SUPER-GRAVITY UNIFICATION WITH BILINEAR R–PARITY VIOLATION

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Bilinear R–parity violation (BRpV) provides the simplest and most meaningful way to include such effects into the Minimal Supersymmetrical Standard Model (MSSM). It is defined by a quadratic superpotential term $\epsilon LH$ which mixes lepton and Higgs superfields and mimics the effects of models with spontaneous breaking. I review some of its main features and show how large $\epsilon$ values can lead to a small neutrino mass radiatively, without any fine-tuning. I discuss the effect of BRpV on gauge and Yukawa unification, showing how bottom–tau unification can be achieved at any value of $\tan \beta$. However, for very large $m_\epsilon$ values the large $\tan \beta$ solution is ruled out.

1 Introduction

Although the Standard Model (SM) works well in describing the phenomenology of the strong and electroweak interactions of the known particles it leaves unanswered some theoretical issues such as the hierarchy problem and the unification of the gauge couplings at $M_{GUT}$. These provide strong impetus to the study of supersymmetric extensions of the SM, the simplest being the Minimal Supersymmetric Standard Model (MSSM). In this case one can show that the unification of gauge couplings at $M_{GUT}$ occurs for acceptable of the effective Supersymmetry breaking as well as unification scales.

It is usual to assign to SM state an R–Parity defined by $R_p = (-1)^{3B+L+2S}$, where $B$ is the baryon number, $L$ is the lepton number and $S$ is the spin. In this way, quarks, leptons and Higgs bosons are R–even, and the supersymmetric particles are R–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.

In contrast, if R–Parity is violated then supersymmetric particles can be singly produced, and the LSP decays into standard quarks and leptons. Moreover, the LSP need not be the lightest neutralino.
2 R–Parity Violation

Possible terms in the superpotential which violate R–Parity are

\[ W_{R_p} = \lambda''_{ijk} \hat{U}^i \hat{D}^j \hat{D}^k + \varepsilon_{ab} \left( \lambda'_{ijk} \hat{L}^i \hat{Q}^j \hat{D}^k + \lambda_{ijk} \hat{L}^i \hat{L}^j \hat{R}^k + \epsilon_i \hat{L}^i \hat{H}^b \right), \]  

(1)

Such terms may arise as residues of unification, e.g. as gravitational effects. The first three terms are Trilinear R–Parity Violation (TRpV) terms. Each of the generation indices \( i, j, k \) runs from 1 to 3, thus implying a very large number of arbitrary parameters. It is impossible to provide a systematic way to analyze the implications of TRpV, the best one can do is to consider one or two \( \lambda \)'s different from zero at a time. Some of these couplings are strongly restricted by proton stability and/or primordial baryon asymmetry survival.

The fourth term in eq. (1) corresponds to Bilinear R–Parity Violation (BRpV) \(^5\), 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 broken spontaneously \(^7\) through vacuum expectation values (vev) of right handed sneutrinos \( \langle \tilde{\nu}^c \rangle = v_R \neq 0 \) generate BRpV (and not TRpV) \(^a\). Such spontaneous R-Parity Violation scenarios are also interesting from the point of view of the electroweak phase transition and baryogenesis \(^8\) as well as phenomenologically, due to the existence of massless pseudoscalar majorons \(^7\) which brings in the possibility of invisibly decaying Higgs bosons \(^9\).

The \( \epsilon_i \) parameters are then equal to some Yukawa coupling times \( v_R \). This provides the main theoretical motivation for introducing explicitly BRpV in the MSSM superpotential. From a practical point of view it provides the most predictive approach to the violation of R–Parity, which renders possible the systematic study of its phenomenological implications \(^10\). Here I will mention the most important features of this model.

For simplicity we set 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

\[ W = \varepsilon_{ab} \left[ h_t \tilde{Q}_3^a \tilde{U}_3^b \hat{H}_2^b + h_\tau \tilde{L}_3^a \tilde{D}_3^b \hat{H}_1^b + h_\tau \tilde{L}_3^b \tilde{R}_3^a \hat{H}_1^a - \mu \tilde{H}_1^a \tilde{H}_2^b + \epsilon_3 \tilde{L}_3^a \tilde{H}_2^b \right], \]  

(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

\(^a\) 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.
\[ \langle \hat{\nu}_\tau \rangle = v_3 / \sqrt{2} \] is non-zero. This in turn generates more R-parity and tau lepton number violating terms which, as we will see later, induce a tau neutrino mass.

It has often been claimed, by looking at the last two terms in the superpotential, that the BRpV term be rotated away from the superpotential by a suitable choice of basis. If this were true the \( \epsilon \) term would be unphysical. Indeed, consider the following rotation of the superfields

\[
\hat{H}_1' = \frac{\mu \hat{H}_1 - \epsilon_3 \hat{L}_3}{\sqrt{\mu^2 + \epsilon_3^2}}, \quad \hat{L}_3' = \frac{\epsilon_3 \hat{H}_1 + \mu \hat{L}_3}{\sqrt{\mu^2 + \epsilon_3^2}}.
\]

The superpotential in the new basis is

\[
W = h_1 \hat{Q}_3 \hat{U}_3 \hat{H}_2 + h_b \frac{\mu}{\mu'} \hat{Q}_3 \hat{D}_3 \hat{H}_1' + h_\tau \hat{L}_3' \hat{R}_3 \hat{H}_2' - \mu' \hat{H}_1' \hat{H}_2 + h_b \frac{\epsilon_3}{\mu'} \hat{Q}_3 \hat{D}_3 \hat{L}_3',
\]

where \( \mu'' = \mu^2 + \epsilon_3^2 \). The first four terms are MSSM-like terms and the last term violates the R–Parity defined in the new basis. Notice that, although the \( \epsilon \) term disappears from the superpotential in the new basis, R–Parity is reintroduced in the form of TRpV. Moreover, supersymmetry must be broken and this is parametrized by soft supersymmetry breaking terms. The soft terms which play an important role in BRpV are the following

\[
V_{soft} = m_{H_1}^2 |H_1|^2 + M_{L_3}^2 |\bar{L}_3|^2 - \left[ B \mu H_1 H_2 - B_2 \epsilon_3 \bar{L}_3 H_2 + h.c. \right] + ...
\]

where \( m_{H_1}^2 \) and \( M_{L_3}^2 \) are the soft masses corresponding to \( H_1 \) and \( \bar{L}_3 \) respectively, while \( B \) and \( B_2 \) are the bilinear soft mass parameters associated to the next-to-last and last terms in the superpotential in eq. (4). 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. (5).

The soft terms in the rotated basis are given by

\[
V_{soft} = \frac{m_{H_1}^2 \mu^2 + M_{L_3}^2 \epsilon_3^2}{\mu''^2} |H_1'|^2 + \frac{m_{H_1}^2 \epsilon_3^2 + M_{L_3}^2 \mu^2}{\mu''^2} |\bar{L}_3'|^2 - \frac{B \mu^2 + B_2 \epsilon_3^2}{\mu''} H_1' H_2
\]

\[ + \frac{\epsilon_3 \mu}{\mu''} (m_{H_1}^2 - M_{L_3}^2) \bar{L}_3' H_1' + \frac{\epsilon_3 \mu}{\mu''} (B_2 - B) \bar{L}_3' H_2 + h.c. + ... \]

(6)

The first three terms are MSSM-like terms, equivalent to the first three terms in eq. (3). In fact, in analogy with the MSSM, the coefficients of \( |H_1'|^2 \) and \( |\bar{L}_3'|^2 \) could be defined in the rotated basis as the soft masses \( m_{H_1}'^2 \) and \( M_{L_3}'^2 \).b

bFor three generations there is also a \( \hat{R} \hat{L} \hat{L}' \) term.
where χ massless isodoublet pseudoscalar majoron.

The linear terms of the scalar potential are then

\[ V = \epsilon_1 v_1^2 + \epsilon_2 v_2^2 + \epsilon_3 v_3^2 + \frac{1}{8} (g^2 + g'^2) v_1 (v_1^2 - v_2^2 + v_3^2) = 0, \]

\[ t_2 = (m^2_{H_2} + \mu^2 + \epsilon_3) v_2 - B_1 v_1 + B_2 \epsilon_3 v_3 - \frac{1}{8} (g^2 + g'^2) v_2 (v_1^2 - v_2^2 + v_3^2) = 0, \]

\[ t_3 = (M^2_{L_3} + \epsilon_3) v_3 - \mu \epsilon_3 v_1 + B_2 \epsilon_3 v_2 + \frac{1}{8} (g^2 + g'^2) v_3 (v_1^2 - v_2^2 + v_3^2) = 0. \]  

(7)

The \( t_i \) are the tree level tadpoles and they are equal to zero at the minimum. The linear terms of the scalar potential are then \( V_{\text{linear}} = t_1 \chi_1^0 + t_2 \chi_2^0 + t_3 \tilde{\nu}_R^0 \), where \( \chi_1^0 = \sqrt{2} \text{Re}(H_1) - v_1 \) and \( \tilde{\nu}_R^0 = \sqrt{2} \text{Re}(\tilde{\nu}_R) - v_3 \). The first two equations reduce to the MSSM minimization conditions after taking the MSSM limit \( \epsilon_3 = v_3 = 0 \), and in this case, the third equation is satisfied trivially. Note that \( \epsilon_3 = 0 \) implies two solutions for \( v_3 \) from the third tadpole in eq. (7), from which only \( v_3 = 0 \) is viable, since the second solution would imply the existence of a massless isodoublet pseudoscalar majoron.

The first two tadpole equations in the rotated basis are

\[ t'_1 = \mu'^2 v'_1 + \frac{m^2_{H_1} \mu^2 + M^2_{L_3} \epsilon_3^2}{\mu'^2} v'_1 - \frac{B_1 \mu^2 + B_2 \epsilon_3^2}{\mu'} v_2 + \frac{(m^2_{H_1} - M^2_{L_3}) \epsilon_3 \mu}{\mu'^2} v'_3 + \frac{1}{8} (g'^2 + g^2) v'_1 (v'_1^2 - v'_2^2 + v'_3^2) = 0, \]

(8)

\[ t'_2 = \mu'^2 v_2 + m^2_{H_2} v_2 - \frac{B_1 \mu^2 + B_2 \epsilon_3^2}{\mu'} v'_1 + \frac{(B_2 - B) \epsilon_3 \mu}{\mu'} v'_3 - \frac{1}{8} (g'^2 + g^2) v_2 (v_1^2 - v_2^2 + v_3^2) = 0, \]

(9)

where \( \langle H'_1 \rangle = v'_1 / \sqrt{2} \) and \( \langle L'_3 \rangle = v'_3 / \sqrt{2} \) with \( v'_1 = (\mu v_1 - \epsilon_3 v_3) / \mu' \) and \( v'_3 = (\epsilon_3 v_1 + \mu v_3) / \mu' \), as suggested by eq. (3). These two equations resemble the MSSM minimization conditions when we set \( v'_3 = 0 \). The third tadpole equation is

\[ t'_3 = (m^2_{H_1} - M^2_{L_3}) \frac{\epsilon_3 \mu}{\mu'^2} v'_1 + \frac{(B_2 - B) \epsilon_3 \mu}{\mu'} v_2 + \frac{m^2_{H_1} \epsilon_3^2 + M^2_{L_3} \mu^2}{\mu'^2} v'_3 + \frac{1}{8} (g'^2 + g^2) v'_3 (v'_1^2 - v'_2^2 + v'_3^2) = 0, \]

(10)
In this equation we note that \( v'_3 = 0 \) if \( \Delta m^2 \equiv m_{H_1}^2 - M_{L_3}^2 = 0 \) and \( \Delta B \equiv B_2 - B = 0 \) at the weak scale. 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 \( M_{GUT} \approx 2 \times 10^{16} \) GeV, but radiative corrections lifts this degeneracy due to the running of the renormalization group equations (RGE) between \( M_{GUT} \) and \( M_{\text{weak}} \). 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

\[
v'_3 \approx -\frac{\epsilon_3 \mu}{\mu'^2 m_{\tilde{\nu}_\tau}^2} (v_1' \Delta m^2 + \mu' v_2 \Delta B)
\]

(11)

where we have introduced

\[
m_{\tilde{\nu}_\tau}^2 \equiv m_{H_1}^2 e_3^2 + M_{L_3}^2 \mu'^2 + \frac{1}{8} (g'^2 + g^2) (v_1'^2 - v_2'^2)
\]

(12)

which reduces to the tau sneutrino mass in the MSSM when we set \( \epsilon_3 = 0 \). Note that eq. (11) implies that the R–parity-violating effects induced by \( v'_3 \) are calculable in terms of the primordial effective R–parity-violating \( \epsilon_3 \).

3 Neutrino Mass

The presence of tau lepton number and BRpV terms, characterized by the parameters \( \epsilon_3 \) and \( v_3 \), leads to a mixing between neutralinos and the tau neutrino \( \tilde{\nu}_\tau \) as a result of which the tau neutrino acquires a mass \( m_{\nu_\tau} \). In the original basis, where \( (\psi^0)^T = (-i \lambda', -i \lambda^3, H_1^0, H_2^0, \nu_\tau) \), the scalar potential contains the following mass terms

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

(13)

where the neutralino/neutrino mass matrix is

\[
M_N = \begin{bmatrix}
  M' & 0 & -\frac{1}{2} g' v_1 & \frac{1}{2} g' v_2 & -\frac{1}{2} g' v_3 \\
  0 & M & \frac{1}{2} g v_1 & -\frac{3}{2} g v_2 & \frac{1}{2} g v_3 \\
  -\frac{1}{2} g' v_1 & \frac{1}{2} g v_1 & 0 & -\mu & x \epsilon_3 \\
  \frac{1}{2} g' v_2 & -\frac{3}{2} g v_2 & -\mu & 0 & \epsilon_3 \\
  -\frac{1}{2} g' v_3 & \frac{1}{2} g v_3 & 0 & \epsilon_3 & 0
\end{bmatrix}
\]

(14)

where \( M \) 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 very small with respect to the weak scale and, in fact, this has often been claimed as a way to dismiss the phenomenological relevance of R–parity violation. However, the cosmological critical density bound \( m_{\nu_\tau} < 920h^2 \text{ eV} \) only holds if neutrinos are stable. In the present BRpV model (where there is no majoron) the \( \nu_\tau \) can decay into 3 neutrinos, via the neutral current \( 13 \), or by slepton exchanges. This mechanism may be efficient in reducing the relic \( \nu_\tau \) abundance below the required level, as long as \( \nu_\tau \) is heavier than about 100 keV or so. On the other hand primordial Big-Bang nucleosynthesis implies that \( \nu_\tau \) is lighter than about an Mev or so. Thus, in addition to the electron-volt neutrino mass range, we obtain another region, say between .1 to 1 MeV where heavy \( \nu_\tau \) masses are cosmologically consistent in the BRpV model. Needless to say, in the spontaneous breaking version of the model all masses up to the LEP limit are cosmologically consistent due to the majoron-induced decay and annihilation channels\[14\].

Let us now compare the cosmologically allowed values of the tau neutrino mass with the theoretically predicted ones. In order to do this we embed our MSSM–BRpV model into supergravity, with universality of scalar \( (m_0) \), gaugino \( (\tilde{M}_{1/2}) \), bilinear \( (B) \), and trilinear \( (A) \) soft mass parameters at the unification scale \( M_X \approx 2 \times 10^{16} \text{ GeV} \). The expected \( m_{\nu_\tau} \) values are illustrated in Fig. 1 where 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 \( v_3 \), imposing the LEP limit on \( m_{\nu_\tau} \) and that the supersymmetric particles are not too light. As one can see there is a strict correlation between the neutrino mass and the magnitude of R–parity-violation given by \( \epsilon_3 \) and \( v_3 \) which need not be small. We have explicitly verified that \( |\epsilon_3| \) can be as large as 400 GeV and that \( |v_3| \) can be close to 100 GeV without conflicting with laboratory limits on the tau neutrino mass. An obvious question arises at this stage: how can we get a small neutrino mass, if so desired? The answer lies in the fact that the induced neutrino mass is \( \propto (\epsilon_3 v_1 + \mu v_3)^2 \), and this last combination is what needs to be small. One can see that the contributions to \( m_{\nu_\tau} \) coming from Higgsino and gaugino mixing, which are proportional to \( \epsilon_3 \) and \( v_3 \) may nearly cancel, leading to a mass that can be very small, in the eV or sub-eV range. How natural is such a cancellation? We have found that if our model is unified a la supergravity 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 order to appreciate this better we rewrite neutral mass matrix in eq. (14)
Figure 1: Tau neutrino mass as a function of the effective R–parity-violating parameter \( \xi \equiv (\epsilon_3 v_1 + \mu v_3)^2 \).

in the rotated basis. This corresponds to 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

\[
m_{\nu_\tau} \approx -\frac{(g^2 M + g'^2 M')\mu' v'^2_3}{4 M M' \mu'^2 - 2(g^2 M + g'^2 M')v'_1 v_2 \mu'}
\]

Now solving the renormalization group equations for the soft mass parameters \( m^2_{H_1}, m^2_{L_3}, B, \) and \( B_2 \), in first approximation we get

\[
m^2_{H_1} - M^2_{L_3} \approx -\frac{3h_b^2}{8\pi^2} \left( m^2_{H_1} + M^2_Q + M^2_D + A_D^2 \right) \ln \frac{M_{GUT}}{m_Z} \\
B_2 - B \approx \frac{3h_b^2}{8\pi^2} A_D \ln \frac{M_{GUT}}{m_Z}
\]

Using eq. (14) we can show that the sneutrino vev \( v'_3 \) given through \( \xi \equiv (\epsilon_3 v_1 + \mu v_3)^2 \) is radiatively generated, with a maximum value of few GeV or so. The resulting tau neutrino mass is given by

\[
m_{\nu_\tau} \approx \frac{[\mu' v_2 A_D - v'_1 \left( m^2_{H_1} + M^2_Q + M^2_D + A_D^2 \right)]^2}{2v'_1 v_2 - 4MM' \mu'/\left( g^2 M + g'^2 M' \right)} \mu'' m^2_{\tilde{\nu}_0} \left( \frac{\epsilon_3 \mu}{\mu'^2} \right)^2 \left( \frac{3h_b^2}{8\pi^2} \ln \frac{M_{GUT}}{m_Z} \right)^2
\]
This mass can be further approximated by

\[ m_{\nu_{\tau}} \approx \frac{m^2_Z}{M_{\text{SUSY}}} \left( \frac{\epsilon_3}{M_{\text{SUSY}}} \right)^2 \tilde{h}_b^4 \sim 1 \text{ eV} \]  

(18)

where we have explicitly indicated that 1 eV is a perfectly viable \( m_{\nu_{\tau}} \) value in this model. This was obtained for \( M_{\text{SUSY}} \sim \epsilon_3 \sim m_Z \) and \( \tilde{h}_b \sim 10^{-2} \). Therefore, \( m_{\nu_{\tau}} \) can be naturally small, even though the R–parity-violating parameters are large. The actual scale on neutrino mass can, of course, be larger, as the smallness of \( m_{\nu_{\tau}} \) is tightly related to our soft SUSY breaking terms universality assumption at the unification scale.

4 Unification of Couplings

Unification of the gauge couplings in our model works basically as in the MSSM. In contrast, Yukawa coupling unification is rather different. To carry out this discussion we start from the basic superpotential in eq. (2). Similarly to neutralino-neutrino mixing, charginos also mix with the tau lepton, forming a set of three charged fermions \( F^\pm_i, i = 1, 2, 3 \). In the original basis where \( \psi^{+T} = (\mp i \lambda^+, \tilde{H}^+_2, \tau^+_R) \) and \( \psi^{-T} = (-i \lambda^-, \tilde{H}^-_2, \tau^-_L) \), the charged fermion mass terms in the Lagrangian are

\[ L_m = -\psi^{-T} M_C \psi^+ \]

with the mass matrix given by

\[ M_C = \begin{bmatrix}
\frac{1}{\sqrt{2}} g v_2 & \mu & -\frac{1}{\sqrt{2}} h\tau v_3 \\
\frac{1}{\sqrt{2}} g v_2 & -\epsilon_3 & \frac{1}{\sqrt{2}} h\tau v_1 \\
\frac{1}{\sqrt{2}} g v_2 & -\frac{1}{\sqrt{2}} h\tau v_3 & -\epsilon_3
\end{bmatrix} \]  

(19)

As a result, the tau Yukawa coupling is not related to the tau mass by the usual MSSM relation. In contrast, \( h\tau \) depends now on the parameters of the chargino sector \( M, \mu, \) and \( \tan \beta \), as well as the BRpV parameters \( \epsilon_3 \) and \( v_3 \), through a formula given in ref. [16]. In addition, the top and bottom quark Yukawa couplings are related to the quark masses in a way different from that of the MSSM due to the non-zero value of \( v_3/v \), i.e.

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

(20)

where \( v = 246 \text{ GeV} \) and we have defined \( \cos \theta \equiv v_3/v \).

The re-scaling in the bottom quark Yukawa term ensures that the same quark mass is obtained with the same Yukawa coupling in the two basis. This re-scaling with respect to the MSSM is non-trivial and has profound consequences in Yukawa unification, as shown in Fig. [17] taken from ref. [17]. 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 $v_3$. The horizontal lines correspond to the $1\sigma$ experimental determination of $m_t$. The plot in Fig. 2 is obtained through a scan over parameter space such that points which satisfy $h_b(M_{\text{GUT}}) = h_\tau(M_{\text{GUT}})$ within 1% are kept, where $M_{\text{GUT}}$ is the gauge coupling unification scale. Each selected point is placed in one of the regions of Fig. 2 according to its $|v_3|$ value. However, one can see that, for very large $v_3 > \sim 40$ GeV the large $\tan \beta$ solution is ruled out. Points with top-bottom-tau unification are concentrated in the diagonal line at high values of $\tan \beta$, analogously to the MSSM case.

In summary, BRpV is the simplest way to introduce R–Parity violation to the MSSM. The model can be embedded into Supergravity models with universality of scalar, gaugino, bilinear and trilinear soft mass parameters. In this case, the induced tau neutrino mass arises radiatively and is naturally small. The BRpV parameters $\epsilon_3$ and $v_3$ need not be small and can be easily of the order of $m_Z$. Another important feature is that BRpV changes the relation between the Yukawa couplings and the masses of the top and bottom quarks and the tau lepton. As a result, bottom-tau Yukawa unification can be achieved for any $\tan \beta$ value, provided we choose appropriately the value of the sneutrino
vev $v_3$. Even in the unlikely limit where the tau neutrino is massless with $\epsilon_3 \neq 0$ (which corresponds to having universality of soft mass parameters at the weak scale!, which is not natural) R-Parity is not conserved. In fact, 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 due to the last term in eq. (4). Thus R-parity violation can be sizeable even if neutrinos turn out to be very light, as indicated by present solar and atmospheric neutrino data. Some of the phenomenological implications of the model have been discussed in ref. [4].

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1. H.P. Nilles, *Phys. Rep.* **110**, 1 (1984); H.E. Haber and G.L. Kane, *Phys. Rep.* **117**, 75 (1985); R. Barbieri, *Riv. Nuovo Cimento* **11**, 1 (1988).
2. U. Amaldi, W. de Boer, and H. Furstenau, *Phys. Lett. B* **260**, 447 (1991); J. Ellis, S. Kelley, and D.V. Nanopoulos, *Phys. Lett. B* **260**, 131 (1991); P. Langacker and M. Luo, *Phys. Rev. D* **44**, 817 (1991); C. Giunti, C.W. Kim and U.W. Lee, *Mod. Phys. Lett.* **A6**, 1745 (1991).
3. P. Langacker and N. Polonsky, *Phys. Rev. D* **47**, 4028 (1993); P.H. Chankowski, Z. Pluciennik, and S. Pokorski, *Nucl. Phys. B* **439**, 23 (1995); P.H. Chankowski, Z. Pluciennik, S. Pokorski, and C.E. Vayonakis, *Phys. Lett. B* **358**, 264 (1995).
4. V. Berezinskii, Anjan S. Joshipura, Jose W.F. Valle. *Phys.Rev.D* **57**, 147-151 (1998)
5. F. de Campos, M.A. García-Jareño, A.S. Joshipura, J. Rosiek, and J.W.F. Valle, *Nucl. Phys.* **B451**, 3 (1995); A.S. Joshipura and M. Nowakowski, *Phys. Rev. D* **51**, 2421 (1995); T. Banks, Y. Grossman, E. Nardi, and Y. Nir, *Phys. Rev. D* **52**, 5319 (1995); F. Vissani and A. Yu. Smirnov, *Nucl.Phys. B460*, 37 (1996); R. Hempfling, *Nucl. Phys. B478*, 3 (1996); F.M. Borzumati, Y. Grossman, E. Nardi, Y. Nir, *Phys. Lett. B* **384**, 123 (1996); H.P. Nilles and N. Polonsky, *Nucl. Phys. B484*, 33 (1997); B. de Carlos, P.L. White, *Phys. Rev. D* **55**, 4222 (1997); E. Nardi, *Phys. Rev. D* **55**, 5772 (1997); S. Roy and B. Mukhopadhyaya, *Phys. Rev. D* **55**, 7020 (1997); A. Faessler, S. Kovalenko, F. Simkovic, hep-ph/9712362; M. Carena, S. Pokorski, and C.E.M. Wagner, hep-ph/9801251; M.E. Gómez and K. Tamvakis, hep-ph/9801348.
6. M.A. Díaz, hep-ph/9712533; M.E. Gómez and K. Tamvakis, hep-ph/9801348.
7. A. Masiero and J.W.F. Valle, *Phys. Lett.* **B251**, 273 (1990); J.C. Romão, A. Ioannissyan and J.W.F. Valle, *Phys. Rev. D** **55**, 427 (1997).
8. T. Multamaki, I. Vilja, [hep-ph/9804371](http://arxiv.org/abs/hep-ph/9804371)
9. A. S. Joshipura, J.W.F. Valle, *Nucl.Phys. B** **397** 105-122 (1993), for a study of the phenomenological implications see F. de Campos, O.J.P. Eboli, J. Rosiek, J.W.F. Valle, *Phys.Rev. D** **55**, 1316-1325 (1997), [hep-ph/9612692](http://arxiv.org/abs/hep-ph/9612692); L3 Collaboration (M. Acciarri et al.). CERN-PPE-97-097, *Phys.Lett. B** **418** 389-398 (1998).
10. J.W.F. Valle, [hep-ph/9712277](http://arxiv.org/abs/hep-ph/9712277), [hep-ph/9802292](http://arxiv.org/abs/hep-ph/9802292) and [hep-ph/9603307](http://arxiv.org/abs/hep-ph/9603307)
11. L. Hall and M. Suzuki, *Nucl.Phys. B** **231**, 419 (1984).
12. G.G. Ross, J.W.F. Valle. *Phys.Lett. B** **151B** 375 (1985); John Ellis, G. Gelmini, C. Jarlskog, G.G. Ross, J.W.F. Valle, *Phys.Lett. B** **150B** 142 (1985); A. Santamaria, J.W.F. Valle, *Phys.Lett. B** **195B** 423 (1987), *Phys.Rev.Lett. 60*:397-400 (1988) and *Phys.Rev. D** **39** 1780-1783 (1989).
13. J. Schechter, J.W.F. Valle, *Phys.Rev. D** **22** 2227 (1980)
14. A.D. Dolgov, S. Pastor, J.C. Romao, J.W.F. Valle. *Nucl.Phys. B** **496** 24-40 (1997), [hep-ph/9610507](http://arxiv.org/abs/hep-ph/9610507)
15. M.A. Díaz, J.C. Romão, and J.W.F. Valle, *Nucl.Phys. B** **524** 23-40 (1998); [hep-ph/9706313](http://arxiv.org/abs/hep-ph/9706313).
16. A. Akeroyd, M.A. Díaz, J. Ferrandis, M.A. García–Jareño, and J.W.F. Valle, [hep-ph/9707393](http://arxiv.org/abs/hep-ph/9707393); M.A. Díaz, [hep-ph/9710233](http://arxiv.org/abs/hep-ph/9710233); *Nucl.Phys. B*, xxx (1998), in press.
17. M.A. Díaz, J. Ferrandis, J.C. Romão, and J.W.F. Valle, [hep-ph/9801391](http://arxiv.org/abs/hep-ph/9801391).