We review some of the main features of Bilinear R–Parity Violation (BRpV), defined by a quadratic term in the superpotential which mixes lepton and Higgs superfields and is proportional to a mass parameter $\epsilon$. We show how large values of $\epsilon$ can induce a small neutrino mass without fine-tuning. We mention the effect on the mass of the lightest Higgs boson. Finally we report on the effect of BRpV on gauge and Yukawa unification, showing that bottom–tau unification can be achieved at any value of $\tan \beta$.

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The Standard Model (SM) works well in describing the phenomenology of the strong and electroweak interactions of the known particles. For this reason, the motivations for studying supersymmetric extensions of the SM are mostly theoretical. The only experimental indication that favors the Minimal Supersymmetric Standard Model (MSSM) \[1\] in comparison with the SM is the unification of gauge couplings at some high scale \(M_{\text{GUT}}\) \[2, 3\].

Supersymmetry \[4\] is the only known way of unifying non–trivially the space-time symmetries of the Poincarè group with some other internal symmetry. This symmetry relates bosons with fermions, and at the same time affects the notion of space-time itself by introducing anti-commuting coordinates which extends the Minkowsky space into a superspace. Superfields are functions of superspace coordinates and the MSSM is constructed with vector and scalar superfields. Vector superfields \(\hat{V}\) contains a spin–1 gauge boson \(v_\mu\) and a fermionic partner \(\lambda\) (for example the photon and the photino in supersymmetric electrodynamics). Scalar superfields \(\hat{\Phi}\) contains a scalar boson \(\phi\) and a fermionic partner \(\psi\) (for example Higgs bosons and higgsinos). The superpotential is a cubic polynomial function of superfields.

It is customary to assign to each component field an R–Parity defined by \(R_p = (-1)^{3B+L+2S}\), where \(B\) is the barion number, \(L\) is the lepton number and \(S\) is the spin. In this way, quarks, leptons and Higgs bosons are R–Parity even, and the supersymmetric particles are R–Parity odd. If R–Parity is conserved, then supersymmetric particles are produced in pairs in the laboratory. In addition, the lightest supersymmetric particle (LSP, the lightest neutralino) is stable.

On the contrary, if R–Parity is not conserved then supersymmetric particles can be single produced, and the LSP decays into standard quarks and leptons. Furthermore, the LSP needs not to be the lightest neutralino. Possible terms in the superpotential which violate R–Parity are

\[
W_{R_p} = \lambda'_{ijk} \hat{U}_i \hat{D}_j \hat{D}_k + \varepsilon_{ab} \left[ \lambda'_{ijk} \hat{L}_a \hat{Q}_b \hat{D}_k + \lambda'_{ijk} \hat{L}_a \hat{L}_b \hat{R}_k + \epsilon_i \hat{L}_a \hat{H}_2^b \right],
\]

(1)

Trilinear R–Parity Violation (TRpV) corresponds to the first three terms and, considering that each of the generation indices \(i, j, k\) run from 1 to 3, they involve a very large number of arbitrary parameters. The only practical way to study TRpV is to consider one or two \(\lambda'\)’s different from zero at a time.

The fourth term in eq. (1) corresponds to Bilinear R–Parity Violation (BRpV) \[3, 4\], and involves only three extra parameters, one \(\epsilon_i\) for each generation. The \(\epsilon_i\) terms also violate lepton number in the \(i\)th generation respectively. Models where R–Parity is spontaneously broken \[5\] through vacuum expectation values (vev) of right handed sneutrinos \(\langle \tilde{\nu}_R \rangle = v_R \neq 0\) generate BRpV (and not TRpV)\[6\]. The \(\epsilon_i\) parameters are then

\(^{*}\) Of course, this is true in the original basis. If we rotate the Higgs and Lepton superfields then TRpV terms are generated, as explained later.
equal to some Yukawa coupling times $v_R$. Motivated by spontaneously broken R–Parity, we introduce explicitly BRpV in the MSSM superpotential and review the most important features of this model.

For simplicity we take from now on $\epsilon_1 = \epsilon_2 = 0$, in this way, only tau–lepton number is violated. In this case, considering only the third generation, the MSSM–BRpV has the following superpotential

$$W = \varepsilon_{ab} \left[ h_t \tilde{Q}_3 \tilde{U}_3 \tilde{H}^b_2 + h_b \tilde{Q}_3 \tilde{D}_3 \tilde{H}^a_1 + h_\tau \tilde{L}_3 \tilde{R}_3 \tilde{H}^a_1 - \mu \tilde{H}^a_1 \tilde{H}^b_2 + \epsilon_3 \tilde{L}_3 \tilde{H}^b_2 \right], \quad (2)$$

where the first four terms correspond to the MSSM. The last term violates tau–lepton number as well as R–Parity.

The presence of the $\epsilon$ term in the superpotential implies that the tadpole equation for the tau sneutrino is non–trivial, i.e., the vacuum expectation value $\langle \tilde{\nu}_{\tau} \rangle = v_3/\sqrt{2}$ is non–zero. This in turn generates more R–parity and tau lepton number violating terms which, in particular, induce a tau neutrino mass as we will see later.

By looking at the last two terms in the superpotential an immediate question arises. Can the BRpV term be rotated away from the superpotential? and consequently, is the $\epsilon$ term physical? Indeed, consider the following rotation of the superfields \[8\]

$$\tilde{H}'_1 = \frac{\mu \tilde{H}_1 - \epsilon_3 \tilde{L}_3}{\sqrt{\mu^2 + \epsilon_3^2}}, \quad \tilde{L}'_3 = \frac{\epsilon_3 \tilde{H}_1 + \mu \tilde{L}_3}{\sqrt{\mu^2 + \epsilon_3^2}}. \quad (3)$$

In the new basis the $\epsilon$ term disappears from the superpotential, nevertheless, R–Parity is reintroduced in the form of TRpV. The superpotential in the new basis is

$$W = h_t \tilde{Q}_3 \tilde{U}_3 \tilde{H}_2' + h_b \frac{\mu}{\mu'} \tilde{Q}_3 \tilde{D}_3 \tilde{H}'_1 + h_\tau \tilde{L}'_3 \tilde{R}_3 \tilde{H}'_1 - \mu' \tilde{H}'_1 \tilde{H}_2' + h_b \frac{\epsilon_3}{\mu'} \tilde{Q}_3 \tilde{D}_3 \tilde{L}_3', \quad (4)$$

where $\mu'^2 = \mu^2 + \epsilon_3^2$. The first four terms are MSSM looking terms and the last term violates the R–Parity defined in the new basis. Note the re-scaling in the bottom quark Yukawa term. Its presence ensures that the same quark mass is obtained with the same Yukawa coupling in the two basis. This re-scaling is non-trivial and has important consequences in Yukawa unification, as shown later.

As we know, supersymmetry must is broken and this is parametrized by soft supersymmetry breaking terms. The soft terms which play an important role in BRpV are the following

$$V_{soft} = m^2_{H_1} |H_1|^2 + M^2_{L_3} |L_3|^2 - \left[ B \mu H_1 H_2 - B_2 \epsilon_3 L_3 H_2 + h.c. \right] + ... \quad (5)$$

where $m^2_{H_1}$ and $M^2_{L_3}$ are the soft masses corresponding to the fields $H_1$ and $L_3$ respectively, and $B$ and $B_2$ are the bilinear soft mass parameters associated to the next-to-last and last terms in the superpotential in eq. \[(2)\]. It is clear, for example, that Higgs vacuum expectation values $\langle H_i \rangle = v_i/\sqrt{2}$ induce a non-trivial tadpole equation and a non-zero vev for the sneutrino through the $B_2$ term in eq. \[(4)\].
The soft terms in the rotated basis are given by
\[ V_{sof} = \frac{m_H^2 + m_{\text{soft}}}{{\mu}^2} |H'|^2 + \frac{m^2_{\tilde{H}_3}}{\mu^2} |\tilde{L}'_3|^2 - \frac{B\mu^2 + 2\epsilon^2_3\mu'}{\mu'} H'H_2 \\
- \frac{\epsilon_3\mu}{\mu^2} (m^2_{\tilde{H}_1} - M^2_{\tilde{L}_2}) \tilde{L}'_3 H' - \frac{\epsilon_3\mu}{\mu'} (B_2 - B) \tilde{L}'_3 H_2 + h.c. \]  

(6)

The first three terms are MSSM like terms equivalent to the first three terms in eq. (6). In fact, in analogy with the MSSM, the coefficients of $|H'|^2$ and $|\tilde{L}'_3|^2$ could be defined in the rotated basis as the soft masses $m^2_{\tilde{H}_1}$ and $M^2_{\tilde{L}_3}$ respectively, and the coefficient of $H'H_2$ would be the new bilinear soft term $B'\mu'$. The last two terms violate R–Parity and tau lepton number, and are equivalent to the last term in eq. (6), i.e., they induce a non-zero vev for the tau sneutrino field in the rotated basis $⟨\tilde{\nu}_r⟩ = v_3/\sqrt{2}$.

Vacuum expectation values are calculated by minimizing the scalar potential, or equivalently, by imposing that the tadpoles are equal to zero. The linear terms of the scalar potential are $V_{\text{linear}} = t_1\chi^0_1 + t_2\chi^0_2 + t_3\tilde{\nu}^R$, where $\chi^0_i = \sqrt{2} Re(H_i^0) - v_i$ and $\tilde{\nu}^R = \sqrt{2} Re(\tilde{\nu}_r) - v_3$. The $t_i$ are the tree level tadpoles and they are equal to zero at the minimum. In the original basis the tadpole equations are

\[ t_1 = (m^2_{\tilde{H}_1} + \mu^2)v_1 - B\mu v_2 - \mu \epsilon_3 v_3 + \frac{1}{8}(g^2 + g'^2) v_1 (v_1^2 - v_2^2 + v_3^2) = 0, \]
\[ t_2 = (m^2_{\tilde{H}_2} + \mu^2 + \epsilon_3^2)v_2 - B\mu v_1 + B_2\epsilon_3 v_3 - \frac{1}{8}(g^2 + g'^2) v_2 (v_1^2 - v_2^2 + v_3^2) = 0, \]
\[ t_3 = (M^2_{\tilde{L}_3} + \epsilon_3^2)v_3 - \mu \epsilon_3 v_1 + B_2\epsilon_3 v_2 + \frac{1}{8}(g^2 + g'^2) v_3 (v_1^2 - v_2^2 + v_3^2) = 0. \]  

(7)

The first two tadpole equations reduce to the MSSM minimization conditions after taking the MSSM limit $\epsilon_3 = v_3 = 0$, and in this case, the third tadpole equation is satisfied trivially. Note that $\epsilon_3 = 0$ implies two solutions for $v_3$ from the third tadpole in eq. (6), from which only $v_3 = 0$ is viable because the second solution implies the existence of a massless pseudoscalar.

The first two tadpole equations in the rotated basis are

\[ t'_1 = \mu'^2 v'_1 + \frac{m^2_{\tilde{H}_1} + \mu^2}{\mu'^2} v'_1 - \frac{B\mu^2 + 2\epsilon_3\mu'}{\mu'} v_2 + \frac{m^2_{\tilde{H}_1} - M^2_{\tilde{L}_3}}{\mu'^2} \epsilon_3 \mu' v'_3 \\
+ \frac{1}{8}(g^2 + g'^2) (v'_1^2 - v_2^2 + v_3^2) = 0 \]  

(8)

\[ t'_2 = \mu'^2 v_2 + \frac{m^2_{\tilde{H}_2} + \mu^2}{\mu'^2} v_2 - \frac{B\mu^2 + 2\epsilon_3\mu'}{\mu'} v'_1 + (B_2 - B) \frac{\epsilon_3 \mu}{\mu'} v'_3 \\
- \frac{1}{8}(g^2 + g'^2) v_2 (v'_1^2 - v_2^2 + v_3^2) = 0 \]  

(9)

where $⟨H'_1⟩ = v'_1/\sqrt{2}$ and $⟨\tilde{L}'_3⟩ = v'_3/\sqrt{2}$, and the following relations hold $v'_1 = (\mu v_1 - \epsilon_3 v_3)/\mu'$ and $v'_3 = (\epsilon_3 v_1 + \mu v_3)/\mu'$, as suggested by eq. (6). These two tadpole equations resemble the MSSM minimization conditions when we set $v'_3 = 0$. The third tadpole equation is

\[ t'_3 = (m^2_{\tilde{H}_1} - M^2_{\tilde{L}_3}) \frac{\epsilon_3 \mu}{\mu'^2} v'_1 + (B_2 - B) \frac{\epsilon_3 \mu}{\mu'} v_2 + \frac{m^2_{\tilde{H}_1} \epsilon_3^2 + M^2_{\tilde{L}_3} \mu^2}{\mu'^2} v'_3 \\
+ \frac{1}{8}(g^2 + g'^2) v'_3 (v'_1^2 - v_2^2 + v_3^2) = 0 \]  

(10)
In this equation we observe that \( v'_3 = 0 \) if \( \Delta m^2 \equiv m^2_{H_1} - M^2_{L_{3}} = 0 \) and \( \Delta B \equiv B_2 - B = 0 \) at the weak scale, which is not true in general. In supergravity models with universality of scalar soft masses and bilinear mass parameters we have \( \Delta m^2 = 0 \) and \( \Delta B = 0 \) at the unification scale \( M_{\text{GUT}} \approx 10^{16} \) GeV, but radiative corrections spoil this degeneracy. In the approximation where \( \Delta m^2 \) and \( \Delta B \) are small we find that \( v'_3 \) is also small and in first approximation given by

\[
v'_3 \approx -\frac{\epsilon_3 \mu}{\mu'^2 m^2_{\tilde{\nu}}} \left( v'_1 \Delta m^2 + \mu' v_2 \Delta B \right)
\]

(11)

where we have introduced

\[
m^2_{\tilde{\nu}} \equiv \frac{m^2_{H_1} \epsilon_3^2 + M^2_{L_{3}} \mu'^2}{\mu'^2} + \frac{1}{8} (g'^2 + g^2) (v'^2_1 - v^2_2)
\]

which reduces to the tau sneutrino mass in the MSSM when we set \( \epsilon_3 = 0 \).

As a consequence of tau lepton number and BRpV terms, characterized by the parameters \( \epsilon_3 \) and \( v_3 \), a mixing between neutralinos and the tau neutrino is generated. This implies that the tau neutrino acquires a mass \( m_{\nu_\tau} \). In the original basis, where \( (\psi^0)^T = (-i\lambda', -i\lambda^3, \tilde{H}_1, \tilde{H}_2, \nu_\tau) \), the scalar potential contains the following mass terms

\[
\mathcal{L}_m = -\frac{1}{2} (\psi^0)^T M_N \psi^0 + \text{h.c.}
\]

(13)

where the neutralino/neutrino mass matrix is

\[
M_N = \begin{bmatrix}
M' & 0 & \frac{1}{2} g' v_1 & \frac{1}{2} g' v_2 & -\frac{1}{2} g' v_3 \\
0 & M & \frac{1}{2} g v_1 & -\frac{1}{2} g v_2 & \frac{1}{2} g v_3 \\
-\frac{1}{2} g' v_1 & \frac{1}{2} g v_1 & 0 & -\mu & 0 \\
\frac{1}{2} g' v_2 & -\frac{1}{2} g v_2 & -\mu & 0 & \epsilon_3 \\
-\frac{1}{2} g' v_3 & \frac{1}{2} g v_3 & 0 & \epsilon_3 & 0
\end{bmatrix}
\]

(14)

Here \( M \) and \( M' \) are the \( SU(2) \) and \( U(1) \) gaugino masses. It can be seen from eq. (14) that mixings between tau neutrino and neutralinos are proportional to \( \epsilon_3 \) and \( v_3 \). Naively one could think that, due to the strong experimental constraint on the tau neutrino mass, the parameters \( \epsilon_3 \) and \( v_3 \) should be small compared with \( m_Z \). This is not the case, and from Fig. 1 we observe that \( |\epsilon_3| \) can be as large as 400 GeV!

Indeed, to make Fig. 1 we have embedded the MSSM–BRpV model into supergravity, with universality of scalar \( (m_0) \), gaugino \( (M_{1/2}) \), bilinear \( (B) \), and trilinear \( (A) \) soft mass parameters at the unification scale \( M_X \approx 10^{16} \) GeV. We have imposed the radiative breaking of the electroweak symmetry by minimizing the scalar potential with the aid of one–loop tadpole equations. We have made a scan over the parameter space, including the BRpV parameters \( \epsilon_3 \) and \( v_3 \). Points that satisfy the constraint \( m_{\nu_\tau} < 30 \) MeV are kept (we also impose that the supersymmetric particles are not too light).

We observe from Fig. 1 that it is easy to satisfy the constraint on the tau neutrino mass, and even \( m_{\nu_\tau} \) of the order of 1 eV can be achieved. The central region where \( \epsilon_3 \) is
Figure 1: Tau neutrino mass as a function of the R–Parity violating parameter $\epsilon_3$.

Figure 2: Tau neutrino mass as a function of the tau sneutrino vacuum expectation value $v_3$ in the original basis.
close to zero is less populated at high values of \( m_{\nu_\tau} \) because in this case we are closer to the MSSM, where the neutrinos are massless.

Similarly, in Fig. 2 we plot \( m_{\nu_\tau} \) as a function of the vacuum expectation value of the sneutrino \( v_3 \). For the same reason we already mention, the central region at high values of \( m_{\nu_\tau} \) is less populated. In this figure we observe that the BRpV parameter \( v_3 \) is not necessarily small, and that \( |v_3| \) can be as high as 100 GeV. The value of \( |v_3| \) cannot be as high as \( \epsilon_3 \) because the sneutrino vev contributes also to the W–boson mass according to

\[
m^2_W = \frac{1}{4} (g^2 + g'^2)(v_1^2 + v_2^2 + v_3^2).
\]

Considering that the mass terms which mix the neutrino with the neutralinos are proportional to \( \epsilon_3 \) and \( v_3 \), an obvious question arises: how can we get a small neutrino mass? The answer lies in the fact that the induced neutrino mass satisfy \( m_{\nu_\tau} \sim (\epsilon_3 v_1 + \mu v_3)^2 \), and this last combination is what needs to be small. Indeed, as we will see below, in models with universality of scalar and bilinear soft mass parameters, the combination \( (\epsilon_3 v_1 + \mu v_3) \) is radiatively induced, and therefore, naturally small.

In Fig. 3 we have the dependence of the tau neutrino mass \( m_{\nu_\tau} \) as a function of the parameter \( \xi \equiv (\epsilon_3 v_1 + \mu v_3)^2 = (\mu' v'_3)^2 \). We see a clear correlation between \( m_{\nu_\tau} \) and \( v'_3 \). The parameter \( |v'_3| \) takes a maximum value of the order of 10 GeV.

The neutralino–neutrino mass matrix in the rotated basis, analogous to eq. (14), is

\[
M'_{\nu} = R(M_{\nu}),
\]

where the rotation \( R \) is defined by eq. (8) or, equivalently, by the substitution \((v_1, v_3, \epsilon_3, \mu) \rightarrow (v'_1, v'_3, 0, \mu')\). In this basis the \( \epsilon \) term is not present, and the only source of mixing responsible for the neutrino mass is the vev \( v'_3 \). In first approximation, valid when \( v'_3 \) is small, we get

\[
m_{\nu_\tau} \approx -\frac{(g^2 M + g'^2 M')(\mu^2 v'_3^2)}{4M'M'\mu^2 - 2(g^2 M + g'^2 M')v'_1 v'_2 \mu'}
\]
On the other hand, considering the renormalization group equations for the soft mass parameters \(m^2_{H^1}, m^2_{L_3}, B, \) and \(B_2\), which solved in first approximation give us

\[
m^2_{H^1} - M_{L_3}^2 \approx -\frac{3h^2}{8\pi^2} \left( m^2_{H^1} + M_Q^2 + M_D^2 + A_D^2 \right) \ln \frac{M_{GUT}}{m_Z}
\]
\[
B_2 - B \approx \frac{3h^2}{8\pi^2} A_D \ln \frac{M_{GUT}}{m_Z}
\]

we can show, using eq. (11), that the tau neutrino mass is radiatively generated and given by

\[
m_{\nu_\tau} \approx \frac{m_Z^2}{M_{SUSY}} \left( \frac{\epsilon_3}{M_{SUSY}} \right)^2 h_b^4 \sim 1 \text{ KeV}
\]

where the 1 KeV was obtained in the case \(M_{SUSY} \sim \epsilon_3 \sim m_Z\) and \(h_b \sim 10^{-2}\). An even lighter \(\nu_\tau\) can be obtained if we increase \(M_{SUSY}\) or decrease \(\epsilon_3\), as can be seen from Fig. 3, where neutrinos as light as \(m_{\nu_\tau} \sim 1 \text{ eV}\) are shown [10].

Another interesting feature of BRpV is that the neutral CP–even Higgs sector now mixes with the real part of the tau sneutrino forming a set of three neutral CP–even scalars \(S_i^0\), \(i = 1, 2, 3\). In the original basis, where \(S^0 = [\chi^0_1, \chi^0_2, \tilde{\nu}_\tau^0]\), the mass matrix is given by

\[
M^2_{S^0} = 
\begin{bmatrix}
B\mu_v + \frac{1}{4}g_Z^2 v_1^2 + \mu_\epsilon_3 v_1 & -B\mu + \frac{1}{4}g_Z^2 v_1 v_2 & -\epsilon_3 + \frac{1}{4}g_Z^2 v_1 v_3 \\
-B\mu - \frac{1}{4}g_Z^2 v_1 v_2 & B\mu_v + \frac{1}{4}g_Z^2 v_2^2 - B_2\epsilon_3 v_2 & B_2\epsilon_3 - \frac{1}{4}g_Z^2 v_2 v_3 \\
-\epsilon_3 + \frac{1}{4}g_Z^2 v_1 v_3 & B_2\epsilon_3 - \frac{1}{4}g_Z^2 v_2 v_3 & \epsilon_3 v_3 - B_2\epsilon_3 v_3 + \frac{1}{4}g_Z^2 v_3^2
\end{bmatrix}
\]

In the MSSM limit, where \(\epsilon_3 = v_3 = 0\), the mass matrix \(M^2_{S^0}\) reduced to a \(2 \times 2\) block corresponding to the normal CP–even Higgs sector of the MSSM, and a decoupled tau sneutrino. A similar effect occurs with the charged Higgs sector, which couples to the stau sector forming a set of four charged scalars, one of them being the unphysical Goldstone boson [11].

We have calculated the lightest CP–even Higgs mass in BRpV and compared it with its mass in the MSSM and the result is plotted in Fig. 4. We include only the largest radiative corrections proportional to \(m_t^4\) [12]. We observe that the lightest Higgs mass is in general decreased due to the mixing with the sneutrino and, of course, the effect disappear as the BRpV parameter \(|v_3|\) approaches to zero. In this case the Higgs \(h\) have R–Parity violating decays because it can “behaves” as a tau sneutrino.
Similarly to the Higgs bosons, charginos mix with the tau lepton forming a set of three charged fermions $F^+_i, i = 1, 2, 3$. In the original basis where $\psi^+ = (−i\lambda^+, \tilde{H}_2^1, \tau^+_R)$ and $ψ^− = (−i\lambda^−, \tilde{H}_1^2, \tau^−_L)$, the charged fermion mass terms in the lagrangian are $L_m = −ψ^− M_C ψ^+$, with the mass matrix given by

$$
M_C = \begin{bmatrix}
M & \frac{1}{\sqrt{2}} g v_2 & 0 \\
\frac{1}{\sqrt{2}} g v_1 & \mu & -\frac{1}{\sqrt{2}} h_\tau v_3 \\
\frac{1}{\sqrt{2}} g v_3 & -\epsilon_3 & \frac{1}{\sqrt{2}} h_\tau v_1
\end{bmatrix}
$$

As a result, the tau Yukawa coupling is not related to the tau mass by the usual MSSM relation. On the contrary, $h_\tau$ depends now on the parameters of the chargino sector $M, \mu, \tan \beta$, and $h_\tau$ as well as the BRpV parameters $\epsilon_3$ and $v_3$, through a formula given in ref. [11].

In addition, the top and bottom quark Yukawa couplings are related to the quark masses by

$$
m_t = h_t \frac{v}{\sqrt{2}} \sin \beta \sin \theta, \quad m_b = h_b \frac{v}{\sqrt{2}} \cos \beta \sin \theta
$$

where $v = 246$ GeV and we have defined $\cos \theta \equiv v_3/v$.

These differences with the MSSM have profound consequences on Yukawa unification as shown in Fig. 5. In this figure we observe that bottom–tau Yukawa unification can be achieved at any value of $\tan \beta$ by choosing appropriately the value of $v_3$ [13]. The plot in Fig. 5 is made with a scan over parameter space such that points which satisfy $h_b(M_{GUT}) = h_\tau(M_{GUT})$ within 1% are kept, where $M_{GUT}$ is the gauge coupling unification scale. Each selected point is placed in one of the regions of Fig. 5 according to its $|v_3|$ value. The diagonal band at high values of $\tan \beta$ corresponds to points where top-bottom-tau unification is achieved [13].

In summary, it is shown that BRpV is the simplest extension of the MSSM which
introduce R–Parity violation. This model can be successfully embedded into Supergravity models with universality of scalar, gaugino, bilinear and trilinear soft mass parameters. In this case, the induced tau neutrino mass is radiatively generated and, therefore, naturally small. It is shown that the BRpV parameters $\epsilon_3$ and $v_3$ do not need to be small, in fact they can be easily of the order of $m_Z$. In addition, in this model the CP–even Higgs bosons couple with the tau sneutrino field, and the effect on the mass of the lightest Higgs is to lower it compared to the MSSM. Finally, BRpV changes the relation between the Yukawa couplings and the masses of the top and bottom quarks and the tau lepton. As a consequence, bottom-tau Yukawa unification can be achieved at any value of the parameter $\tan \beta$ provided we choose appropriately the value of the sneutrino vev $v_3$. Top-bottom-tau unification is achieved in a slightly wider region at high $\tan \beta$. We would like to stress the fact that, even in the unlikely limit where the tan neutrino is massless with $\epsilon_3 \neq 0$ (if $v_3' = 0$, obtained when there is universality of soft mass parameters at the weak scale, which is not natural) R–Parity is not conserved, and even though the neutralinos decouple from the tau neutrino, the lightest neutralino decays for example to $b\bar{b}\nu_\tau$ through an intermediate sbottom.

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