero since one of the eigenvectors corresponds to the charged Goldstone boson with zero eigenvalue.

For completeness, we give the neutral Higgs sector mass matrices. The quadratic scalar potential includes

$$V_{quadratic} = \tfrac{1}{2}[\varphi_1^0, \varphi_2^0, \tilde{\nu}_\tau^I]\boldsymbol{M}_{P^0}^2 \begin{bmatrix} \varphi_1^0 \\ \varphi_2^0 \\ \tilde{\nu}_\tau^I \end{bmatrix} + \tfrac{1}{2}[\chi_1^0, \chi_2^0, \tilde{\nu}_\tau^R]\boldsymbol{M}_{S^0}^2 \begin{bmatrix} \chi_1^0 \\ \chi_2^0 \\ \tilde{\nu}_\tau^R \end{bmatrix} + \ldots \tag{13}$$

where the CP-odd neutral scalar mass matrix is

$$\boldsymbol{M}_{P^0}^2 = \begin{bmatrix} B\mu\frac{v_2}{v_1} + \mu\epsilon_3\frac{v_3}{v_1} + \frac{t_1}{v_1} & B\mu & -\mu\epsilon_3 \\ B\mu & B\mu\frac{v_1}{v_2} - B_2\epsilon_3\frac{v_3}{v_2} + \frac{t_2}{v_2} & -B_2\epsilon_3 \\ -\mu\epsilon_3 & -B_2\epsilon_3 & \mu\epsilon_3\frac{v_1}{v_3} - B_2\epsilon_3\frac{v_2}{v_3} + \frac{t_3}{v_3} \end{bmatrix} \tag{14}$$

This matrix also has a vanishing determinant after the tadpoles are set to zero, and the zero eigenvalue corresponds to the mass of the neutral Goldstone boson. The usual MSSM mass matrix of the pseudoscalar Higgs sector is recovered in the limit $v_3 = \epsilon_3 = 0$ in the upper–left 2×2 block, and the third component corresponding to the imaginary part of the sneutrino decouples from it. Note that in this limit, the MSSM mass of the sneutrino is recovered provided we replace the expression for the tadpole in eq. (7) before taking the limit.

The neutral CP-even scalar sector mass matrix in eq. (13) is given by [†]

$$\boldsymbol{M}_{S^0}^2 = \tag{15}$$

$$\begin{bmatrix} B\mu\frac{v_2}{v_1} + \frac{1}{4}g_Z^2 v_1^2 + \mu\epsilon_3\frac{v_3}{v_1} + \frac{t_1}{v_1} & -B\mu - \frac{1}{4}g_Z^2 v_1 v_2 & -\mu\epsilon_3 + \frac{1}{4}g_Z^2 v_1 v_3 \\ -B\mu - \frac{1}{4}g_Z^2 v_1 v_2 & B\mu\frac{v_1}{v_2} + \frac{1}{4}g_Z^2 v_2^2 - B_2\epsilon_3\frac{v_3}{v_2} + \frac{t_2}{v_2} & B_2\epsilon_3 - \frac{1}{4}g_Z^2 v_2 v_3 \\ -\mu\epsilon_3 + \frac{1}{4}g_Z^2 v_1 v_3 & B_2\epsilon_3 - \frac{1}{4}g_Z^2 v_2 v_3 & \mu\epsilon_3\frac{v_1}{v_3} - B_2\epsilon_3\frac{v_2}{v_3} + \frac{1}{4}g_Z^2 v_3^2 + \frac{t_3}{v_3} \end{bmatrix}$$

where we have defined $g_Z^2 \equiv g^2 + g'^2$. In the upper–left 2×2 block, in the limit $v_3 = \epsilon_3 = 0$, the reader can recognize the MSSM mass matrix corresponding to the CP–even neutral

[†]Note that in the non-diagonal entries of this matrix the terms involving gauge couplings correct from those given in ref. [11] by a factor 2.

Higgs sector. Similar to the previous case, in this limit the third component decouples from the other two and corresponds to the real part of the sneutrino, which become degenerate with the imaginary part of the sneutrino. Another way of looking at the separation of the sneutrino field into real and imaginary parts is through a sneutrino–anti-sneutrino 45 degrees mixing [27].

In the general case, there will be a mixing between the Higgs sector and the stau sector. The three mass matrices in eqs. (9), (14), and (15) are diagonalized by rotation matrices which define the eigenvectors

$$\begin{bmatrix} G^+ \\ H^+ \\ \tilde{\tau}_1^+ \\ \tilde{\tau}_2^+ \end{bmatrix} = \boldsymbol{R}_{S^\pm} \begin{bmatrix} H_1^+ \\ H_2^+ \\ \tilde{\tau}_L^+ \\ \tilde{\tau}_R^+ \end{bmatrix}, \qquad \begin{bmatrix} G^0 \\ A \\ \tilde{\nu}_\tau^I \end{bmatrix} = \boldsymbol{R}_{P^0} \begin{bmatrix} \varphi_1^0 \\ \varphi_2^0 \\ \tilde{\nu}_\tau^I \end{bmatrix}, \qquad \begin{bmatrix} h \\ H \\ \tilde{\nu}_\tau^R \end{bmatrix} = \boldsymbol{R}_{S^0} \begin{bmatrix} \chi_1^0 \\ \chi_2^0 \\ \tilde{\nu}_\tau^R \end{bmatrix}, \qquad (16)$$

and the eigenvalues $\mathrm{diag}(0, \mathrm{m}^2_{\mathrm{H}^\pm}, \mathrm{m}^2_{\tilde{\tau}_1^\pm}, \mathrm{m}^2_{\tilde{\tau}_2^\pm}) = \boldsymbol{R}_{\mathrm{S}^\pm} \boldsymbol{M}^2_{\mathrm{S}^\pm} \boldsymbol{R}^{\mathrm{T}}_{\mathrm{S}^\pm}$ for the charged scalar sector, $\mathrm{diag}(0, \mathrm{m}^2_{\mathrm{A}}, \mathrm{m}^2_{\tilde{\nu}_\tau^{\mathrm{I}}}) = \boldsymbol{R}_{\mathrm{P}^0} \boldsymbol{M}^2_{\mathrm{P}^0} \boldsymbol{R}^{\mathrm{T}}_{\mathrm{P}^0}$ for the CP–odd neutral scalar sector, and $\mathrm{diag}(\mathrm{m}^2_{\mathrm{h}}, \mathrm{m}^2_{\mathrm{H}}, \mathrm{m}^2_{\tilde{\nu}_\tau^{\mathrm{R}}}) = \boldsymbol{R}_{\mathrm{S}^0} \boldsymbol{M}^2_{\mathrm{S}^0} \boldsymbol{R}^{\mathrm{T}}_{\mathrm{S}^0}$ for the CP–even neutral scalar sector.

The labelling for the different eigenstates is as follows. In the charged scalar sector the eigenstate with zero mass is denoted $G^\pm$. Among the other three, we call staus the two eigenstates $S_i^\pm$ with the biggest stau component calculated with $(R^{i3}_{S^\pm})^2 + (R^{i4}_{S^\pm})^2$, and by convention $m_{\tilde{\tau}_1^\pm} < m_{\tilde{\tau}_2^\pm}$. The remaining eigenstate is called charged Higgs $H^\pm$. In the neutral CP–odd sector, G^0 is the eigenstate P_i^0 with zero mass. Among the other two eigenstates, the one with largest stau component $(R^{i3}_{P^0})^2$ is called $\tilde{\nu}_\tau^I$, and the remaining state is the CP–odd Higgs A. Similarly, in the neutral CP–even sector, the state S_i^0 with the largest stau component $(R^{i3}_{S^0})^2$ is called sneutrino $\tilde{\nu}_\tau^R$. The other two are the neutral Higgs bosons h and H, with $m_h < m_H$. With this notation $H^\pm$ is the field which is mostly the MSSM charged Higgs, but with a small component of stau, and similarly for the neutral Higgs bosons.

If a 3×3 matrix $\boldsymbol{M}$ has a zero eigenvalue, then the other two eigenvalues satisfy

$$m_\pm = \frac{1}{2}\mathrm{Tr}\boldsymbol{M} \pm \frac{1}{2}\sqrt{(\mathrm{Tr}\boldsymbol{M})^2 - 4(M_{11}M_{22} - M_{12}^2 + M_{11}M_{33} - M_{13}^2 + M_{22}M_{33} - M_{23}^2)} \quad (17)$$

The CP-odd neutral scalar mass matrix eq. (14) has a zero determinant, so that its eigenvalues $m^2_{\tilde{\nu}_\tau^R}$ and m^2_A can be calculated exactly with the previous formula. The same can be done with the charged scalar mass matrix in the limit $h_\tau = 0$. In this case, the right stau decouples, and the eigenvalues $m^2_{H^\pm}$ and $m^2_{\tilde{\tau}_L^\pm}$ can be also calculated with eq. (17). Note that this limit is taken here only for the sake of illustration. In all numerical calculations we have used the realistic value for h_τ which is fixed through an exact tree-level relation eq. (28) given in the Appendix.

One can determine in the tree-level approximation the minimum of the scalar potential by imposing the condition of vanishing tadpoles in eq. (7). One–loop corrections change these equations to

$$t_i = t_i^0 - \delta t_i + T_i(Q) \tag{18}$$

where t_i, with $i = 1, 2, 3$, are the renormalized tadpoles, t_i^0 are the tree level tadpoles given in eq. (7), δt_i are the tadpole counter-terms, and $T_i(Q)$ are the sum of all one–loop contributions to the corresponding one–point functions with zero external momentum. The contribution from quarks and squarks to these tadpoles in our model can be found in ref. [10]. In an on shell scheme we identify the tree level tadpoles with the renormalized ones. Therefore, to find the correct minima we use eq. (7) unchanged, where now all the parameters are understood to be renormalized quantities.

We have used the diagrammatic (tadpole) method. Although equivalent to the effective potential method for minimization purposes [28], the diagrammatic method is better than the effective potential when it comes to calculating the one–loop corrected scalar masses [29]. Following ref. [30] (see also ref. [31]), we work in an on–shell scheme where by definition the tree level CP-odd Higgs mass m_A and the tree level W–boson mass m_W correspond to the respective pole masses. In this case, the renormalized charged Higgs mass $m_{H^\pm}$ is equal to the tree level charged Higgs mass calculated in the previous section, plus the following one–loop contributions:

$$\Delta m^2_{H^\pm} = \mathrm{Re}\left[A_{H^+H^-}(m_A^2 + m_W^2) - A_{AA}(m_A^2) - A_{WW}(m_W^2)\right] \tag{19}$$

where $A_{SS}(p^2)$, with $S = H^\pm, A, W$, are self energies. Each self energy is infinite by itself, but $\Delta m^2_{H^\pm}$ is finite.

Before turning to the numerical study of the charged Higgs mass spectrum, let us mention that, throughout this paper, except the case where the parameters have been fixed (Fig. 5 and 6), we have taken the MSSM parameters varying in the range :

$$\begin{array}{rcl} 0.5 < & \tan\beta & < 90 \\ 0\,\mathrm{GeV} < & M, M' & < 1000\,\mathrm{GeV} \\ 0\,\mathrm{GeV} < & m_{R_3}, m_{L_3} & < 300\,\mathrm{GeV} \\ -500\,\mathrm{GeV} < & A_\tau & < 500\,\mathrm{GeV} \\ -200\,\mathrm{GeV} < & \mu, B & < 0\,\mathrm{GeV} \end{array} \tag{20}$$

and the two R–Parity violating parameters varying as:

$$\begin{array}{rcl} -\,200\,\mathrm{GeV} < & \epsilon_3 & < 200\,\mathrm{GeV} \\ -90\,\mathrm{GeV} < & v_3 & < 90\,\mathrm{GeV} \end{array} \tag{21}$$

Note that $m^2_{H_1}$, $m^2_{H_2}$, and B_2 are fixed through the tadpole equations given in eq. (7). No big differences are observed if we take $\mu \geq 0$, and the sign of B is equal to the sign of μ because at the weak scale we have $m_A^2 \propto \mu B$. We are interested in relatively light charged Higgs , so we take $|\mu|$, $|B| \leq 200$ GeV . Similarly we are interested in relatively light staus , and that is why we take m_{R_3} , $m_{L_3} \leq 300$ GeV.

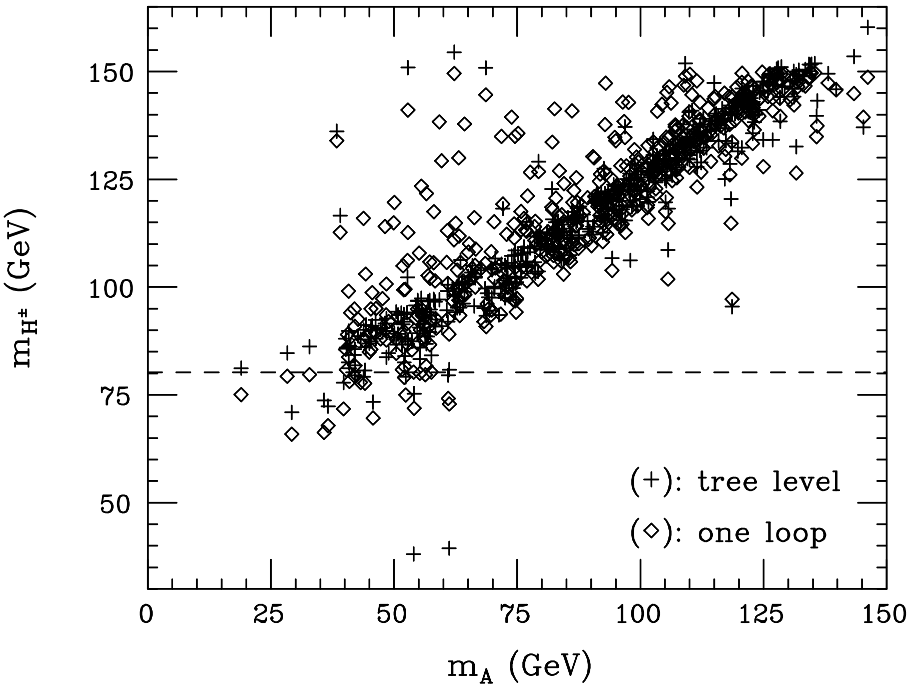

Figure 2: Tree level and one–loop charged Higgs boson mass as a function of the CP–odd Higgs mass m_A. The variation of parameters in the scan is indicated in the text. The horizontal dashed line corresponds to the W-boson mass.

We now turn to the numerical study of the lowest-lying charged scalar boson mass. Our results are illustrated in Fig. 2. In Fig. 2 we display allowed values of charged Higgs boson mass as a function of the CP-odd neutral Higgs boson mass m_A. Here the main point to note is that $m_{H^\pm}$ can be lower than expected in the MSSM, even before including radiative corrections. This is due to negative contributions arising from the R-parity violating stau-Higgs mixing, controlled by the parameter ϵ_3. We have varied the relevant model parameters in the range given by eq. (20) and eq. (21).

The one loop correction $\Delta m^2_{H^\pm}$ in eq. (19) depends on the soft squark masses M_Q, M_U and M_D, and the tri-linear soft masses A_t and A_b . Since they appear only through radiative corrections , for simplicity we have taken them degenerate at the weak scale : $M_Q = M_U = M_D = 1$ TeV and $A_t = A_b = A_\tau$.

An alternative way to display the influence of ϵ_3 parameter on the charged Higgs boson mass can be seen in Fig. 3. In this figure the curves corresponding to different ϵ_3 and v_3 values delimit the minimum theoretically allowed charged Higgs boson mass corresponding to those specific values. These curves are found in a scan where the MSSM parameters are varied according to eq. (20) and the R–parity violating parameters ϵ_3 and v_3 are varied according to the label in Fig. 3. Below each curve no points are found.

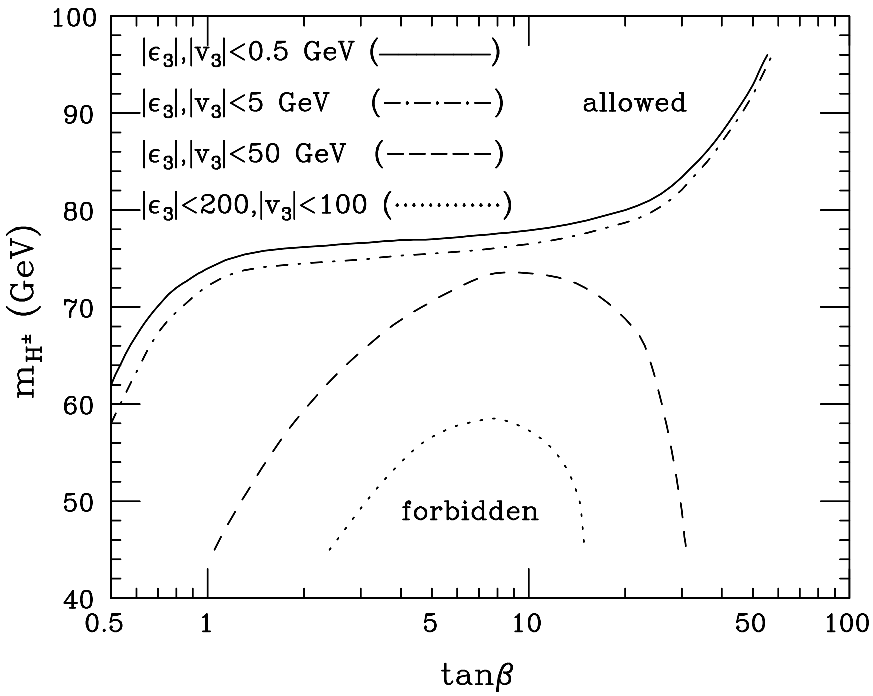

Figure 3: Minimum charged Higgs boson mass versus $\tan\beta$. Each curve corresponds to a different range of variation of the R–parity violating parameters ϵ_3 and v_3

The radiatively corrected MSSM prediction is recovered, as expected, when $\epsilon_3 = 0 = v_3$. This is indicated by the solid line in the figure. For larger values of the R-parity violating parameters one sees that the mass can be substantially lower than the MSSM expectation. This is due, again, to the negative contributions arising from the R-parity violating stau-Higgs mixing, as discussed above.

4 Production and Decays of Charged Scalars

Charged scalar pair production cross section can be calculated with the aid of the $ZS_i^+S_j^-$ Feynman rule, which is equal to $i\lambda^{ij}_{ZS^+S^-}(p+p')^\mu$ where p and p' are the momenta in the direction of the positive electric charge flow. The λ couplings are equal to $\boldsymbol{\lambda}_{ZS^+S^-} = \boldsymbol{R}_{S^\pm}\boldsymbol{\lambda}'_{ZS^+S^-}\boldsymbol{R}^T_{S^\pm}$ where the couplings in the unrotated basis are

$$\boldsymbol{\lambda}'_{ZS^+S^-} = \frac{g}{2c_W}\begin{bmatrix} -c_{2W} & 0 & 0 & 0 \\ 0 & -c_{2W} & 0 & 0 \\ 0 & 0 & -c_{2W} & 0 \\ 0 & 0 & 0 & 2s_W^2 \end{bmatrix} \tag{22}$$

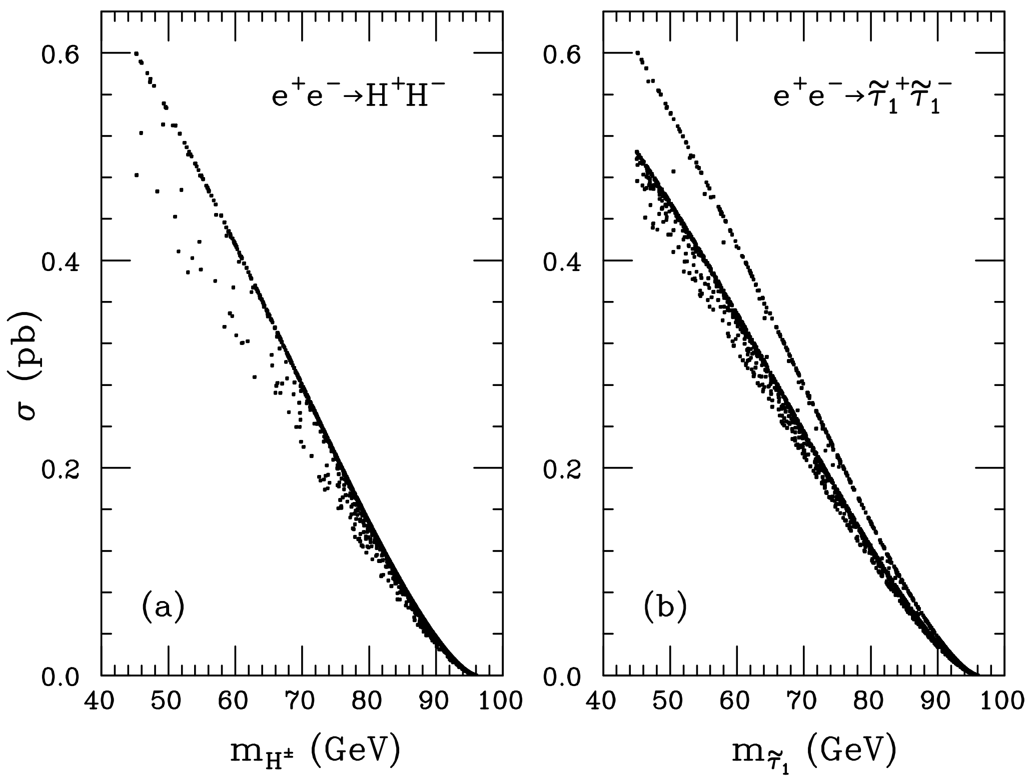

Figure 4: Total production cross section of a pair of (a) charged Higgs bosons and (b) light staus, as a function of their mass. The centre of mass energy is 192 GeV.

The differential cross section for the production of two charged scalars is

$$\frac{d\sigma}{d\cos\theta_{cm}}(e^+e^- \longrightarrow S_i^+ S_j^-) = \frac{1}{32\pi s}\lambda^{3/2}(1, m^2_{S_i^\pm}/s, m^2_{S_j^\pm}/s)\sin^2\theta_{cm} \times \left[\frac{e^4}{2}\delta_{ij} - \frac{ge^2}{2c_W}g_V^e \lambda^{ij}_{ZS^+S^-}\delta_{ij}\frac{s}{s-m_Z^2} + \frac{g^2}{8c_W^2}(g_V^{e2}+g_A^{e2})\lambda^{ij2}_{ZS^+S^-}\frac{s^2}{(s-m_Z^2)^2}\right] \tag{23}$$

where $g_V^e = \frac{1}{2} - 2s_W^2$ and $g_A^e = \frac{1}{2}$ and $\lambda(a,b,c) = a^2 + b^2 + c^2 - 2ab - 2ac - 2bc$.

In Fig. 4 we plot the production cross section of a pair of charged scalars as a function of its mass, considering $\sqrt{s} = 192$ GeV. In Fig. 4a we have $\sigma(e^+e^- \longrightarrow H^+H^-)$ and most of the points fall on the MSSM curve. The points that deviate from the main curve are due to mixing between charged Higgs and right stau. In fact, if the right stau were decoupled from the rest of the charged scalars, the charged Higgs pair production cross section would be identical to the MSSM for any value of ϵ_3 or v_3, and the reason is that the upper–left 3×3 relevant sub-matrix of $\boldsymbol{\lambda}'_{ZS^+S^-}$ in eq. (22) is proportional to the identity.

In Fig. 4b we have $\sigma(e^+e^- \longrightarrow \tilde\tau_1^+\tilde\tau_1^-)$ as a function of $m_{\tilde\tau_1^\pm}$. The points concentrate around the two MSSM curves corresponding to the cases where $\tilde\tau_1^\pm$ is mainly left stau (upper curve) and where $\tilde\tau_1^\pm$ is mainly right stau (lower curve). The smallness of the right-handed stau cross section relative to the left-handed one is understandable because

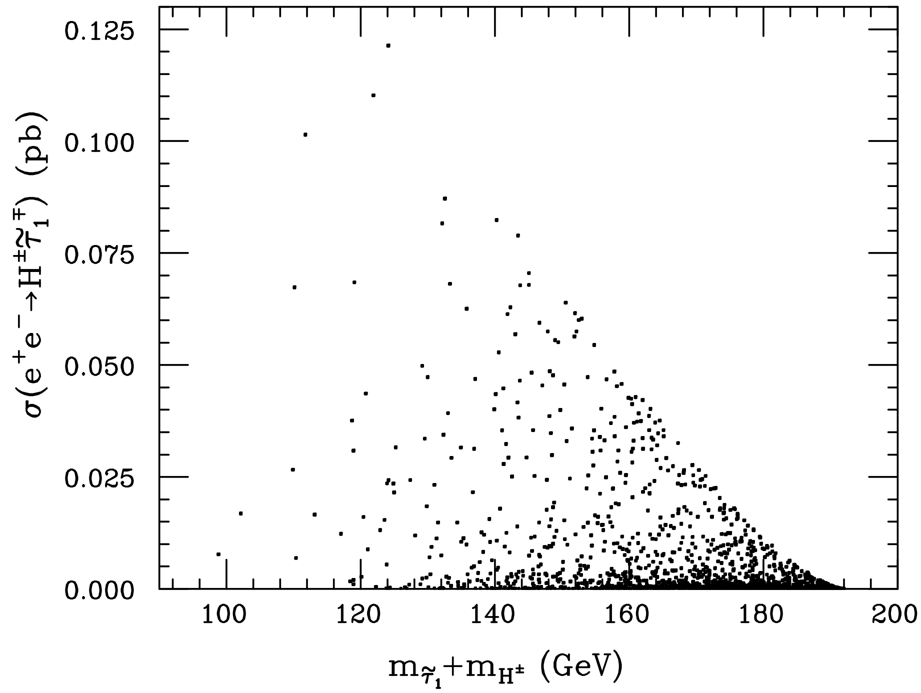

Figure 5: Mixed $H^{\pm}\tilde{\tau}_1^{\mp}$ production cross section as a function of $m_{H^{\pm}} + m_{\tilde{\tau}_1^{\mp}}$ at 192 GeV centre-of-mass energy.

the right-handed stau is an $SU(2)_L$ singlet. Again, the points which deviate from these curves are due to the mixing between the left and right staus.

An interesting characteristic of our model is the mixed production $e^+e^- \longrightarrow H^{\pm}\tilde{\tau}_1^{\mp}$ which is absent in the MSSM. In Fig. 5 we plot the total mixed production cross section, defined by $\sigma(e^+e^- \longrightarrow H^{\pm}\tilde{\tau}_1^{\mp}) \equiv \sigma(e^+e^- \longrightarrow H^+\tilde{\tau}_1^-) + \sigma(e^+e^- \longrightarrow H^-\tilde{\tau}_1^+)$, as a function of the sum of the final product masses $m_{H^{\pm}} + m_{\tilde{\tau}_1^{\pm}}$ for $\sqrt{s} = 192$ GeV. This mixed production cross section can be sizable, with a maximum value of the order of 0.12 pb.

As we have already seen, our model allows strong charged Higgs stau mixing, and this can substantially affect both the masses and the couplings. As a result the decay branching patterns of the charged scalar bosons can be significantly affected.

We now turn to a discussion of the charged scalar boson decays. Our first result here relates to the stau. In Fig. 6 we display the stau decay branching ratios below and past the neutralino threshold. In this figure we have fixed the parameters as

$$\begin{aligned} \mu &= 370 \ \text{GeV} & \tan\beta &= 2 \\ v_3 &= -4.7 \ \text{GeV} & M = 2M' &= 170 \ \text{GeV} \\ B &= 40 \ \text{GeV} & \epsilon_3 &= 10 \ \text{GeV} \\ A_\tau &= 500 \ \text{GeV} & m_{R_3} &= 400 \ \text{GeV} \end{aligned} \tag{24}$$

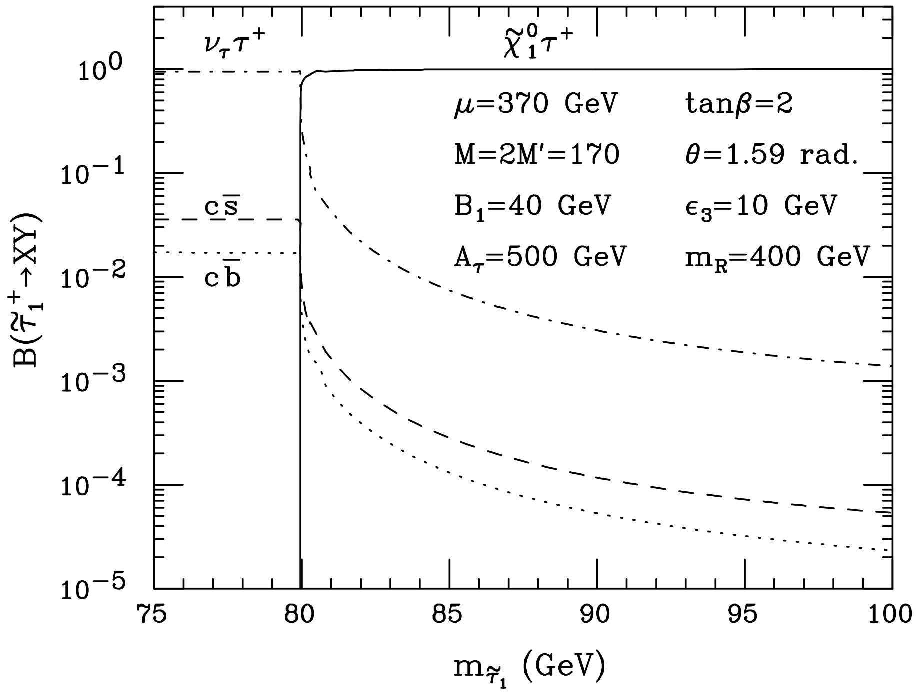

Figure 6: Stau branching ratios possible in our model for a particular choice of parameters. Note the neutralino threshold below which only R–parity violating decays are present.

in such a way as to ensure that the lightest neutralino is about 80 GeV in mass and thus may be produced as a decay product of a stau produced at LEP II energies.

In addition, we have chosen $\epsilon_3 = 10$ GeV and $v_3 = -4.7$ GeV ($\theta = 1.59$ rad) to demonstrate that we don't need large R–parity violating parameters to obtain sizeable effects. Otherwise, the choice of parameters is arbitrary .

Below the neutralino threshold, the stau is the LSP and will have only R-parity violating decays, therefore totally un-suppressed, even for the case of small R–parity breaking mixing. The main modes of stau decay in this case are into $\nu_\tau\tau$, $c\bar{s}$ and $c\bar{b}$, as clearly seen from Fig. 6. Moreover, one sees that for typical values of the R–parity breaking parameters, the stau will decay inside the detector.

Finally, for the case of the R-parity violating charged Higgs boson decays one can see from Fig. 7 that the branching ratios into supersymmetric channels can be comparable or even bigger than the R-parity conserving ones, even for relatively small values of ϵ and v_3.

Indeed it is explicitly seen from Fig. 7 that, in the region of small $\tan\beta$, the R-parity violating Higgs boson decay branching ratios can exceed the conventional ones and may reach values close to 100 %, since the R-parity–conserving decay $H^\pm \to \tau\nu_\tau$ is proportional

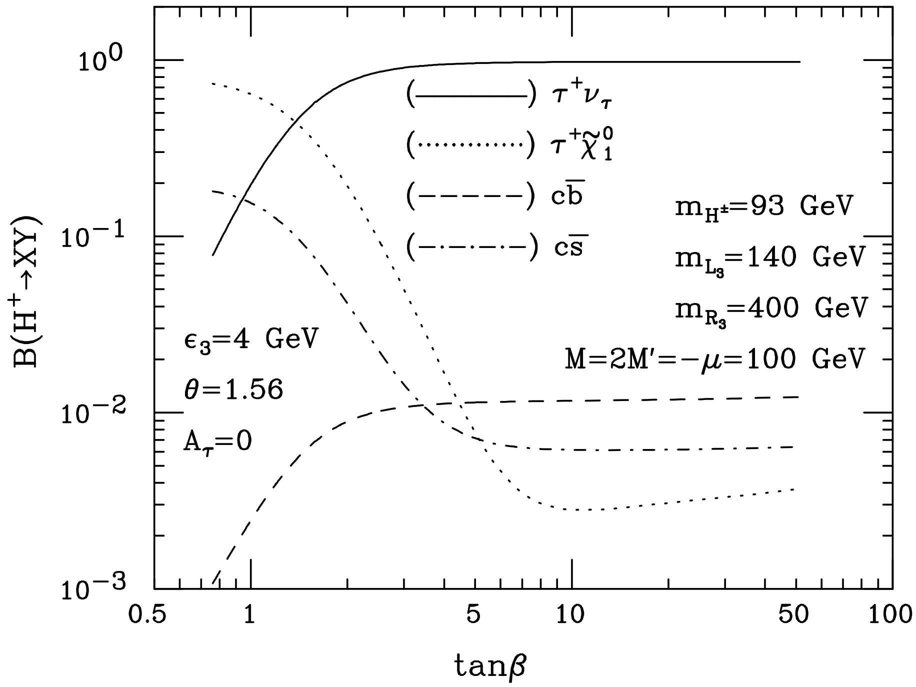

Figure 7: Charged Higgs branching ratios possible in our model for a particular choice of parameters. The R–parity violating decay dominates at low $\tan\beta$.

to $\tan^2\beta$ and so is usually dominant for larger $\tan\beta$. This figure was obtained for a fixed choice of parameters, given as

$$\begin{aligned} \mu &= -100 \text{ GeV} & m_{L_3} &= 140 \text{ GeV} \\ v_3 &= 2.66\text{GeV} & M = 2M' &= 100 \text{ GeV} \\ m_{H^\pm} &= 93 \text{ GeV} & \epsilon_3 &= 4 \text{ GeV} \\ A_\tau &= 0 \text{ GeV} & m_{R_3} &= 400 \text{ GeV} \end{aligned} \tag{25}$$

Note that in eq. (25) we have chosen μ with the opposite sign of eq. (24). This is done to see that the trends in Fig. 6 are not a peculiarity of that particular choice of parameters. Again we have taken small R–parity violating parameters: $\epsilon_3 = 4\,\text{GeV}$ and $v_3 = 2.66$ GeV. In order to get a ligth charged Higgs with $m_{H^\pm} = 93\,\text{GeV}$ we take $|\mu|$ and $|B|$ small.

We have also checked that a cosmologically safe 1 MeV ν_τ would not modify appreciably our conclusions. For example, we have verified that Fig. 7 remains unchanged if we vary m_{ν_τ} up to 1 MeV. The reason is that, although typically correlated with m_{ν_τ} the R–parity violating branching ratio may be large if the Higgs and stau masses are relatively close to each other, even for much smaller m_{ν_τ} values.

Another way to see that the dominance of R–parity-violating Higgs boson decays is not an accident of the above parameter choice is illustrated in Fig. 8. The various curves

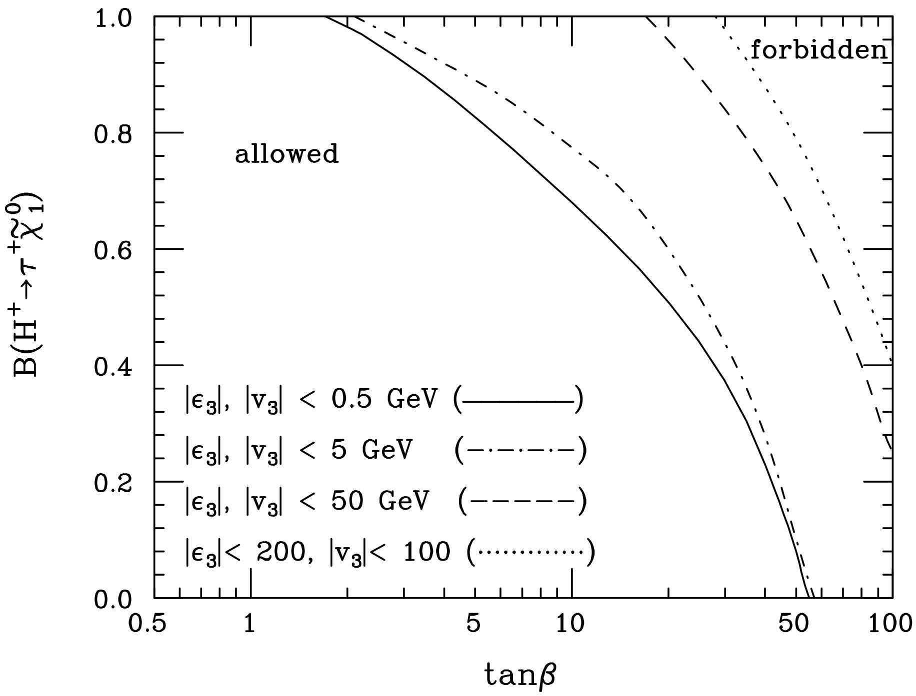

Figure 8: The curves denote the maximum attainable R-parity-violating charged Higgs branching ratio versus $\tan\beta$.

denote the maximum attainable values for the R–parity-violating Higgs boson branching ratio $B(H^+ \longrightarrow \tau^+\tilde{\chi}_1^0)$. The parameter space is scanned and the curves represent the boundaries above which no points are found. The MSSM parameters are varied according to eq. (20) and the range of variation of the R–parity violating parameters ϵ_3 and v_3 is indicated in the figure. In this way and in absolute generality, we demonstrate that even for very small R–parity violating parameters the branching ratio $B(H^+ \longrightarrow \tau^+\tilde{\chi}_1^0)$ can be close to unity, and that in the region of $\tan\beta \gg 1$ the decay $H^+ \longrightarrow \tau^+\nu_\tau$ dominates.

5 Discussion

We have seen in the last section [see Fig. 6] that if the lightest stau $\tilde{\tau}_1^\pm$ is the LSP it will decay only through R-parity-violating interactions, to cs or $\tau\nu_\tau$. As a result it leads to decay signatures which are identical to those of the charged Higgs boson in the MSSM. However, if it is not the LSP the $\tilde{\tau}_1^\pm$ is more likely to have standard R-parity-conserving decays such as neutralino plus τ, leading to signals that can be drastically different from those expected in the MSSM and which would arise from $\tilde{\chi}^0 \longrightarrow \nu_\tau Z^*$ or $\tilde{\chi}^0 \longrightarrow \tau W^*$. Unless ϵ_3 and v_3 are extremely small, the neutralino will decay inside the detector. For the case of stau pair production in e^+e^- colliders, such as LEP II, this would imply a plethora of new high fermion-multiplicity events (multi-jets and/or multi-leptons) . For example, di–tau + 4 jets + missing energy if both neutralinos decay into jets through neutral currents, or 4 taus + 4 jets if both neutralinos decay into jets through charged currents. Such processes are expected to have high rates and negligible background. As for hadron colliders, we also can have very high leptonic multiplicity events such as six leptons of which at least two are taus, plus missing momentum. This should be easy to see at the LHC, due again to the negligible standard model background. For a recent discussion see, for example, ref. [32].

If the charged Higgs is the lightest charged scalar boson one can distinguish two scenarios. If the R-parity-violating parameter ϵ_3 is small, the charged Higgs boson mainly decays to standard model fermions, thus conserving R-parity. However (as pointed out in ref. [11] for the case of neutral Higgs) it is possible to obtain very large branching ratios for R-parity-violating charged Higgs boson decays, even for moderate ϵ_3 values, as long as the mass difference between Higgs and staus is not too large. This happens because the R-parity-violating decay rates are governed by gauge strength interactions, whereas the R-parity-conserving ones are determined by Yukawa couplings.

In the opposite case of large $\epsilon_3 \sim m_W$ we always expect large branching ratios for R-parity-violating charged Higgs boson decay modes. As a result, one expects again a large number of novel signatures arising from neutralino decays. These would be the same large multiplicity events that we mentioned above for stau pair production.

Contrary to the MSSM, where stau and charged Higgs boson signal topologies are different, in our model they can be identical, for suitably chosen parameters. If one of these signals is observed, it will be difficult to know if it comes from a charged Higgs or a stau. To disentangle them, measurements of masses and decay rates of both particles would be required. Alternatively, less experimental information would be required to separate the origin of the signals if the particle mass spectrum were predicted from theory, as happens in supergravity versions of this model [10].

Another interesting feature of our model is the mixed production $e^+e^- \longrightarrow H^\pm\tilde{\tau}^\mp$. If $m_{\chi_1^0} < m_{\tilde{\tau}_1^\pm}$ then one can produce interesting signatures like di-tau + di-jets + missing energy. This is obtained when $H^\pm \longrightarrow \tau^\pm\nu_\tau$ and $\tilde{\tau}^\mp \longrightarrow \tau^\mp\tilde{\chi}_1^0 \longrightarrow \tau^\mp q\bar{q}\nu_\tau$. Although the cross section for this case is typically smaller, one can see from Fig. 5 that for many choices of parameters it may be non-negligible.

6 Conclusion

In summary we have considered the most salient aspects of the phenomenology of the charged scalar boson sector in the simplest effective low-energy R-parity breaking model characterized by a bilinear violation of R-parity in the superpotential. We have shown that the mass of the charged Higgs boson can be lower than expected in the MSSM, even before including radiative corrections. We have also studied the charged scalar boson decay branching ratios and show that the R-parity violating decay rates can be comparable or even bigger than the R-parity conserving ones. Moreover, if the stau is the LSP it will have only decays into standard model fermions, therefore totally un-suppressed. In the opposite case where it is heavier than the lightest neutralino one expects a plethora of exotic high fermion multiplicity events for which the standard model backgrounds should be rather small or absent. A detailed analysis of the detectability prospects of the related signatures at future accelerators lies outside the scope of the present paper and will be taken up elsewhere.

Acknowledgements

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Appendix

For completeness, in this Appendix we collect the Feynman rules relevant for the study of charged scalar decays into two fermions, where these two fermions are a chargino-tau and a neutralino-neutrino. We also give the exact formula, eq. (28), for the tau-lepton Yukawa coupling h_τ which differs in this model from that of the MSSM. First we set our conventions in the fermionic sector.

As we have already learned, the charged Higgs states H_1^2 and H_2^1 mix with the stau states $\tilde{\tau}_L^+$ and $\tilde{\tau}_R^+$ and form a set of four charged scalar eigenstates $S_i^\pm$ with $i = 1, 2, 3, 4$. The charged scalar mass matrix $\boldsymbol{M^2_{S^\pm}}$ in eq. (9) is diagonalized by a 4×4 rotation matrix $\boldsymbol{R}_{S^\pm}$ defined in eq. (16).

In a similar way, charginos mix with the tau lepton forming a set of three charge fermions $F_i^\pm$, $i = 1, 2, 3$. In a basis where $\psi^{+T} = (-i\lambda^+, \widetilde{H}_2^1, \tau_R^+)$ and $\psi^{-T} = (-i\lambda^-, \widetilde{H}_1^2, \tau_L^-)$, the charged fermion mass terms in the lagrangian are

$$\mathcal{L}_m = -\frac{1}{2}(\psi^{+T}, \psi^{-T}) \begin{pmatrix} \mathbf{0} & \boldsymbol{M}_C^T \\ \boldsymbol{M}_C & \mathbf{0} \end{pmatrix} \begin{pmatrix} \psi^+ \\ \psi^- \end{pmatrix} + h.c. \tag{26}$$

where the chargino/tau mass matrix is given by

$$\boldsymbol{M}_C = \begin{bmatrix} M & \frac{1}{\sqrt{2}}gv_2 & 0 \\ \frac{1}{\sqrt{2}}gv_1 & \mu & -\frac{1}{\sqrt{2}}h_\tau v_3 \\ \frac{1}{\sqrt{2}}gv_3 & -\epsilon_3 & \frac{1}{\sqrt{2}}h_\tau v_1 \end{bmatrix} \tag{27}$$

and M is the $SU(2)$ gaugino soft mass. We note that 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 $\boldsymbol{U}$ and $\boldsymbol{V}$ The tau Yukawa coupling h_τ is chosen such that one of the eigenvalues is equal to the tau mass. This is calculated from the vacuum expectation values of the model through an exact tree level relation given by

$$h_\tau^2 = \frac{2M_\tau^2}{v_1^2}\left(\frac{f + g(\varepsilon_3, v_3)}{f - \frac{2}{v_1^2}h(\varepsilon_3, v_3)}\right) \tag{28}$$

where :

$$f = g^2\left(\frac{1}{2}M_\tau^2(v_1^2+v_2^2) + M\mu v_1 v_2 - \frac{1}{4}g^2 v_1^2 v_2^2\right) + (\mu^2 - M_\tau^2)(M_\tau^2 - M^2)$$

$$\begin{aligned} g(\varepsilon_3, v_3) = {} & \tfrac{1}{2}v_3^2 g^2\left(M_\tau^2 - \mu^2 - \tfrac{1}{2}g^2 v_2^2\right) - \varepsilon_3^2\left(M^2 - M_\tau^2 + \tfrac{1}{2}g^2 v_1^2\right) \\ & -\varepsilon_3 v_3 g^2 (M v_2 + \mu v_1) \end{aligned}$$

$$\begin{aligned} h(\varepsilon_3, v_3) = {} & \varepsilon_3 v_3\left(\mu v_1 (M_\tau^2 - M^2) + \tfrac{1}{2}g^2 M v_1^2 v_2\right) + \tfrac{1}{2}\varepsilon_3^2 v_3^2 \left(M^2 - M_\tau^2\right) \\ & +\tfrac{1}{2}v_3^2\left(M_\tau^2(M_\tau^2 - M^2) - g^2 M_\tau^2 (v_1^2 + \tfrac{1}{2}v_2^2) - g^2 M\mu v_1 v_2 + \tfrac{1}{2}g^4 v_1^2 v_2^2\right) \\ & +\tfrac{1}{2}\varepsilon_3 v_3^3 g^2 M v_2 - \tfrac{1}{4}v_3^4 g^2 (M_\tau^2 - \tfrac{1}{2}g^2 v_2^2) \end{aligned}$$

In our model, the tau neutrino aquires mass, and this is due to a mixing between the neutralino sector and the neutrino–tau, forming a set of five neutral fermion F_i^0, $i = 1, ...5$. In the basis $\psi^{0T} = (-i\lambda', -i\lambda^3, \widetilde{H}_1^1, \widetilde{H}_2^2, \nu_\tau)$ the neutral fermions mass terms in the lagrangian are given by

$$\mathcal{L}_m = -\frac{1}{2}(\psi^0)^T \boldsymbol{M}_N \psi^0 + h.c. \tag{29}$$

where the neutralino/neutrino mass matrix is

$$\boldsymbol{M}_N = \begin{bmatrix} M' & 0 & -\frac{1}{2}g'v_1 & \frac{1}{2}g'v_2 & -\frac{1}{2}g'v_3 \\ 0 & M & \frac{1}{2}gv_1 & -\frac{1}{2}gv_2 & \frac{1}{2}gv_3 \\ -\frac{1}{2}g'v_1 & \frac{1}{2}gv_1 & 0 & -\mu & 0 \\ \frac{1}{2}g'v_2 & -\frac{1}{2}gv_2 & -\mu & 0 & \epsilon_3 \\ -\frac{1}{2}g'v_3 & \frac{1}{2}gv_3 & 0 & \epsilon_3 & 0 \end{bmatrix} \tag{30}$$

and M' is the $U(1)$ gaugino soft mass. This neutralino/neutrino mass matrix is diagonalized by a 5×5 rotation matrix $\boldsymbol{N}$ such that

$$\boldsymbol{N}^* \boldsymbol{M}_N \boldsymbol{N}^{-1} = \mathrm{diag}(m_{\chi_1^0}, m_{\chi_2^0}, m_{\chi_3^0}, m_{\chi_4^0}, m_{\nu_\tau}) \tag{31}$$

where by definition the eigenstate F_5^0 is the neutrino–tau, i.e., with the largest tau component $(N_{i5})^2$.

Now we are ready to work out the $S_i^{\pm} F_j^{\mp} F_k^0$ Feynman rules. These vertices are denoted by

$$S_i^+ F_j^- F_k^0 \longrightarrow i\lambda^{L\,ijk}_{S^+F^-F^0} P_L + i\lambda^{R\,ijk}_{S^+F^-F^0} P_R \tag{32}$$

where $P_L = \frac{1}{2}(1-\gamma_5)$ and $P_R = \frac{1}{2}(1+\gamma_5)$ are the usual left and right proyection operators. For simplicity, we work in a basis where the chaged scalars are unrotated. In this basis $S'^+ = (H_1^{2*}, H_2^1, \tilde{\tau}_L^+, \tilde{\tau}_R^+)$ the vertices are denoted by

$$S'^+_i F_j^- F_k^0 \longrightarrow i\lambda'^{L\,ijk}_{S^+F^-F^0} P_L + i\lambda'^{R\,ijk}_{S^+F^-F^0} P_R \tag{33}$$

The relation between λ and λ' is given by

$$\lambda^{L\,ijk}_{S^+F^-F^0} = \mathbf{R}^{il}_{S^\pm} \lambda'^{L\,ljk}_{S^+F^-F^0}\,, \qquad \lambda^{R\,ijk}_{S^+F^-F^0} = \mathbf{R}^{il}_{S^\pm} \lambda'^{R\,ljk}_{S^+F^-F^0} \tag{34}$$

and each of the λ' can be read from the following Feynman rules

$$\begin{aligned}
H_1^{2*} F_i^- F_j^0 &\longrightarrow i\frac{g}{2}\left[-U^*_{i1}N^*_{j3} + \frac{1}{\sqrt{2}}U^*_{i2}\left(N^*_{j2} + \frac{g'}{g}N^*_{j1}\right)\right](1-\gamma_5) + i\frac{h_\tau}{2}N_{j5}V_{i3}(1+\gamma_5) \\
H_2^1 F_i^- F_j^0 &\longrightarrow -i\frac{g}{2}\left[V_{i1}N_{j4} + \frac{1}{\sqrt{2}}V_{i2}\left(N_{j2} + \frac{g'}{g}N_{j1}\right)\right](1+\gamma_5) \\
\tilde{\tau}_L^+ F_i^- F_j^0 &\longrightarrow -i\frac{g}{2}\left[U^*_{i1}N^*_{j5} - \frac{1}{\sqrt{2}}U^*_{i3}\left(N^*_{j2} + \frac{g'}{g}N^*_{j1}\right)\right](1-\gamma_5) - i\frac{h_\tau}{2}N_{j3}V_{i3}(1+\gamma_5) \\
\tilde{\tau}_R^+ F_i^- F_j^0 &\longrightarrow -i\frac{g'}{\sqrt{2}}V_{i3}N_{j1}(1+\gamma_5) + i\frac{h_\tau}{2}\left(N^*_{j5}U^*_{i2} - N^*_{j3}U^*_{i3}\right)(1-\gamma_5)
\end{aligned} \tag{35}$$

The reader can check that from eq. (35) we can recover the MSSM Feynman rules taking the appropriate limits.

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