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Tau Lepton Mixing with Charginos and its Effects on Chargino searches at e^+e^- Colliders

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Abstract

In bilinear R-Parity violating models where a term $\epsilon_3 L_3 H_2$ is introduced in the superpotential, the tau lepton can mix with charginos. We show that this mixing is fully compatible with LEP1 precision measurements of the $Z\tau\tau$ and $W\tau\nu_\tau$ couplings even for large values of ϵ_3 and of the induced vacuum expectation value v_3 of the tau-sneutrino. The single production of charginos at e^+e^- colliders is possible in this case and we present numerical values of the cross-section at LEP1, LEP2 and an NLC. We find maximum values of 10 pb at LEP1 and 1 fb at NLC, while the corresponding values at LEP2 are too small to observe.

1 Introduction

The Minimal Supersymmetric Standard Model (MSSM) [1] is the most popular extension of the Standard Model (SM) and its phenomenology has received much attention in recent years. The MSSM assumes the conservation of R-parity [2], a discrete symmetry given by $R_p = (-1)^{(3B+L+2S)}$, where L is the lepton number, B is the baryon number and S is the spin of the state. Two important consequences of such a symmetry are that all supersymmetric particles must be pair-produced, with the lightest of them being stable. In recent years growing attention has been given to models in which the conservation of R-Parity is relaxed [3] by adding explicit R-parity violating terms in the superpotential. In the literature one finds many studies using models which possess cubic Rp violating superpotential terms, which in turn introduce a very large number of arbitrary Yukawa couplings [4]. Such models suffer from the fact that their phenomenological study involves a large number of free parameters. Some of these must be strongly suppressed in order to avoid conflict with nucleon stability. In this work we shall first focus on Bilinear R-Parity Violation (BRpV) which has been advocated in a number of previous papers [5, 6, 7, 8]. Such models allow one to map out the phenomenological potential of the model in a systematic fashion [9, 10]. They have the theoretical advantage of being a low energy approximation of spontaneous R-parity violating models. These may have either gauged [11] or ungauged [12] lepton number. In the latter case the particle spectrum will include a massless majoron [13]. The presence of such a particle allows one to avoid the stringent cosmological nucleosynthesis bound of $m_{\tau\nu} \leq 1 - 2$ MeV or so [9], and instead employ the weaker mass limit found from direct searches at LEP1, $m_{\tau\nu} \leq 18$ MeV. However, in order to avoid excessive stellar cooling through majoron emission the value of the sneutrino vacuum expectation value (v_3) is then constrained to small values (≤ 100 MeV) [13]. In contrast, in the simplest model which *explicitly* breaks R-parity via a bilinear superpotential term v_3 is only constrained by the W mass formula, and may be comfortably varied up to values of order 100 GeV.

The Tau lepton in the BRpV or in spontaneous Rp breaking models is allowed to mix with the charginos, an effect dependent on the values of ϵ_3 and v_3 . Such mixing, if substantial, would naively be expected to affect the theoretical value of the $Z\tau\tau$ coupling which is measured to very high accuracy at LEP1 (error $\approx 0.25\%$). Thus it is important to check if one may constrain the allowed parameter space of ϵ_3 and v_3 , and we do this in the first part of this paper. A consequence of the above mixing is the possibility of single chargino production at e^+e^- colliders, a process not possible in Rp conserving models. We update and improve previous analyses [11, 14] of this channel at LEP1 and investigate the phenomenology at LEP2 and a NLC. The paper is organized as follows. In Section 2 the essential properties of the BRpV model are reviewed and we present the relevant

mass matrices. In Section 3 we study the effect of the chargino–tau mixing on the $Z\tau\tau$ and $W\tau\nu_\tau$ couplings, while Section 4 analyses the prospects of single chargino production at e^+e^- colliders. Section 5 contains our conclusions.

2 The Model

In the simplest BRpV model the supersymmetric Lagrangian is specified by the superpotential W given by

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

where $i, j = 1, 2, 3$ are generation indices, $a, b = 1, 2$ are $SU(2)$ indices. The last term in eq. (1) is the only one which violates R -parity. Such a term arises in spontaneous R -parity breaking models, with ϵ_i generated by the product of a new Dirac-type neutrino Yukawa coupling and some new vacuum expectation value (VEV) of an $SU(2)$ singlet sneutrino field. The parameters ϵ_i have the dimension of mass, and for simplicity we consider only ϵ_3 non-zero. Of particular interest to us is the chargino/tau mass matrix, which is given by

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

and M is the $SU(2)$ gaugino soft mass. The form of this matrix is common to all models with spontaneous breaking of Rp, as well as in the simplest truncation of these models provided by the BRpV model. We note that the chargino sector decouples from the tau sector in the limit $\epsilon_3 = v_3 = 0$. As in the MSSM, the chargino mass matrix is diagonalized by two rotation matrices \mathbf{U} and \mathbf{V}

$$\mathbf{U}^* \mathbf{M}_C \mathbf{V}^{-1} = \begin{bmatrix} m_{\chi_1^\pm} & 0 & 0 \\ 0 & m_{\chi_2^\pm} & 0 \\ 0 & 0 & m_\tau \end{bmatrix}. \quad (3)$$

The lightest eigenstate of this mass matrix must be the tau lepton (τ^\pm) and so the mass is constrained to be 1.777 GeV. To obtain this the tau Yukawa coupling becomes a function of the parameters in the mass matrix, and the full expression is given in [6]. The composition of the tau is given by:

$$\tau_R^+ = V_{3j} \psi_j^+, \quad \tau_L^- = U_{3j} \psi_j^- \quad (4)$$

where $\psi^{+T} = (-i\lambda^+, \widetilde{H}_2^1, \tau_R^{0+})$ and $\psi^{-T} = (-i\lambda^-, \widetilde{H}_1^2, \tau_L^{0-})$. The two component Weyl spinors τ_R^{0-} and τ_L^{0+} are weak eigenstates and, similarly, the two component Weyl spinors τ_R^+ and τ_L^- are mass eigenstates.

It follows easily from eq. (3) that the matrix $\mathbf{M}_C \mathbf{M}_C^T$ is diagonalized by \mathbf{U} and the matrix $\mathbf{M}_C^T \mathbf{M}_C$ is diagonalized by \mathbf{V} . It is instructive to write the matrices $\mathbf{M}_C \mathbf{M}_C^T$ and $\mathbf{M}_C^T \mathbf{M}_C$ explicitly since they differ significantly in appearance, particularly in the off diagonal elements which depend on the R-parity violating couplings:

$$\mathbf{M}_C \mathbf{M}_C^T = \begin{bmatrix} M^2 + \frac{1}{2}g^2v_2^2 & \frac{1}{\sqrt{2}}g(Mv_1 + \mu v_2) & \frac{1}{\sqrt{2}}g(Mv_3 - \epsilon_3 v_2) \\ \frac{1}{\sqrt{2}}g(Mv_1 + \mu v_2) & \mu^2 + \frac{1}{2}g^2v_1^2 + \frac{1}{2}h_\tau^2v_3^2 & -\mu\epsilon_3 + \frac{1}{2}v_1v_3(g^2 - h_\tau^2) \\ \frac{1}{\sqrt{2}}g(Mv_3 - \epsilon_3 v_2) & -\mu\epsilon_3 + \frac{1}{2}v_1v_3(g^2 - h_\tau^2) & \frac{1}{2}h_\tau^2v_1^2 + \epsilon_3^2 + \frac{1}{2}g^2v_3^2 \end{bmatrix} \quad (5)$$

and

$$\mathbf{M}_C^T \mathbf{M}_C = \begin{bmatrix} M^2 + \frac{1}{2}g^2(v_1^2 + v_3^2) & \frac{1}{\sqrt{2}}g(Mv_2 + \mu v_1 - \epsilon_3 v_3) & 0 \\ \frac{1}{\sqrt{2}}g(Mv_2 + \mu v_1 - \epsilon_3 v_3) & \mu^2 + \frac{1}{2}g^2v_2^2 + \epsilon_3^2 & -\frac{1}{\sqrt{2}}h_\tau(\mu v_3 + \epsilon_3 v_1) \\ 0 & -\frac{1}{\sqrt{2}}h_\tau(\mu v_3 + \epsilon_3 v_1) & \frac{1}{2}h_\tau^2(v_1^2 + v_3^2) \end{bmatrix} \quad (6)$$

One notices a pair of zeros in the R-Parity violating elements of the matrix $\mathbf{M}_C^T \mathbf{M}_C$. The second pair of elements which violate R-Parity turn out to be small. This can be seen from the fact that the term $(\mu v_3 + v_1 \epsilon_3)$ is naturally small since its square is proportional to the mass of ν_τ [6, 7, 10]. Indeed, the above combination satisfy $\mu v_3 + v_1 \epsilon_3 = \mu' v'_3$, where $\mu' = \sqrt{\mu^2 + \epsilon_3^2}$ and v'_3 is the vacuum expectation value (VEV) of the sneutrino field in a rotated basis where the ϵ_3 term disappear from the superpotential (although reintroduced in the soft terms). In models with universality of soft masses at the unification scale, this VEV v'_3 is radiatively generated and proportional to the bottom quark Yukawa coupling squared. On the other hand, the mixing between the neutralinos and the tau-neutrino in the rotated basis is proportional to v'_3 , therefore a neutrino mass is induced and satisfy $m_{\nu_\tau} \sim v_3'^2$. In summary, $m_{\nu_\tau} \sim v_3'^2 \sim (\mu v_3 + v_1 \epsilon_3)^2$ is naturally small since it is one-loop induced and proportional to h_b^2 . This makes the R-Parity violating couplings V_{31} and V_{32} also naturally small and controlled by m_{ν_τ} .

In contrast, the R-Parity violating elements in the matrix $\mathbf{M}_C \mathbf{M}_C^T$ may be of greater magnitude. In this way, U_{31} and U_{32} may be larger than their similars V_{31} and V_{32} . Nevertheless, a closer look tells us that in the limit $(\mu v_3 + v_1 \epsilon_3) \rightarrow 0$ (*i.e.*, massless neutrino) we find $U_{31} \rightarrow 0$. Therefore, U_{31} is also small and controlled by m_{ν_τ} . On the contrary, U_{32} can be larger. In the next section we show that this is not in conflict with the τ or ν_τ couplings to gauge bosons.

3 The τ couplings to gauge bosons

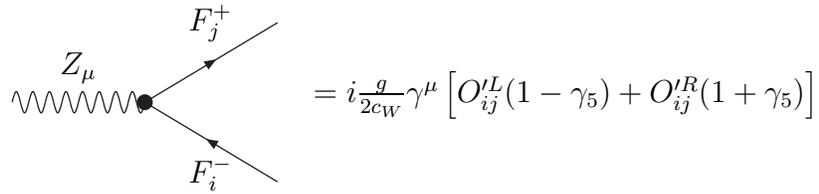
The pair production of τ leptons in the SM proceeds via two tree level Feynman diagrams, *i.e.*, the s-channel mediated by a photon and a Z boson. In BRpV there is an extra diagram, namely that of t-channel production mediated by a tau-sneutrino. This diagram

arises because in general the τ has a gaugino component. One may attempt to constrain the chargino/tau mixing in eq. (2) by using precision measurements at LEP1 of the $Z\tau\tau$ and $W\tau\nu_\tau$ couplings. The current experimental values of the axial vector (g_A^τ) and the vector (g_V^τ) parts of the $Z\tau\tau$ coupling are [15]:

$$g_A^\tau = -0.5009 \pm 0.0013, \quad g_V^\tau = 0.0374 \pm 0.0022 \quad (7)$$

whose tree level values are $g_A^\tau = -\frac{1}{2}$ and $g_V^\tau = \frac{1}{2} - 2s_W^2$. It is customary to write the coupling $Z\tau^+\tau^-$ in the MSSM in terms of the constants $g_L^\tau = -\frac{1}{2}(g_A^\tau + g_V^\tau)$ and $g_R^\tau = \frac{1}{2}(g_A^\tau - g_V^\tau)$ which are respectively the coupling strengths of the left and right handed τ to Z .

In BRpV g_L^τ and g_R^τ are the third diagonal elements of the matrices $O'_{ij}{}^L$ and $O'_{ij}{}^R$, which are 3×3 generalizations of the analogous 2×2 matrices in the MSSM. If we call F_i^\pm , $i = 1, 2, 3$, the three charged fermion mass eigenstates, of which the first two are the charginos and the third one is the tau lepton, then the $ZF_j^+F_i^-$ interactions are



$$= i \frac{g}{2c_W} \gamma^\mu \left[O'_{ij}{}^L (1 - \gamma_5) + O'_{ij}{}^R (1 + \gamma_5) \right]$$

with the couplings $O'_{ij}{}^L$ and $O'_{ij}{}^R$ given by

$$\begin{aligned} O'_{ij}{}^L &= -V_{i1}V_{j1}^* - \frac{1}{2}V_{i2}V_{j2}^* + \delta_{ij}s_W^2, \\ O'_{ij}{}^R &= -U_{i1}^*U_{j1} - \frac{1}{2}U_{i2}^*U_{j2} - \frac{1}{2}U_{i3}^*U_{j3} + \delta_{ij}s_W^2. \end{aligned} \quad (8)$$

In our notation the τ lepton is the “third chargino” and so to obtain the $Z\tau\tau$ coupling it is sufficient to put $i = j = 3$. The axial part g_A^τ of the coupling $Z\tau\tau$ is given by $O'_{33}{}^R - O'_{33}{}^L$ while the vector part g_V^τ is given by $-O'_{33}{}^R - O'_{33}{}^L$. Written explicitly one finds:

$$\begin{aligned} g_A^\tau &= -|U_{31}|^2 - \frac{1}{2}|U_{32}|^2 - \frac{1}{2}|U_{33}|^2 + |V_{31}|^2 + \frac{1}{2}|V_{32}|^2, \\ g_V^\tau &= |U_{31}|^2 + \frac{1}{2}|U_{32}|^2 + \frac{1}{2}|U_{33}|^2 + |V_{31}|^2 + \frac{1}{2}|V_{32}|^2 - 2s_W^2. \end{aligned} \quad (9)$$

Of course, in the R-parity conserving limit, that is $V_{31} = V_{32} = U_{31} = U_{32} = 0$, one recovers the formula for τ couplings in the MSSM.

We have evaluated the numerical value of these couplings for 10^4 randomly chosen points. All the points satisfy the following experimental mass limits:

$$m_{\chi^+} \geq 70 \text{ GeV} [16], \quad m_{\nu_\tau} \leq 18 \text{ MeV} [17] \quad m_{\chi^0} \geq 20 \text{ GeV}. \quad (10)$$

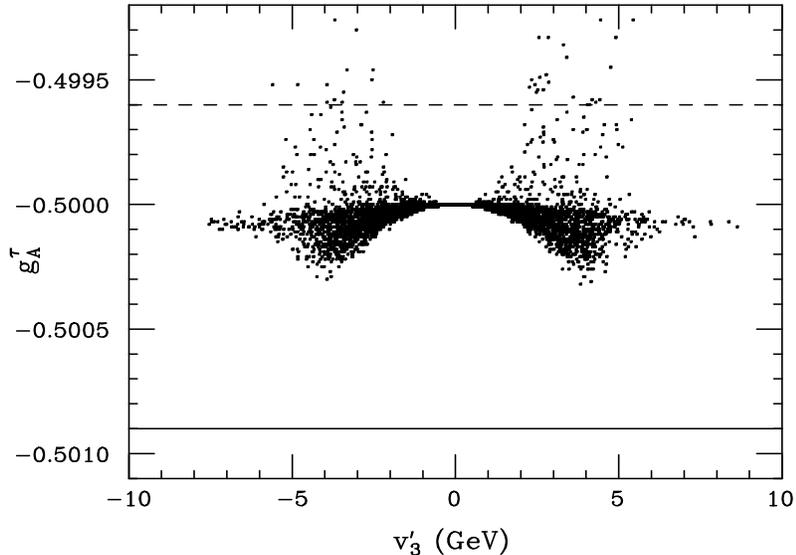


Figure 1: Axial vector coupling of the τ to a Z gauge boson as a function of the BRpV parameter v'_3 . The solid line is the experimental central value and the dashed line corresponds to 1σ deviation.

Note that, to be conservative, we applied assumed a lower bound on the neutralino mass of 20 GeV. This is reasonable in view of the work presented in ref. [18]¹. The couplings g_A^τ and g_V^τ are functions of 6 independent parameters which are varied in the following ranges:

$$\begin{aligned}
 -200 \text{ GeV} &\leq \mu, \epsilon_3 \leq 200 \text{ GeV} \\
 -100 \text{ GeV} &\leq v_3 \leq 100 \text{ GeV} \\
 0.5 &\leq \tan \beta \leq 90 \\
 30 \text{ GeV} &\leq M \leq 200 \text{ GeV} \\
 60 \text{ GeV} &\leq m_{\tilde{\nu}} \leq 200 \text{ GeV}
 \end{aligned} \tag{11}$$

In Fig. 1 we plot the axial vector coupling g_A^τ as a function of the sneutrino vacuum expectation value v'_3 in the rotated basis. The central value is given by $g_A^\tau = -0.5009$ (solid line) and the horizontal dashed line corresponds to 1σ deviation. Clearly, the great majority of the points fall within the LEP1 bound on g_A^τ at the 1σ level. Note that the values of ϵ_3 and v_3 can be large and constrained only by $v'_3 \lesssim 10$ GeV. Similarly, in Fig. 2 we plot the vector coupling g_V^τ as a function of v'_3 . The solid horizontal line at $g_V^\tau = 0.0374$ is the central value, and the dashed line denotes the 1σ deviation. In this case all points

¹Strictly speaking, however, there is not yet a published determination on the neutralino bounds from LEP2 in the bilinear model of broken Rp.

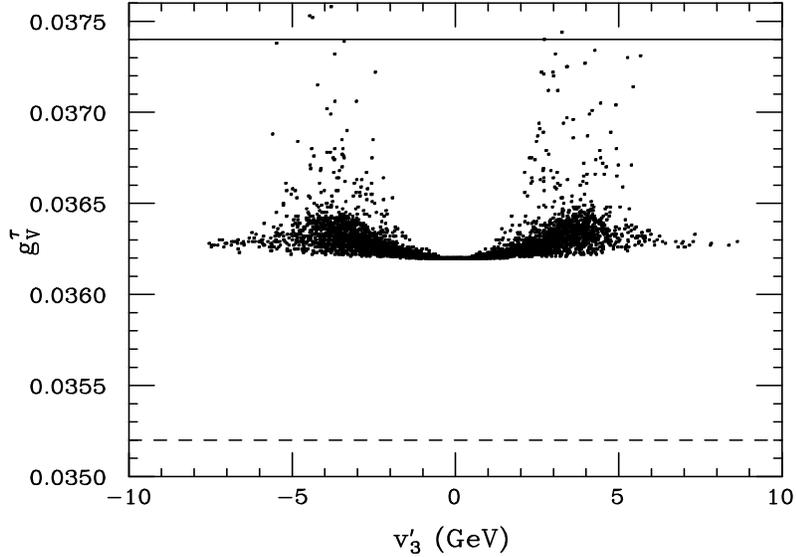


Figure 2: Vector coupling of the τ to a Z gauge boson as a function of the BRpV parameter v'_3 . The solid line is the experimental central value and the dashed line corresponds to 1σ deviation.

fall inside the 1σ region.

Let us now consider how the tau mixing with the charginos affects the $W\tau\nu_\tau$ vertex. Any change in this coupling would affect the τ decay widths which are measured very accurately. In our model there are five neutral fermions mass eigenstates which we denote F_i^0 , $i = 1, \dots, 5$, with the lightest (corresponding to $i = 5$) being the tau neutrino and the first four being the “neutralinos”. The 5×5 neutralino mass matrix is diagonalized by the matrix \mathbf{N}_{ij} . The vertex $WF_j^\pm F_i^0$ is

$$= i\frac{g}{2}\gamma^\mu \left[O_{ij}^L(1 - \gamma_5) + O_{ij}^R(1 + \gamma_5) \right]$$

with the generalized O^L and O^R couplings given by

$$\begin{aligned} O_{ij}^L &= -\frac{1}{\sqrt{2}}N_{i4}V_{j2}^* + N_{i2}V_{j1}^*, \\ O_{ij}^R &= \frac{1}{\sqrt{2}}N_{i3}^*U_{j2} + N_{i2}^*U_{j1} + \frac{1}{\sqrt{2}}N_{i5}^*U_{j3}. \end{aligned} \quad (12)$$

For the vertex $W\tau\nu_\tau$ one sets $i = 5$ and $j = 3$. In the Rp-conserving limit one would

obtain

$$\begin{aligned}
-\frac{1}{\sqrt{2}}N_{54}V_{32}^* + N_{52}V_{31}^* &\longrightarrow 0, \\
\frac{1}{\sqrt{2}}N_{53}^*U_{32} + N_{52}^*U_{31} + \frac{1}{\sqrt{2}}N_{55}^*U_{33} &\longrightarrow \frac{1}{\sqrt{2}}.
\end{aligned}
\tag{13}$$

With the same 10^4 randomly chosen points we find that the deviations from the SM values of the $W\tau\nu_\tau$ couplings are well inside the experimental error, and less pronounced than the corresponding deviations for the $Z\tau\tau$ vertex.

The reason why the deviations of the tau couplings to gauge bosons are small with respect to the SM predictions can be traced to two facts. First the tau–neutrino mass is small, and second the Higgs superfield H_1 and the tau–lepton superfield L_3 both possess the same $SU(2)$ quantum numbers.

Summarizing the results of this section, we conclude that the couplings $Z\tau\tau$ and $W\tau\nu_\tau$ can be easily maintained within the experimental bounds even for large values of ϵ_3 and v_3 . This has important consequences for the phenomenology of the BRpV model in general [9], with the most salient effect being the existence of Rp violating branching ratios for sparticles, the most dramatic example being the LSP. Moreover, Rp violation may also affect the expected mass spectrum in comparison with the Rp conserving case. One example recently considered in ref. [6] is the charged Higgs mass can be lowered below 80 GeV if ϵ_3 is large.

4 Single Chargino Production

In this section we consider the single chargino production in electron–positron annihilation. We study in turn the phenomenology at LEP1, LEP2, and NLC ($\sqrt{s} = 500$ GeV). In general the total cross–section consists of three distinct contributions given by:

$$\sigma(e^+e^- \rightarrow Z, \tilde{\nu}_\tau \rightarrow \tilde{\chi}_1^\pm \tau^\mp) = \sigma_Z + \sigma_{\tilde{\nu}} + \sigma_{\tilde{\nu}Z}
\tag{14}$$

Note that an intermediate photon does not contribute. Explicit expressions for these formulae can be found in ref. [19]. The generalization to BRpV is straightforward and is obtained by replacing the 2×2 matrices O'^L , O'^R , V and U by 3×3 matrices and summing over three “charginos”.

First, we consider the single chargino production at LEP1. In this case, the terms involving $\tilde{\nu}_\tau$ are negligible, which is expected at the Z peak. In addition, the sneutrino contribution is proportional to $|V_{31}|^2$ which is small, as we mentioned in the previous section.

The magnitude of σ_Z depends crucially on the parameters $O'_{13}{}^L$ and $O'_{13}{}^R$ which are given by

$$\begin{aligned} O'_{13}{}^L &= -V_{11}V_{31}^* - \frac{1}{2}V_{12}V_{32}^* \\ O'_{13}{}^R &= -U_{11}^*U_{31} - \frac{1}{2}U_{12}^*U_{32} - \frac{1}{2}U_{13}^*U_{33} \end{aligned} \quad (15)$$

Since $O'_{13}{}^L$ depends on the R-parity violating V_{31} and V_{32} elements (which are small) one would expect the potentially larger contribution to come from $O'_{13}{}^R$, since the element U_{32} may have larger values. Nevertheless, it is possible to show that in the limit $m_{\nu_\tau} \rightarrow 0$ we get $U_{12}^*U_{32} - U_{13}^*U_{33} \rightarrow 0$. Therefore, the whole cross section is controlled by the neutrino mass. Numerically, we find that the elements V_{31} and V_{32} have a greater contribution to the cross-section.

At LEP1 despite the small values for the couplings $O'_{13}{}^L$ and $O'_{13}{}^R$ and the relatively low luminosity of 82 pb^{-1} , we benefit from being at the Z peak. A study of the cross-section at LEP1 was carried out in Ref. [11, 14] and it was found maximum cross-sections of order 1 pb, using lighter values for the chargino mass (45 GeV) and allowing a heavier tau neutrino (35 MeV). We improve that discussion by including a full scan of the parameter space as well as imposing the latest limits on the sparticle masses as well as m_{ν_τ} . The latest mass limits [see eq. (10)] restrict the parameter space for large cross-sections, so much so that the results of Ref. [14, 11] suggest that only cross-sections of order 0.1 pb are now possible, which would be unobservable at LEP1. Here we show that this was underestimated and explain the reasons.

Fig. 3 shows regions of attainable cross-sections in the $m_{\nu_\tau} - m_{\chi_1}$ plane. We consider $\tan\beta \leq 20$ and so is an update of the figure 7 in Ref. [11]. The reason why we get larger cross sections in this figure with respect to the previous reference is that we consider larger values of ϵ_3 . Inside the largest triangular region lie the points with $\sigma > 0.1$ pb, and inside the smaller triangular region we have $\sigma > 0.4$ pb, values which are in the limit of observability at LEP1. The vertical and horizontal sides of the triangles correspond to two of the limits in eq. (10).

In Fig. 4 we consider $\tan\beta$ up to 90 and this allows much larger cross-sections, for a given tau neutrino mass. The reason is that the cross section depends crucially on the tau Yukawa coupling: larger values of $\tan\beta$ increase the value of h_τ and this in turn increases σ . This region of large $\tan\beta$ was not explored in ref. [11]. In this figure we have three regions with the total cross section larger than 0.1, 0.4, and 4 pb.

To understand better the relation between the cross section and $\tan\beta$ we show in Fig. 5 the explicit cross-section dependence on $\tan\beta$. It shows clearly a steep climb for $\tan\beta \gtrsim 30$. In addition we show two curves, one for $m_{\nu_\tau} < 18$ MeV and another one with $m_{\nu_\tau} < 1$ MeV. They clearly show that the cross section is controlled by m_{ν_τ} and

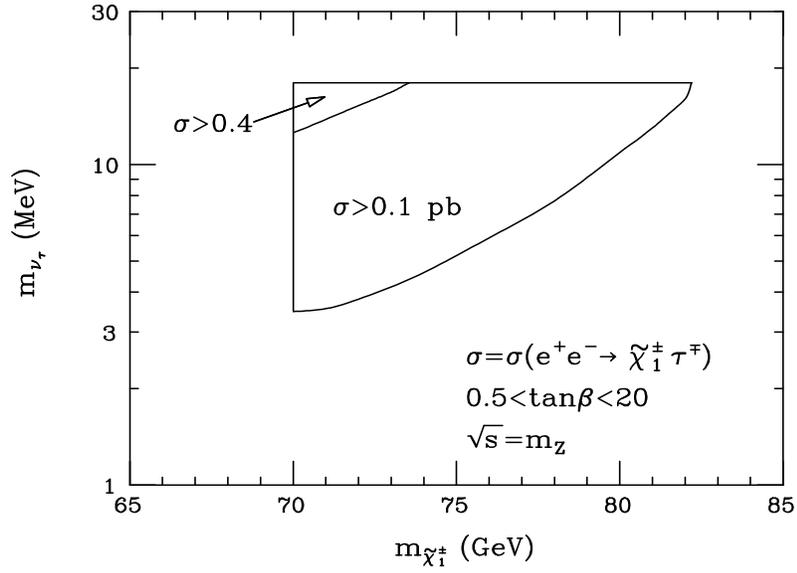


Figure 3: Regions of attainable cross section in BRpV in the plane tau neutrino mass ν/s chargino mass for moderate values of $\tan\beta$.

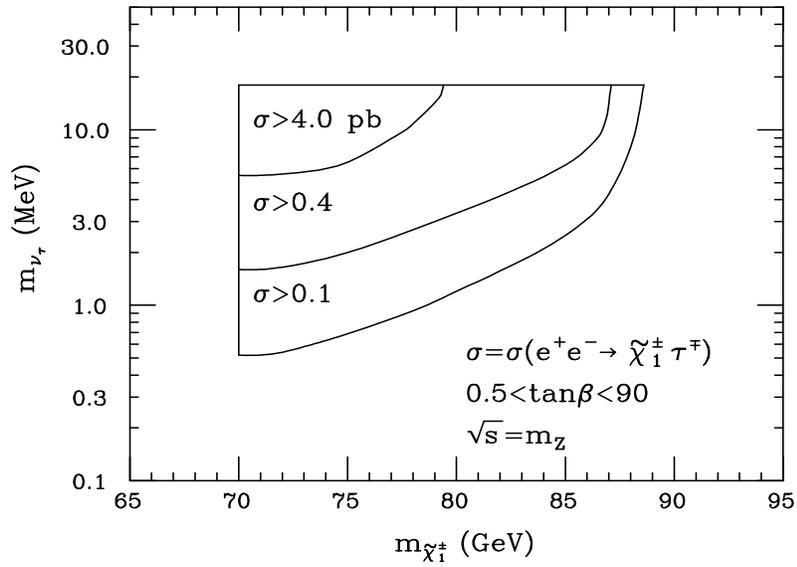


Figure 4: Regions of attainable cross section in BRpV in the plane tau neutrino mass ν/s chargino mass including large values of $\tan\beta$.

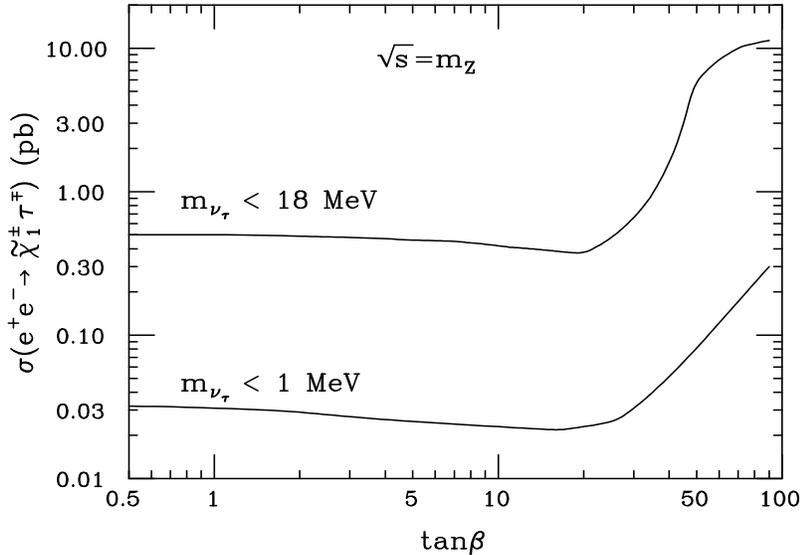


Figure 5: Maximum single chargino production cross section in BRpV as a function of $\tan\beta$ for two different upper bounds on the tau neutrino mass.

will approach zero as the neutrino mass goes to zero, as expected in the model, see, e.g. discussions given in [9].

Cross-sections as large 10 pb can be obtained for $\tan\beta \rightarrow 90$. Of interest is the region $55 \lesssim \tan\beta \lesssim 60$, which is required if top–bottom–tau Yukawa coupling unification is imposed [8]. Within this region we find maximum single chargino production cross section values of 8 pb. The cross-section has a direct dependence on the Tau Yukawa coupling and this can be explained as follows. As mentioned above, the coupling O_{13}^L gives the major contribution to the cross-section and O_{13}^L is a function of V_{31} and V_{32} , which in turn have a strong dependence on h_τ . One can infer from the matrix $M_C^T M_C$ that the largest values of h_τ require low v_3 and large $\tan\beta$, i.e., if the term $v_1^2 + v_3^2$ is made smaller then h_τ must be increased to keep the τ mass at 1.777 GeV. Since h_τ takes its largest values for $v_3 = 0$ and large $\tan\beta$ we can explain the $\tan\beta$ dependence found in the cross-section in Fig. 5. We note that in Ref. [11] v_3 was equated to zero, although here we explicit show that $v_3 = 0$ maximizes the cross-section. This can be seen explicitly in a graphic form as depicted in Fig. 6. This result has important implications if one considers the BRpV model as a low energy form of a spontaneous Rp violating model, in which a massless majoron is present. Such models require $v_3 \leq 100$ MeV. Our past discussion as well as Fig. 6 shows that it is precisely for such small v_3 values that we have maximum cross-sections. In addition the presence of a majoron would allow one to use the laboratory bound on m_{ν_τ} , which is less stringent than the nucleosynthesis bound, thus

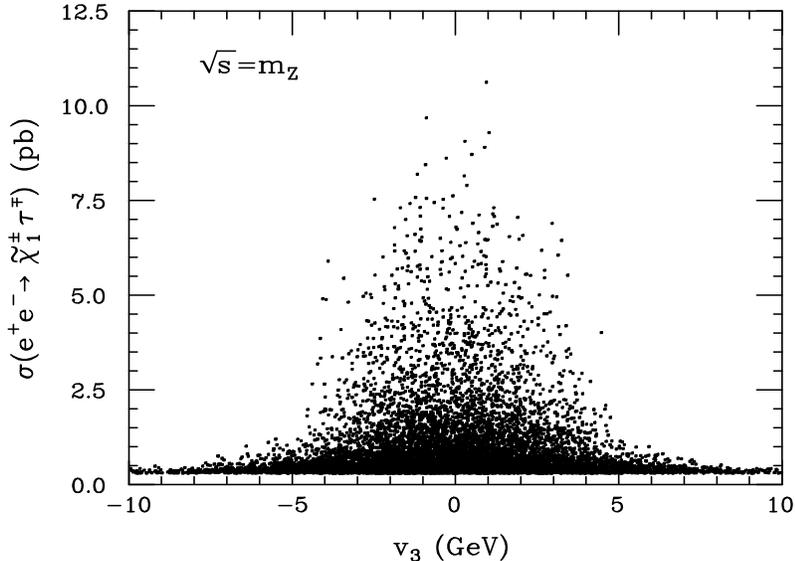


Figure 6: Maximum single chargino production cross section in BRpV as a function of v_3 .

allowing cross-sections of 1 to 10 pb in a limited range of parameter space. In the case of no majoron, one is obliged to use the tighter bound $m_{\nu_\tau} \lesssim 2$ MeV, which only allows maximum cross-sections of order 0.4 pb at LEP1.

At LEP2 one moves away from the Z peak and so the cross-section falls. With 4×10^4 random points chosen the maximum cross-section that we found had a value of 7.4 fb, and so 3.7 events would be expected at LEP2 with a luminosity of 500 pb^{-1} . Hence we conclude that LEP2 has no chance of obtaining a signal in this channel.

At an NLC of $\sqrt{s} = 500$ GeV one finds even smaller cross-sections although we benefit from the higher luminosity ($30 \rightarrow 100 \text{ fb}^{-1}$). Production of the heavier chargino χ_2 is now possible and we show in Fig. 7 the maximum values of the cross-section of both $\tilde{\chi}_1$ and $\tilde{\chi}_2$ as a function of chargino mass. We find that for $\tilde{\chi}_1$ the cross section starts at $m_{\chi_1} = 70$ GeV with 1.3 fb and then drops. In addition, the cross section for $\tilde{\chi}_2$ peaks at around $m_{\chi_2} = 250$ GeV with 1.2 fb and then falls. One would expect a smaller cross section for the heavier chargino since its mass is larger, but this is not the case. One can explain this as follows. For the cross-section of $\tilde{\chi}_1$ there exists a cancellation between the terms $\sigma_{\tilde{\nu}}$ and $\sigma_{Z\tilde{\nu}}$, which are of similar magnitude (each one is larger than σ_Z and with opposite relative sign). Hence one is left with the contribution σ_Z which gives the shape in Fig. 7. For $\tilde{\chi}_2$ the story is different. We find numerically that the coupling O'_{23} is smaller than O'_{13} . This implies that σ_Z and $\sigma_{Z\tilde{\nu}}$ are smaller in the case of $\tilde{\chi}_2$ compared with $\tilde{\chi}_1$. Therefore the shape of the cross-section in Fig. 7 for $\tilde{\chi}_2$ is dictated by the t-channel contribution $\sigma_{\tilde{\nu}}$.

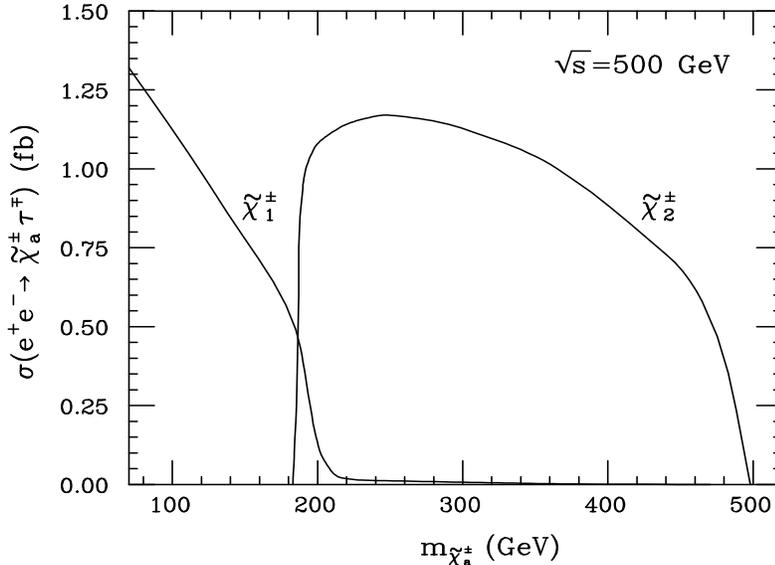


Figure 7: Maximum single chargino production cross section as a function of the chargino mass at NLC in BRpV. Light and heavy charginos are displayed.

5 Discussion and Conclusions

We have studied the charged and neutral current couplings of the tau lepton in models with spontaneous or bilinear breaking of R-parity. We showed that precision measurements of the $Z\tau\tau$ coupling at LEP1 allow relatively large values of the effective Rp breaking parameter ϵ_3 ($\rightarrow 200$ GeV) since such values only induce large mixing between τ^- and \tilde{H}_1^- which share the same $SU(2)$ quantum numbers. We then moved to see the possible manifestations of Rp breaking as single production of charginos in models with spontaneous or bilinear breaking of R-parity at various e^+e^- accelerators such as LEP1, LEP2, and a NLC with $\sqrt{s} = 500$ GeV. The numerical values of the cross-sections are small at LEP2 and at NLC, with maximum values of around 7.4 fb and 1.3 fb respectively. With large luminosity at the latter (say, 30 fb^{-1} or greater) there would be some chance of detection in this channel. On the contrary, the signal is un-observable at LEP2. At LEP1 prospects are vastly improved, and we found that the large values $\tan\beta$ (≈ 90) allow cross-sections as great as 10 pb. Assuming top-bottom-tau Yukawa unification (which requires $55 \lesssim \tan\beta \lesssim 60$) we find maximum cross-sections of around 8 pb.

The cross section decreases with decreasing m_{ν_τ} . In addition, there is a direct correlation between the tau Yukawa coupling (h_τ) and the cross-section, with large values of the latter only possible for large values of the former. The coupling h_τ itself has a strong dependence on v_3 and $\tan\beta$, with its greatest values being found at $v_3 = 0$ and

large $\tan\beta$.

Our results are to a large extent model independent, since they depend only of the structure of the chargino-tau mass matrix and this is universal in models with spontaneous breaking of R-parity as well as their effective truncation in terms of a bilinear explicit Rp violating superpotential term (BRpV model). The model dependence arises due to the following:

1. The bounds on m_{ν_τ} . This depends on whether or not the majoron exists. As already mentioned, if it does, we should apply the laboratory limit from LEP1, which allows for example, much larger single chargino production cross sections at the Z peak, due to the strict correlation between Rp violation observables and m_{ν_τ} . In models without the majoron, such as those in ref. [11, 12], or the BRpV model [5, 7] one must apply the tighter limits of about one MeV to the tau neutrino mass that follows from primordial Big-Bang nucleosynthesis and we loose the large cross sections.
2. The bounds on $m_{\tilde{\chi}_1}$. This also depends on whether or not the majoron exists. If it does, we should note that the chargino may be lighter than in the MSSM, the lower bound being about 70 GeV from LEP2 data [16], and therefore the left part in figures 3 and 4 is justified. If the majoron is absent, then we expect a tighter limit on the lighter chargino mass, similar to that which holds in the MSSM, though, strictly speaking, there has been no dedicated chargino search taking into account the decay modes expected in models with broken R-parity which would affect at least its cascade decays, since the lightest neutralino would decay.
3. The allowed values of v_3 . This also depends on whether or not the majoron exists. If it does, we should note that in order to avoid excessive stellar cooling by majoron emission one needs small v_3 . Fortunately, as seen in Fig. 6 our results for single production at LEP1 are better for small v_3 .

We conclude that the most favourable model which allows a sizeable single chargino cross section at the Z peak is the simplest $SU(2) \times U(1)$ spontaneously broken R-parity model, which is characterized by the existence of a majoron and a strict correlation between R-parity breaking observables and neutrino mass. Possible signatures will be model-dependent. For example, in the majoron version there will be di-tau + missing momentum, coming from the $e^+e^- \rightarrow \chi\tau$ with $\chi \rightarrow \tau + J$, which has been discussed in ref. [16] in the context of chargino pair-production at LEP2. In the explicit BRpV version the main effect of Rp violation is in the lightest neutralino decay which would lead to a rich cascade decay pattern characterized by large fermion multiplicities.

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References

- [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] G. Farrar and P. Fayet, *Phys. Lett.* **76B**, 575 (1978); *Phys. Lett.* **79B**, 442 (1978).
- [3] C.S. Aulakh, R.N. Mohapatra, *Phys.Lett.* **119B** 136, (1982); L. Hall and M. Suzuki, *Nucl.Phys.* **B231**, 419 (1984); G.G. Ross, J.W.F. Valle, *Phys.Lett.* **151B** 375 (1985); John Ellis, G. Gelmini, C. Jarlskog, G.G. Ross, J.W.F. Valle *Phys.Lett.* **150B** 142, (1985)
- [4] For a recent compilation of bounds see G. Bhattacharyya, hep-ph/9709395, *Nucl.Phys.Proc.Suppl.* 52A (1997) 83; bounds from $\mu - e$ conversion in nuclei are given in the recent paper by K. Huitu, J. Maalampi, M. Raidal, A. Santamaria, FTUV/97-45, HIP-1997-68/TH.
- [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.* **D51**, 2421 (1995); T. Banks, Y. Grossman, E. Nardi, and Y. Nir, *Phys. Rev. D* **52**, 5319 (1995); R. Hempfling, *Nucl. Phys.* **B478**, 3 (1996); F. Vissani and A. Yu. Smirnov, *Nucl.Phys.* **B460**, 37 (1996); H. P. Nilles and N. Polonsky, *Nucl. Phys.* **B484**, 33 (1997); B. de Carlos, P. L. White, *Phys.Rev.* **D55**, 4222 (1997); S. Roy and B. Mukhopadhyaya, *Phys. Rev. D* **55**, 7020 (1997); M. Carena, S. Pokorski, C. E. M. Wagner CERN-TH/97-373; M. Bisset, O.C.W. Kong, C. Macesanu, and L.H. Orr, hep-ph/9804282.
- [6] A. Akeroyd, M.A. Díaz, J. Ferrandis, M.A. Garcia-Jareño, and Jose W.F. Valle, hep-ph/9707395, *Nucl. Phys.* **Bxx**, xx (1998), in press.
- [7] M.A. Díaz, J.C. Romão, and J.W.F. Valle, hep-ph/9706315, *Nucl. Phys.* **Bxx**, xx (1998), in press
- [8] M.A. Díaz, J. Ferrandis, J.C. Romão, and J.W.F. Valle, hep-ph/9801391.
- [9] J.W.F. Valle, proceedings of the *Workshop on Physics Beyond the Standard Model: Beyond the Desert: Accelerator and Nonaccelerator Approaches*, Tegernsee, Germany, 8-14 Jun 1997, hep-ph/9712277.

- [10] M.A. Díaz, talks given at the International Workshop on Quantum Effects in the MSSM, Barcelona, Spain, Nov. 1997, hep-ph/9711435, and at International Europhysics Conference on High-Energy Physics, Jerusalem, Israel, 19-26 Aug 1997, hep-ph/9712213
- [11] M.C. Gonzalez-Garcia and J.W.F. Valle, *Nucl. Phys. B* **355**, 330 (1991).
- [12] K. Huitu, J. Maalampi, K. Puolamaki, e-Print Archive: hep-ph/9705406
- [13] A. Masiero and J.W.F. Valle, *Phys. Lett.* **B251**, 273 (1990); J.C. Romão, C.A. Santos, J.W.F. Valle, *Phys. Lett. B* **288**, 311 (1992); J.C. Romão, A. Ioannissyan and J.W.F. Valle, *Phys. Rev. D***55**, 427 (1997).
- [14] P. Nogueira, J.C. Romao, J.W.F. Valle, *Phys.Lett.* **B251** 142-14 (1990)
- [15] Review of Particle Physics, *Phys. Rev. D* **54**, 1 (1996).
- [16] F. de Campos, O.J.P. Eboli, M.A. Garcia-Jareno, J.W.F. Valle hep-ph/9710545, submitted to Nucl. Phys.
- [17] ALEPH collaboration, *European Physical Journal* **C2** (1998) 3, 395.
- [18] J. C. Romao, F. de Campos, M. A. Garcia-Jareno, M. B. Magro, J. W. F. Valle, *Nucl.Phys.* **B482** 3-23 (1996)
- [19] A. Bartl, H. Fraas, and W. Majerotto, *Z. Phys.* **C30**, 441 (1986).