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Spontaneous R-parity breaking at hadron supercolliders *

M.C. Gonzalez-Garcia [1]

Physics Department, University of Wisconsin, Madison, WI 53706, USA

J.C. Romão [2]

Centro de Física da Matéria Condensada, INIC, Av. Prof. Gama Pinto, 2, 1699 Lisbon, Codex, Portugal

J.W.F. Valle [3]

Instituto de Física Corpuscular, CSIC, Departament de Física Teòrica, Universitat de València, 46100 Burjassot, València, Spain

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If supersymmetry (SUSY) is realized with spontaneous breaking of R-parity, there can be large rates for single sparticle production. We make a detailed study of the rates for Drell–Yan-like single chargino and single neutralino production at a hadron supercollider. The attainable rates are promising in view of the luminosities expected at these facilities. All observational restrictions from cosmology, astrophysics and laboratory are taken into account in our analysis, including the recent LEP results as well as neutrino physics constraints.

1. Introduction

Most studies of supersymmetric phenomenology have been made in the framework of the Minimal Supersymmetric Standard Model (MSSM) which assumes the conservation of a discrete symmetry called R-parity [1]. Under this symmetry all the standard model particles are R-even while their superpartners are R-odd. R-parity is related to the spin (S), total lepton (L), and baryon (B) number

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[1] E-mail: CONCHA@WISCPHEN or 47375::CONCHA

[2] E-mail: ROMAO@PTIFM

[3] E-mail: VALLE@EVALUN11 or 16444::VALLE

according to $R_p = (-1)^{(3B+L+2S)}$. Therefore the requirement of baryon and lepton number conservation implies the conservation of R-parity. Under this assumption the supersymmetric (SUSY) particles must be pair-produced, the lightest of them being absolutely stable. These two features underlie all the experimental searches for new supersymmetric states. In addition, the assumption of R-parity conservation may have important cosmological implications. e.g. related to the question of dark matter.

However, neither gauge invariance nor SUSY require R-parity conservation. Moreover, whether or not R-parity is a good symmetry, and to what extent, is ultimately a dynamical question, related to physics at a more fundamental scale. For example, R-parity violation could emerge as the residual effect of a more unified theory scheme or, alternatively, it could be generated in a spontaneous way. At the present state-of-the-art theory can not decide. It is therefore of great interest to pursue the phenomenological consequences of alternative scenarios. Specially in view of the fact that the associated effects may well be accessible to experimental verification.

The most general supersymmetric extension of the standard model contains explicit R-parity violating interactions [2]

$$W = \lambda^{(1)} u^c u^c d^c + \lambda^{(2)} l l e^c + \lambda^{(3)} Q l d^c. \tag{1}$$

These are consistent with both gauge invariance and supersymmetry. These R-parity violating interactions involve a large number of arbitrary constants, since each of the $\lambda^{(i)}$ has three generation indices, which have been omitted. Detailed analysis of the constraints on these models and their possible signals have been made in refs. [3,4]. Many of these couplings must be set to zero in order, for example, to avoid having too fast proton decay rates. Restricted structures of the generation indices may, however, be phenomenologically acceptable. These could arise by imposing some global and/or discrete symmetry. In addition to proton decay constraints there are, however, other stringent restrictions on R-parity violating couplings. These may affect in fact, *all* of these couplings, if all of them are present, in the most general form. These follow from cosmological arguments related to the baryon asymmetry of the universe [5]. The point is that the interactions in eq. (1) mediate $(B-L)$-violating decays of squarks and sleptons of the kind $\tilde{\mathrm{u}} \to \bar{\mathrm{d}}\bar{\mathrm{d}}$, $\tilde{\mathrm{u}} \to \bar{\ell}\mathrm{d}$, $\tilde{\ell} \to \ell\nu$, etc. The rate for these decay processes may be estimated as

$$\Gamma_{\mathrm{D}} \approx \frac{\lambda^2}{4\pi} \frac{\tilde{m}^2}{\left(\tilde{m}^2 + T^2\right)^{1/2}}, \tag{2}$$

where λ denotes the appropriate coupling constant in eq. (1) and $\tilde{m}$ the corresponding supersymmetric particle mass relevant for the decay process of interest.

Now, it has been realized [6] that at temperatures T above $\sim m_W/\alpha_{\text{weak}}$, transitions that violate B and L number will occur rapidly since these quantum numbers are not conserved due to the electroweak anomaly. These transitions may then erase any primordial B-asymmetry, which could have been generated at some grand-unification energy scale. In order that a pre-existing B-asymmetry survive, an excess of the anomaly-free $B-L$ combination must have existed at very early times, and this may be easily generated in most GUTs. However, in this case it is crucial that the $B-L$ asymmetry not be eliminated through some other mechanism, such as as the interactions present in eq. (1). Requiring that no such interaction comes into equilibrium after the $B-L$ asymmetry is produced, i.e. that Γ_D is smaller than the Hubble constant at any $T > \tilde{m}$, leads to [5]

$$\lambda \leqslant \mathscr{O}(10^{-7}) \left(\frac{\tilde{m}}{\text{TeV}} \right)^{1/2}. \tag{3}$$

Barring the imposition of additional symmetries that may restrict the flavour structure of the R_p violating couplings [7] this bound holds for each of the coupling constants of the operators in eq. (1), and rules out the possibility of observing new signals associated to *explicitly* R-parity violating interactions in collider experiments. However, we caution the reader that, although this argument has been used in the literature as stated above, it is not really inescapable. For example, explicit R_p violating interactions could be tolerated in the presence of a mechanism that could generate a nonzero baryon asymmetry at low energy, as suggested in ref. [8].

It seems to us more attractive to focus our attention to the possibility that R-parity can be an exact symmetry of the lagrangian, broken spontaneously through the Higgs mechanism [9–12]. This may occur via nonzero vacuum expectation values (VEVs) for scalar neutrinos, such as

$$v_R = \langle \tilde{\nu}_{R\tau} \rangle, \qquad v_L = \langle \tilde{\nu}_{L\tau} \rangle. \tag{4}$$

If lepton number is part of the gauge symmetry there is an additional Z′ gauge boson which acquires mass via the Higgs mechanism. In this case there is no physical Goldstone boson and the scale of R-parity violation is the same that characterizes the new gauge interaction, around TeV scale. Consequently, its effects can be large [12].

On the other hand, if spontaneous R-parity violation occurs in absence of any additional gauge symmetry, it leads to the existence of a physical massless Nambu–Goldstone boson, called majoron (J) [9,13]. Thus, *the lightest SUSY particle is the majoron*. In these models there is a new decay mode for the Z:

$$Z \to \rho + J, \tag{5}$$

where ρ is a light scalar of mass of order $v \ll M_{\mathrm{W}}$ where v is the scale of the spontaneous R-parity breaking. This decay mode would increase the invisible width of the Z by an amount of at least $\Delta\Gamma^{\mathrm{inv}} = 85$ MeV. The LEP measurement of the Z width [14] is enough to exclude any model where the majoron is not mainly an isosinglet [15]. The simplest way to avoid this limit is to extend the minimal supersymmetric model, so that the R-parity breaking is driven by isosinglet lepton VEVs, so that the majoron is mainly singlet [10]. This requires the existence of right-handed neutrinos.

In this paper we analyse the phenomenology of spontaneously broken R-parity at hadron colliders such as LHC, and SSC. The new single-production and decay mechanisms for the supersymmetric particles are studied in detail. We demonstrate that, for integrated luminosities in the range from $L \sim 10^4$ pb^{-1} to $L \sim 10^5$ pb^{-1} in a year, the rates for single production of the lightest SUSY fermions can be sizeable at the planned LHC and SSC hadron colliders, without conflicting any laboratory, cosmological or astrophysical observation. For example at LHC it is possible to have more than 10 events for neutralino masses in the range $M_{\chi^0} \lesssim$ 160–280 GeV and chargino mass $M_{\chi^+} \lesssim$ 180–320 GeV. The lower masses correspond to $L = 10^4$ pb^{-1} and the upper to $L = 10^5$ pb^{-1}. The corresponding masses that can be explored at SSC are in the range $M_{\chi^0} \lesssim$ 200–360 GeV to $M_{\chi^+} \lesssim$ 220–400 GeV. We have also analysed the possible decay channels of the lowest lying SUSY fermions in spontaneously broken R-parity and identified the corresponding signals.

2. A model for spontaneous *R*-parity breaking [10]

In order to set up our notation we recall the basic ingredients of the model for spontaneous violation of R-parity and lepton number proposed in ref. [10]. The superpotential is given by

$$h_{\mathrm{u}} Q H_{\mathrm{u}} u^{\mathrm{c}} + h_{\mathrm{d}} H_{\mathrm{d}} Q d^{\mathrm{c}} + h_{\mathrm{e}} \ell H_{\mathrm{d}} e^{\mathrm{c}} + \left(h_0 H_{\mathrm{u}} H_{\mathrm{d}} - \epsilon^2\right)\Phi + h_\nu \ell H_{\mathrm{u}} \nu^{\mathrm{c}} + h \Phi S \nu^{\mathrm{c}} + \mathrm{h.c.} \tag{6}$$

This superpotential conserves *total* lepton number and R-parity. The superfields (Φ, ν_i^{c}, S_i) are singlets under $\mathrm{SU}_2 \otimes \mathrm{U}(1)$ and carry a conserved lepton number assigned as (0, -1, 1) respectively. All couplings h_{u}, h_{d}, h_{e}, h_ν, h are described by arbitrary matrices in generation space which explicitly break flavor conservation. The additional singlets may appear in different theoretical frameworks [16,17], and some of their possible phenomenological consequences were discussed in refs. [18–20].

As was shown in refs. [10,21] these singlets may drive the spontaneous violation of R-parity leading to the existence of a majoron given by the imaginary part of

$$\frac{v_L^2}{Vv^2}(v_u H_u - v_d H_d) + \frac{v_L}{V}\tilde{\nu}_\tau - \frac{v_R}{V}\tilde{\nu}_\tau^c + \frac{v_S}{V}\tilde{S}_\tau \tag{7}$$

where the isosinglet VEVs

$$v_R = \langle \tilde{\nu}_{R\tau} \rangle, \qquad v_S = \langle \tilde{S}_\tau \rangle \tag{8}$$

with $V = \sqrt{v_R^2 + v_S^2}$ characterize R-parity or lepton number breaking and the isodoublet VEVs

$$v_u = \langle H_u \rangle, \qquad v_d = \langle H_d \rangle \tag{9}$$

drive electroweak breaking and the fermion masses. The combination $v^2 = v_u^2 + v_d^2$ is fixed by the W-mass. Finally there is a small seed of R-parity breaking in the doublet sector, i.e.

$$v_L = \langle \tilde{\nu}_{L\tau} \rangle \tag{10}$$

whose magnitude is now related to the Yukawa coupling h_ν. Since this vanishes as $h_\nu \to 0$, we can naturally obey the limits from stellar energy loss [22], a possibility which is not available in the minimal model of ref. [9].

Several phenomenological aspects of this model relevant for LEP [11], solar neutrinos [24], as well as intense μ and τ sources [25], such as expected at PSI and/or at a tau factory [23] were previously studied. Here we are concerned with the effects at a hadron supercollider. For phenomenological studies in spontaneously broken R-parity we need the forms of the relevant *chargino* and *neutralino* mass matrices. The *chargino* mass matrix is given by

$$\begin{array}{c|ccc} & e_j^+ & \tilde{H}_u^+ & -i\tilde{W}^+ \\ \hline e_i & h_{eij}v_d & -h_{\nu ij}v_{Rj} & \sqrt{2}g_2 v_{Li} \\ \tilde{H}_d^- & -h_{eij}v_{Li} & \mu & \sqrt{2}g_2 v_d \\ -i\tilde{W}^- & 0 & \sqrt{2}g_2 v_u & M_2 \end{array} \tag{11}$$

In some models, such as the one in ref. [10], the effective Higgsino mixing parameter μ may be given as $\mu = h_0\langle \Phi \rangle$, where $\langle \Phi \rangle$ is the VEV of an appropriate singlet scalar. The form of this matrix is common to a wide class of SU(2) ⊗ U(1) SUSY models with spontaneously broken R-parity. In contrast, the mass matrix for the neutral leptons is more model dependent. In the case of interest, it is specified by a 14×14 mass matrix. Under reasonable approximations, we can remove away all heavy isosinglet leptons that may be present. This approximate diagonalization leads to an effective 7×7 *neutralino* mass matrix of the following form [10]:

$$
\begin{array}{c|ccccc}
 & \nu_i & \tilde{H}_{\mathrm{u}} & \tilde{H}_{\mathrm{d}} & -i\tilde{W}_3 & -i\tilde{B} \\
\hline
\nu_i & 0 & h_{\nu ij} v_{\mathrm{R}j} & 0 & g_2 v_{\mathrm{L}i} & -g_1 v_{\mathrm{L}i} \\
\tilde{H}_{\mathrm{u}} & h_{\nu ij} v_{\mathrm{R}j} & 0 & -\mu & -g_2 v_{\mathrm{u}} & g_1 v_{\mathrm{u}} \\
\tilde{H}_{\mathrm{d}} & 0 & -\mu & 0 & g_2 v_{\mathrm{d}} & -g_1 v_{\mathrm{d}} \\
-i\tilde{W}_3 & g_2 v_{\mathrm{L}i} & -g_2 v_{\mathrm{u}} & g_2 v_{\mathrm{d}} & M_2 & 0 \\
-i\tilde{B} & -g_1 v_{\mathrm{L}i} & g_1 v_{\mathrm{u}} & -g_1 v_{\mathrm{d}} & 0 & M_1
\end{array} \tag{12}
$$

In the above two equations $M_{1,2}$ denote the supersymmetry breaking gaugino mass parameters and $g_{1,2}$ are the $\mathrm{SU(2)} \otimes \mathrm{U(1)}$ gauge couplings divided by $\sqrt{2}$. We assume the canonical relation $M_1/M_2 = \frac{5}{3}\tan^2\theta_{\mathrm{W}}$.

Two matrices U and V are needed to diagonalize the 5×5 (non-symmetric) chargino mass matrix

$$\chi_i^+ = V_{ij}\psi_j^+, \qquad \chi_i^- = U_{ij}\psi_j^- \qquad (13),(14)$$

where the indices i and j run from 1 to 5 and $\psi_j^+ = (e_1^+, e_2^+, e_3^+, \tilde{H}_{\mathrm{u}}^+, -i\tilde{W}^+)$ and $\psi_j^- = (e_1^-, e_2^-, E_3^-, \tilde{H}_{\mathrm{d}}^-, -i\tilde{W}^-)$.

Similarly, the neutralino mass matrix (symmetric, due to the Pauli exclusion principle) is diagonalized by a single 7×7 matrix N, i.e.

$$\chi_i^0 = N_{ij}\psi_j^0 \tag{15}$$

where $\psi_j^0 = (\nu_i, \tilde{H}_{\mathrm{u}}, \tilde{H}_{\mathrm{d}}, -i\tilde{W}_3, -i\tilde{B})$, with ν_i denoting weak-eigenstate neutrinos. Here the indices i and j run from 1 to 7.

From the diagonalization of the previous matrices in the range of parameters given in eqs. (30) and (31), and after imposing the constraints in sect. 4, it follows that

(i) the lightest supersymmetric fermion is *always* a neutralino;

(ii) the mass of the lightest neutralino lies in the range 25 GeV $\lesssim M_{\chi^0} \lesssim$ 500 GeV. The lower bound is determined by the experimental limit on the chargino mass;

(iii) the second lightest supersymmetric fermion is most of the times the lightest chargino.

In addition to the determination of the lepton and SUSY fermion mass spectrum, the diagonalization of these matrices also leads to the mixing of the standard leptons with the supersymmetric ones, which is the origin of R-parity violation in these models. In what follows we will give explicit expressions for the couplings of the SUSY particles in terms of these diagonalizing matrices.

3. *R*-parity breaking couplings

Unlike the models with explicit R-parity breaking in the Yukawa couplings of the superpotential, in models with spontaneously broken R-parity the R-violating couplings appear when writing all interactions in the lagrangian in terms of the mass eigenstates. The mixing of the standard leptons with the supersymmetric ones is what violates R-parity in these models. R-parity violation occurs mostly in the gauge couplings, with subdominant R-parity breaking effects also occurring in the Yukawa couplings.

3.1. GAUGE COUPLINGS

Using the diagonalizing matrices one can write the electroweak currents in terms of mass-eigenstate fermions. For example, the charged current lagrangian describing the charged lepton/neutral lepton weak interaction may be written as

$$\frac{g}{\sqrt{2}} W_\mu \bar{\chi}_i^- \gamma^\mu (K_{\mathrm{L}ik} P_\mathrm{L} + K_{\mathrm{R}ik} P_\mathrm{R}) \chi_k^0 + \mathrm{h.c.}, \tag{16}$$

where $P_{\mathrm{L,R}}$ are the two chiral projectors and the 5×7 coupling matrices $K_{\mathrm{L,R}}$ may be written as

$$K_{\mathrm{L}ik} = \eta_i \left(-\sqrt{2}\, U_{i5} N_{k6} - U_{i4} N_{k5} - \sum_{m=1}^{3} U_{im} N_{km} \right), \tag{17}$$

$$K_{\mathrm{R}ik} = \epsilon_k \left(-\sqrt{2}\, V_{i5} N_{k6} + V_{i4} N_{k4} \right). \tag{18}$$

The matrix $K_{\mathrm{L}ik}$ is the analogous of the matrix K introduced in ref. [26]. These couplings break R-parity for $i = 1, \ldots, 3$ and $k = 4, \ldots, 7$, and $i = 4$, 5 and $k = 1, \ldots, 3$.

Similarly, the neutral current lagrangian describing the charged lepton/charged lepton and neutral lepton/neutral lepton weak interaction may be written as

$$\frac{g}{\cos\theta_\mathrm{W}} Z_\mu \{ \bar{\chi}_i^- \gamma^\mu (\eta_i \eta_k O'_{\mathrm{L}ik} P_\mathrm{L} + O'_{\mathrm{R}ik} P_\mathrm{R}) \chi_k^- + \tfrac{1}{2} \bar{\chi}_i^0 \gamma^\mu (\epsilon_i \epsilon_k O''_{\mathrm{L}ik} P_\mathrm{L} + O''_{\mathrm{R}ik} P_\mathrm{R}) \chi_k^0 \}, \tag{19}$$

where the 7×7 coupling matrices $O'_{\mathrm{L,R}}$ and $O''_{\mathrm{L,R}}$ are given by

$$O'_{\mathrm{L}ik} = \tfrac{1}{2} U_{i4} U_{k4} + U_{i5} U_{k5} + \tfrac{1}{2} \sum_{m=1}^{3} U_{im} U_{km} - \delta_{ik} \sin^2\theta_\mathrm{W}, \tag{20}$$

$$O'_{\mathrm{R}ik} = \tfrac{1}{2} V_{i4} V_{k4} + V_{i5} V_{k5} - \delta_{ik} \sin^2\theta_{\mathrm{W}}, \tag{21}$$

$$O''_{\mathrm{L}ik} = \tfrac{1}{2} \left\{ N_{i4} N_{k4} - N_{i5} N_{k5} - \sum_{m=1}^{3} N_{im} N_{km} \right\} = -O''_{\mathrm{R}ik}. \tag{22}$$

In writing these couplings we have assumed CP conservation. Under this assumption the diagonalizing matrices can be chosen to be real. The η_i and ϵ_k factors are sign factors, related with the relative CP parities of these fermions, that follows from the diagonalization of their mass matrices. These couplings break R-parity for $i = 1,\dots,3$ and $k = 4, 5$ in the case of the charged leptons, and $i = 4,\dots,7$ and $k = 1,\dots,3$ for the neutral leptons.

The diagonal couplings for the lightest neutralino and the lightest chargino are of the same order as those in the MSSM. The coupling of the lightest chargino to the Z is maximum when it is mainly a gaugino. In this case $\mu \gg M_2$ and therefore $|V_{45}| \approx |U_{45}| \approx 1$ and $|V_{4i}|$, $|U_{4i}| \ll 1$ for $i \neq 5$. On the other hand it is minimum when its larger component is along the higgsino. In this case $\mu \ll M_2$ and therefore $|V_{44}| \approx |U_{44}| \approx 1$ and $|V_{4i}|$, $|U_{4i}| \ll 1$ for $i \neq 4$. Including these values in eqs. (20) and (21) one gets for the allowed range

$$0.27 \lesssim |O'_{\mathrm{L}44}|,\ |O'_{\mathrm{R}44}| \lesssim 0.77. \tag{23}$$

The R-parity conserving W-coupling of the lightest chargino to the lightest neutralino is maximal when the chargino consists mainly of a higgsino ($|V_{44}| \approx |U_{44}| \approx 1$ and $|V_{4i}|$, $|U_{4i}| \ll 1$ for $i \neq 4$). In that case there are two light neutralinos which form a quasi-Dirac state [27] ($|N_{44}| \approx |N_{45}| \approx 1/\sqrt{2}$). Including these values of the diagonalizing matrices in eqs. (17) and (18) one gets the upper bound on these couplings. Experimental constraints (see sect. 4) determine the lower limit. With all that we have the allowed range

$$10^{-4} \lesssim |K_{\mathrm{L}44}|,\ |K_{\mathrm{R}44}| \lesssim \frac{1}{\sqrt{2}}. \tag{24}$$

For the neutral current couplings of the lightest neutralino the situation is not so clear, because they vanish in both limits $\mu \gg M_2$ and $\mu \ll M_2$. In the first case since the neutralino is mainly a gaugino it does not couple to the Z. In the second case the lightest neutralino is a quasi-Dirac fermion and the couplings cancel almost exactly. After imposing the experimental constraints detailed in sect. 4 we get

$$|O''_{\mathrm{L}44}| \lesssim 0.1. \tag{25}$$

In what concerns the R-parity breaking couplings, the biggest ones correspond to the standard lepton belonging to the third family. In fig. 1 we have plotted the

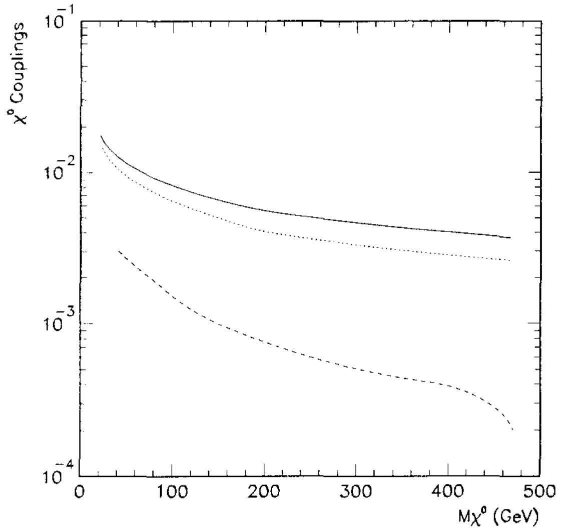

Fig. 1. Attainable R-parity breaking coupling strengths for the lightest neutralino versus its mass: $|K_{L34}|$ (full line), $|K_{R34}|$ (dash line) and $|O''_{L34}|$ (dotted line).

value of these couplings for the lightest neutralino as a function of its mass. In ref. [11] it was noted that this neutral current coupling (dotted line) is responsible for interesting rare zen-events observable at LEP. In fig. 2 we have plotted the corresponding couplings for the lightest chargino. In ref. [11] it was noted that

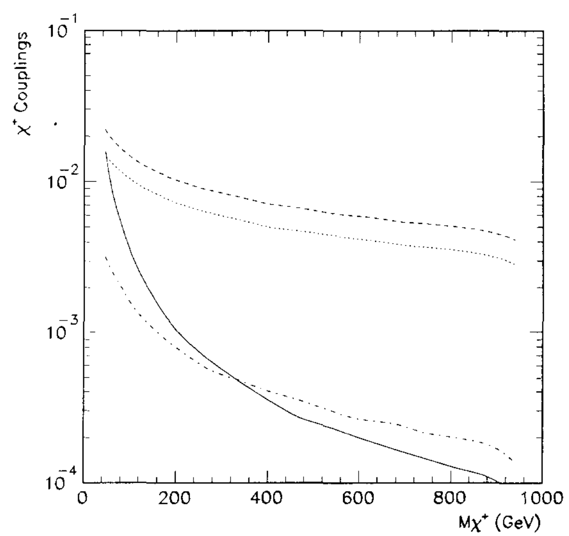

Fig. 2. Attainable R-parity breaking coupling strengths for the lightest chargino versus its mass: $|K_{L43}|$ (full line), $|K_{R43}|$ (dash line), $|O'_{L34}|$ (dotted line) and $|O'_{LR4}|$ (dash–dotted line).

these neutral current couplings (dotted and dash–dotted line) are responsible for inducing the decay $Z \to \tau + \chi$ with rates observable at LEP. As we will see the above couplings can also give rise to observable single production of SUSY particles in hadron supercolliders.

In our estimate of the attainable couplings we have taken into account all of the experimental constraints (see below).

3.2. YUKAWA COUPLINGS

The superpotential (6) includes the following set of Yukawa couplings terms involving two fermions:

$$
\begin{aligned}
& h_{\mathrm{u}ij}\left[\tilde{u}_i\tilde{H}_{\mathrm{u}}^0 u_j^{\mathrm{c}} + u_i\tilde{H}_{\mathrm{u}}^0\tilde{u}_j^{\mathrm{c}} - \tilde{d}_i\tilde{H}_{\mathrm{u}}^+ u_j^{\mathrm{c}} - d_i\tilde{H}_{\mathrm{u}}^+\tilde{u}_j^{\mathrm{c}}\right] \\
& + h_{\mathrm{d}ij}\left[\tilde{H}_{\mathrm{d}}^0\tilde{d}_i d_j^{\mathrm{c}} + \tilde{H}_{\mathrm{d}}^0 d_i\tilde{d}_j^{\mathrm{c}} - \tilde{H}_{\mathrm{d}}^-\tilde{u}_i d_j^{\mathrm{c}} - \tilde{H}_{\mathrm{d}}^- u_i\tilde{d}_j^{\mathrm{c}}\right] \\
& + h_{\mathrm{e}ij}\left[\tilde{H}_{\mathrm{d}}^0\tilde{e}_i e_j^{\mathrm{c}} + \tilde{H}_{\mathrm{d}}^0 e_i\tilde{e}_j^{\mathrm{c}} - \tilde{H}_{\mathrm{d}}^-\tilde{\nu}_i e_j^{\mathrm{c}} - \tilde{H}_{\mathrm{d}}^-\nu_i\tilde{e}_j^{\mathrm{c}}\right] \\
& + h_{\nu ij}\left[\nu_i\tilde{H}_{\mathrm{u}}^0\tilde{\nu}_j^{\mathrm{c}} - e_i^+\tilde{H}_{\mathrm{u}}^+\tilde{\nu}_j^{\mathrm{c}}\right]. \qquad (26)
\end{aligned}
$$

We now write the weak eigenstates fermions in terms of the mass eigenstates. We will assume that fermions and sfermions are simultaneously diagonalizable. In this approximation we can take

$$
\begin{aligned}
h_{\mathrm{e}ij} &= \frac{m_{\mathrm{e}i}}{v_{\mathrm{d}}\left(1 - Ah_{\nu i3}^2 v_{\mathrm{R}}^2\right)}\delta_{ij}, \\
h_{\mathrm{u}ij} &= \frac{m_{\mathrm{u}i}}{v_{\mathrm{u}}}\delta_{ij}, \\
h_{\mathrm{d}ij} &= \frac{m_{\mathrm{d}i}}{v_{\mathrm{d}}}\delta_{ij}, \qquad A = \frac{g^2 v_{\mathrm{d}}^2 + M_2^2}{2\left(v_{\mathrm{u}}v_{\mathrm{d}}g^2 - \mu M_2\right)^2} \qquad (27)
\end{aligned}
$$

for $i = 1,\ldots,3$. With $d_i^m = C_{ik}d_k^w$, where C is the CKM matrix. We have then the R-parity breaking couplings

$$
\begin{aligned}
& \lambda_{1ijk}^{\mathrm{u}}\left(\nu_i u_j^{\mathrm{c}}\tilde{u}_k + \nu_i u_j\tilde{u}_k^{\mathrm{c}}\right) - \lambda_{2ijk}^{\mathrm{u}} e_i^+ u_j^{\mathrm{c}}\tilde{d}_k - \lambda_{2ikj}^{\mathrm{u}} e_i^+ d_j\tilde{u}_k^{\mathrm{c}} \\
& + \lambda_{1ijk}^{\mathrm{d}}\left(\nu_i d_j^{\mathrm{c}}\tilde{d}_k + \nu_i d_j\tilde{d}_k^{\mathrm{c}}\right) - \lambda_{2ikj}^{\mathrm{d}} e_i^- d_j^{\mathrm{c}}\tilde{u}_k - \lambda_{2ijk}^{\mathrm{d}} e_i^- u_j\tilde{d}_k^{\mathrm{c}} \\
& + \lambda_{1ijk}^{\mathrm{e}}\nu_i e_j^+\tilde{e}_k + \lambda_{2ijk}^{\mathrm{e}}\nu_i e_j^-\tilde{e}_k^{\mathrm{c}} - \lambda_{3ijk}^{\mathrm{e}} e_i^- e_j^+\tilde{\nu}_k - \lambda_{1ijk}^{\nu}\nu_i\nu_j\tilde{\nu}_k^{\mathrm{c}} + \lambda_{2ijk}^{\nu} e_i^- e_j^+\tilde{\nu}_k^{\mathrm{c}} \qquad (28)
\end{aligned}
$$

TABLE 1
Upper limits for the R-parity breaking Yukawa couplings in the quark sector

I	J	K	λ^{u}_{1ijk}	λ^{u}_{2ijk}	λ^{d}_{1ijk}	λ^{d}_{2ijk}
1	1	1	2.5×10^{-21}	1.5×10^{-14}	1.2×10^{-14}	1.4×10^{-8}
1	1	2	0	3.4×10^{-15}	0	3.2×10^{-9}
1	1	3	0	1.7×10^{-16}	0	5.8×10^{-11}
1	2	1	0	8.1×10^{-13}	0	6.4×10^{-8}
1	2	2	$6\ \times10^{-19}$	3.6×10^{-12}	2.5×10^{-13}	2.8×10^{-7}
1	2	3	0	1.6×10^{-13}	0	1.3×10^{-8}
1	3	1	0	1.6×10^{-12}	0	$8\ \times10^{-8}$
1	3	2	0	1.8×10^{-11}	0	3.2×10^{-7}
1	3	3	6.6×10^{-17}	4.1×10^{-10}	6.2×10^{-12}	7.3×10^{-6}
2	1	1	1.5×10^{-12}	5.4×10^{-10}	3.3×10^{-7}	2.1×10^{-6}
2	1	2	0	1.2×10^{-10}	0	4.8×10^{-7}
2	1	3	0	6.1×10^{-12}	0	8.7×10^{-9}
2	2	1	0	2.9×10^{-8}	0	9.7×10^{-6}
2	2	2	3.7×10^{-10}	1.3×10^{-7}	6.6×10^{-6}	4.3×10^{-5}
2	2	3	0	5.9×10^{-9}	0	1.9×10^{-6}
2	3	1	0	5.9×10^{-8}	0	1.2×10^{-5}
2	3	2	0	6.5×10^{-7}	0	4.8×10^{-5}
2	3	3	4.1×10^{-8}	1.5×10^{-5}	1.6×10^{-4}	1.1×10^{-3}
3	1	1	1.8×10^{-7}	1.7×10^{-7}	3.3×10^{-5}	3.2×10^{-5}
3	1	2	0	3.9×10^{-8}	0	7.2×10^{-6}
3	1	3	0	$2\ \times10^{-9}$	0	1.3×10^{-7}
3	2	1	0	9.5×10^{-6}	0	1.5×10^{-4}
3	2	2	4.4×10^{-5}	4.2×10^{-5}	6.6×10^{-4}	2.9×10^{-5}
3	2	3	0	1.9×10^{-6}	0	1.3×10^{-7}
3	3	1	0	1.9×10^{-5}	0	1.8×10^{-4}
3	3	2	0	2.1×10^{-4}	0	7.2×10^{-4}
3	3	3	4.9×10^{-3}	4.8×10^{-3}	1.7×10^{-2}	1.7×10^{-2}

with

$$\lambda^{u}_{1ijk}=h_{uk}N^*_{i4}\delta_{jk}, \qquad \lambda^{u}_{2ijk}=h_{uj}V^*_{i4}C^*_{kj},$$

$$\lambda^{d}_{1ijk}=h_{dk}N^*_{i5}\delta_{jk}, \qquad \lambda^{d}_{2ijk}=h_{dj}U^*_{i4}C_{jk},$$

$$\lambda^{\nu}_{1ijk}=\sum_{i=1}^{3}h_{\nu mk}N^*_{im}N^*_{j4}, \qquad \lambda^{\nu}_{2ijk}=\sum_{i=1}^{3}h_{\nu mk}U^*_{im}V^*_{j4},$$

$$\lambda^{e}_{1ijk}=h_{ek}N^*_{i5}V^*_{jk}, \qquad \lambda^{e}_{2ijk}=h_{ek}\left(N^*_{i5}U^*_{jk}-N^*_{ik}U^*_{j4}\right), \qquad \lambda^{e}_{3ijk}=h_{ek}V^*_{jk}U^*_{i4}. \tag{29}$$

The maximum value for these couplings is given in the tables 1 and 2. The couplings $\lambda^{u}_{1,2}$ and $\lambda^{d}_{1,2}$ (table 1) are all small except for the ones with one quark of squark belonging to the third family which are not relevant for our purposes here.

TABLE 2
Upper limits for the R-parity breaking Yukawa couplings in the leptonic sector

I	J	K	λ^{ν}_{1ijk}	λ^{ν}_{2ijk}	λ^{e}_{1ijk}	λ^{e}_{2ijk}	λ^{e}_{3ijk}
1	1	1	7.1×10^{-29}	4.8×10^{-22}	6.3×10^{-16}	6.5×10^{-10}	7.5×10^{-10}
1	1	2	4.2×10^{-28}	1.6×10^{-20}	2.6×10^{-24}	1.5×10^{-7}	1.6×10^{-16}
1	1	3	4.9×10^{-27}	3.3×10^{-15}	5.4×10^{-22}	3.2×10^{-13}	$1\ \times10^{-14}$
1	2	1	$2\ \times10^{-23}$	3.2×10^{-18}	1.2×10^{-26}	1.1×10^{-7}	7.6×10^{-19}
1	2	2	4.3×10^{-21}	$7\ \times10^{-16}$	1.3×10^{-13}	1.3×10^{-7}	1.5×10^{-7}
1	2	3	$9\ \times10^{-22}$	4.2×10^{-11}	6.5×10^{-19}	1.9×10^{-13}	$6\ \times10^{-12}$
1	3	1	6.6×10^{-17}	3.5×10^{-15}	1.5×10^{-25}	1.6×10^{-6}	2.9×10^{-18}
1	3	2	3.9×10^{-16}	$1\ \times10^{-14}$	3.9×10^{-20}	3.5×10^{-4}	3.5×10^{-13}
1	3	3	$5\ \times10^{-15}$	2.4×10^{-8}	2.2×10^{-12}	3.9×10^{-11}	2.6×10^{-6}
2	1	1	1.7×10^{-22}	1.6×10^{-20}	1.7×10^{-8}	1.7×10^{-8}	1.1×10^{-7}
2	1	2	3.8×10^{-20}	3.4×10^{-18}	6.6×10^{-15}	1.3×10^{-7}	3.4×10^{-14}
2	1	3	8.6×10^{-21}	$2\ \times10^{-13}$	5.2×10^{-16}	1.4×10^{-6}	2.9×10^{-14}
2	2	1	4.7×10^{-13}	$7\ \times10^{-16}$	3.2×10^{-17}	6.5×10^{-10}	1.6×10^{-16}
2	2	2	1.4×10^{-10}	3.4×10^{-13}	3.5×10^{-6}	3.9×10^{-6}	2.3×10^{-5}
2	2	3	2.2×10^{-11}	1.5×10^{-8}	1.3×10^{-11}	3.9×10^{-4}	1.5×10^{-10}
2	3	1	$5\ \times10^{-10}$	6.9×10^{-15}	1.5×10^{-19}	1.7×10^{-6}	8.1×10^{-18}
2	3	2	1.7×10^{-7}	2.8×10^{-13}	7.5×10^{-13}	3.3×10^{-4}	$9\ \times10^{-12}$
2	3	3	3.7×10^{-8}	1.9×10^{-7}	5.9×10^{-5}	3.7×10^{-5}	3.9×10^{-4}
3	1	1	1.9×10^{-21}	3.3×10^{-15}	1.7×10^{-6}	1.7×10^{-6}	1.7×10^{-6}
3	1	2	1.1×10^{-20}	$2\ \times10^{-13}$	3.4×10^{-14}	8.2×10^{-8}	2.9×10^{-14}
3	1	3	2.4×10^{-18}	$7\ \times10^{-12}$	2.2×10^{-11}	2.6×10^{-6}	2.2×10^{-11}
3	2	1	3.5×10^{-14}	4.2×10^{-11}	1.7×10^{-16}	$4\ \times10^{-10}$	1.4×10^{-16}
3	2	2	3.6×10^{-12}	1.5×10^{-8}	3.5×10^{-4}	3.5×10^{-4}	3.5×10^{-4}
3	2	3	1.4×10^{-10}	1.1×10^{-8}	3.3×10^{-8}	3.7×10^{-5}	3.3×10^{-8}
3	3	1	1.6×10^{-8}	2.4×10^{-8}	6.1×10^{-15}	8.5×10^{-10}	6.1×10^{-15}
3	3	2	$2\ \times10^{-7}$	1.8×10^{-7}	1.9×10^{-9}	3.9×10^{-6}	1.9×10^{-9}
3	3	3	3.5×10^{-5}	5.8×10^{-5}	$6\ \times10^{-3}$	3.9×10^{-4}	$6\ \times10^{-3}$

However they can have interesting effects in the top physics [29]. The couplings $\lambda^{\nu}_{1,2}$ include the couplings to the majoron after proper normalization. For instance λ^{ν}_{13i3} and λ^{ν}_{1i33} are responsible for the decay $\nu_\tau \to \nu_i J$ and the couplings λ^{ν}_{2ij3} can give rise to flavour violating decays of the type $\ell_i \to \ell_j$ J and other effects analysed in ref. [25].

4. Experimental constraints

All supersymmetric extensions of the standard model, are constrained by data that follow from several collider experiments, such as the recent LEP data on Z decays and by $\bar{p}p$ collider data, e.g. on $W^{\pm}$, Z and gluino production. This certainly applies to the spontaneously broken R-parity models. In addition to these constraints, there are important restrictions, characteristic of broken R-parity models, related to weak interactions and neutrino mass considerations [28]. These follow from laboratory, astrophysics and cosmology. They play a very important role in restricting spontaneously broken R-parity because they are found to exclude many parameter choices that are otherwise allowed by the collider constraints, while the converse is not true. The most relevant constraints have been listed in ref. [11] and include neutrinoless double beta decay and neutrino oscillation limits, direct searches for anomalous peaks at meson decays, limits on the neutrino masses, cosmological limits on neutrino lifetime, limits on muon and tau lifetimes, universality constraints, and limits on lepton flavour violating decays.

All of these constraints have to be taken into account. In order to do that systematically we have assumed the following values for the model parameters:

$$
\begin{aligned}
& v_{\mathrm{L}} = v_{\mathrm{L}3} = 100\ \mathrm{MeV}, && v_{\mathrm{L}1} = v_{\mathrm{L}2} = 0. \\
& v_{\mathrm{R}} = v_{\mathrm{R}3} = 1000\ \mathrm{GeV}, && v_{\mathrm{R}1} = v_{\mathrm{R}2} = 0, \\
& v_{\mathrm{S}} = 1000\ \mathrm{GeV}, && \tan\beta = v_{\mathrm{u}}/v_{\mathrm{d}} = 3,
\end{aligned}
\tag{30}
$$

and we have varied randomly the SUSY parameters μ, M_2, and the parameters $h_{i,3}$ in the interesting range given by

$$
\begin{aligned}
& -1000 \leqslant \frac{\mu}{\mathrm{GeV}} \leqslant 1000, && 20 \leqslant \frac{M}{\mathrm{GeV}} \leqslant 1000, \\
& 10^{-10} \leqslant h_{\nu 13},\ h_{\nu 23} \leqslant 10^{-1}, && 10^{-5} \leqslant h_{\nu 33} \leqslant 10^{-1}.
\end{aligned}
\tag{31}
$$

We have performed a careful sampling of the points in our parameter space that are allowed by all constraints discussed above in order to evaluate the attainable value of the couplings. The allowed values for the diagonal (R-parity conserving) couplings are given in eqs. (23)–(25). The results for the R-parity breaking couplings are plotted in figs. 1 and 2 for the gauge couplings (with the standard fermion in the third family) and listed in tables 1 and 2 for the Yukawa couplings.

5. Spontaneously broken R-parity at hadron colliders

The presence of R-parity breaking couplings modifies both the production mechanism for the supersymmetric particles and their decays. Concerning the production, one is no longer forced to produce the SUSY particles in pairs. Again, unlike the models with explicitly broken R-parity, in the spontaneous breaking case both the standard pair production and the single production take place through the gauge couplings.

Therefore at hadron colliders the main production mechanism is of the Drell–Yan type,

$$ab \to \chi_i \chi_j X, \tag{32}$$

where χ_i and χ_j can be both supersymmetric particles (standard SUSY pair production) or one standard and one supersymmetric (R-parity breaking single production). We label a and b the colliding hadrons. These processes are mediated by $\mathrm{W}^{\pm}$ or Z exchange.

The differential cross section of these processes is given by

$$\mathrm{d}\sigma(a+b\to\chi_i\chi_j X)$$
$$=\sum_{n,m}\int \mathrm{d}x_a\,\mathrm{d}x_b\, f_n^a(x_a, M^2) f_m^b(x_b, M^2)\,\mathrm{d}\sigma(n+m\to\chi_i\chi_j), \tag{33}$$

where x_a and x_b are the momentum fractions carried by the interacting partons in the protons a and b, respectively. M is the invariant mass of the pair χ_i, χ_j, the functions $f_{n,m}^{a,b}(x_{a,b}, M^2)$ are the parton distribution functions describing the probability of finding the parton $n(m)$ in the hadron $a(b)$ with momentum fraction $x_{a(b)}$.

For the case of Z mediated processes one gets

$$\frac{\mathrm{d}\sigma}{\mathrm{d}M\,\mathrm{d}y\,\mathrm{d}\cos\theta}$$
$$=K(M^2)\frac{Mg^4}{96\pi c_{\mathrm{w}}^4}\sqrt{M_{ij}}\sum_q\left[g_q^{\mathrm{S}}(y, M) S_q(M,\cos\theta)\right.$$
$$\left.+g_q^{\mathrm{A}}(y, M) A_q(M,\cos\theta)\right], \tag{34}$$

where

$$g_q^{\mathrm{S,A}}(y, M)=x_a x_b\left[f_q^a(x_a, M^2) f_{\bar{q}}^b(x_b, M^2)\pm f_{\bar{q}}^a(x_a, M^2) f_q^b(x_b, M^2)\right], \tag{35}$$

$$S_q=P\left(v_q^2+a_q^2\right)\left\{C_{ij}\left[N_{ij}+M_{ij}\cos^2\theta\right]+8x_i x_j D_{ij}\right\},$$
$$A_q=4P\cos\theta\; a_q v_q E_{ij}\sqrt{M_{ij}}, \tag{36}$$

with

$$v_q=\frac{T_3}{2}-Q_q s_w^2, \qquad a_q=-\frac{T_3}{2}, \tag{37}$$

and

$$P=\frac{1}{\left[\left(M^2-M_{\mathrm{Z}}^2\right)^2+\left(\Gamma_{\mathrm{Z}} M_{\mathrm{Z}}\right)^2\right]}. \tag{38}$$

The phase space functions N_{ij} and M_{ij} are

$$N_{ij} = 1 - \left(x_i^2 - x_j^2\right)^2,$$

$$M_{ij} = 1 - 2\left(x_i^2 + x_j^2\right) + \left(x_i^2 - x_j^2\right)^2, \tag{39}$$

with $x_a = m_a/M$. The angle θ is the scattering angle in the parton center-of-mass system (the angle between the parton in a and the final positive lepton or with the neutralino). $K(M^2)$ factorizes the QCD radiative corrections

$$K(M^2) = \exp\tfrac{2}{3}\pi\alpha_s(M^2),$$

$$\alpha_s(M^2) = \frac{12\pi}{33 - 2n_f}\frac{1}{\log(M^2/\Lambda^2)},$$

where n_f is the number of flavours and Λ is the QCD constant. For the case of charged leptons:

$$C_{ij} = O'^2_{Lij} + O'^2_{Rij},$$

$$D_{ij} = O'_{Lij}O'_{Rij}\eta_i\eta_j, \tag{40}$$

$$E_{ij} = O'^2_{Lij} - O'^2_{Rij},$$

whereas for neutral leptons:

$$C_{ij} = \frac{2-\delta_{ij}}{2}\left(O''^2_{Lij} + O''^2_{Rij}\right) = (2-\delta_{ij})O''^2_{Lij},$$

$$D_{ij} = \frac{2-\delta_{ij}}{2}O''_{Lij}O''_{Rij}\epsilon_i\epsilon_j = -\frac{2-\delta_{ij}}{2}\left(O''_{Lij}\right)^2\epsilon_i\epsilon_j,$$

$$E_{ij} = \frac{2-\delta_{ij}}{2}\left(O''^2_{Rij} - O''^2_{Lij}\right) = 0. \tag{41}$$

The asymmetric part of the cross section of the neutralinos vanishes because they are majorana particles, which means that their coupling to the Z is purely axial or purely vectorial (under the assumption of CP conservation).

For the case of $W^{\pm}$ mediated processes we have

$$\frac{d\sigma}{dM\, dy\, d\cos\theta}\left(ab \to \chi_i^+ \chi_j^0 X\right)$$

$$= \frac{K(M^2) M g^4}{192 \times 4\pi} \sqrt{M_{ij}} \sum_{qq'} |V_{qq'}|^2 \left[g^{S}_{qq'}(y, M) S(M, \cos\theta) \right.$$

$$\left. + g^{A}_{qq'}(y, M) A(M, \cos\theta) \right] \tag{42}$$

with

$$g^{S,A}_{qq'}(y, M) = x_a x_b \left[f_q^a(x_a, M^2) f_{\bar{q}'}^b(x_b, M^2) \pm f_{\bar{q}'}^a(x_a, M^2) f_q^b(x_b, M^2) \right], \tag{43}$$

$$S_q = P\left\{ \left(K^2_{Lij} + K^2_{Rij} \right) \left[N_{ij} + M_{ij} \cos^2\theta \right] + 8 x_i x_j K_{Lij} K_{Rij} \right\},$$

$$A_q = 2P \cos\theta \left(K^2_{Lij} - K^2_{Rij} \right) \sqrt{M_{ij}}\,. \tag{44}$$

q is an up-type quark and q' is a down-type quark, and $V_{qq'}$ is the CKM matrix. P is the corresponding propagator factor for the W. In this case we define θ to be the angle between χ_i^+ and the parton in the hadron a.

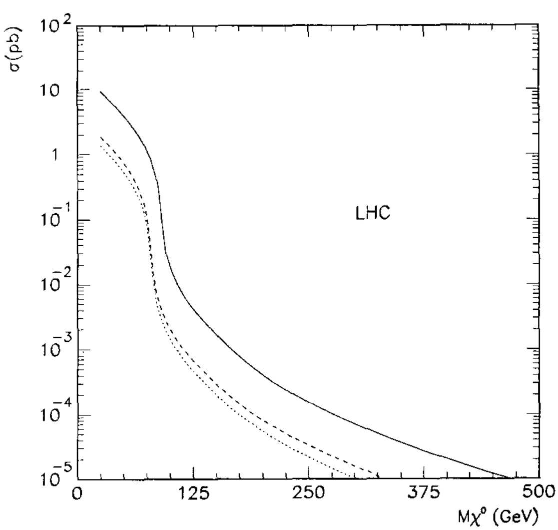

Fig. 3. Attainable cross section for the single production of the lightest neutralino at LHC. The full line corresponds to $pp \to \chi^0 \nu_\tau X$, the dashed line corresponds to $pp \to \chi^0 \tau^+ X$ and the dotted line corresponds to $pp \to \chi^0 \tau^- X$.

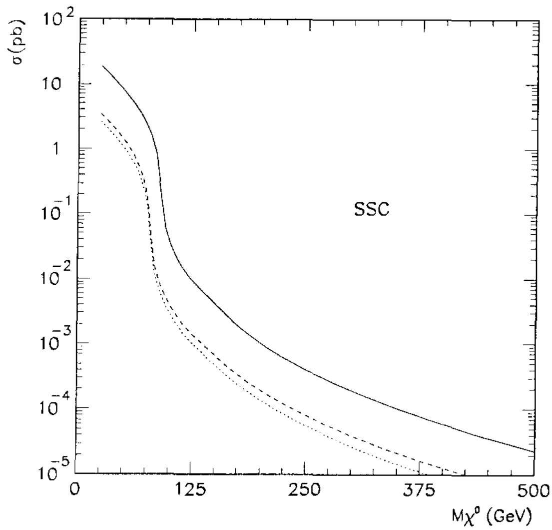

Fig. 4. The same as fig. 3 for SSC.

For the W^- mediated process we have

$$\frac{d\sigma}{dM\,dy\,d\cos\theta}\left(ab\to\chi_i^-\chi_j^0X\right)=\frac{d\sigma}{dM\,dy\,d\cos\theta}\left(ab\to\chi_i^+\chi_j^0X\right)(q\leftrightarrow q'). \quad (45)$$

In figs. 3–6 we plot the single production cross section for the lightest chargino and the lightest neutralino both at LHC and SSC. The expected integrated

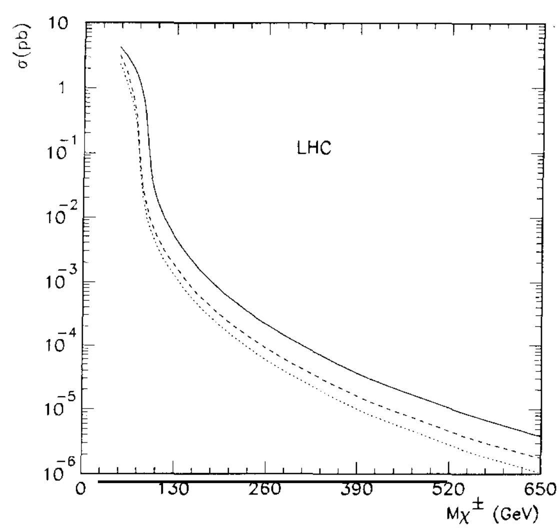

Fig. 5. Attainable cross section for the single production of the lightest chargino at LHC. The full line corresponds to $pp\to\chi^+\tau^-X$, the dashed line corresponds to $pp\to\chi^+\nu_\tau X$ and the dotted line corresponds to $pp\to\chi^-\nu_\tau X$.

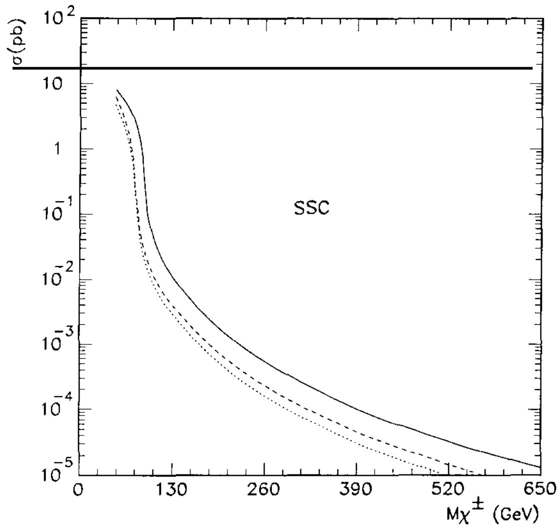

Fig. 6. The same as fig. 5 for SSC.

luminosity at these colliders is $L = 10^4$–10^5 pb^{-1}. Therefore at LHC it is possible to have more than 10 events for neutralino mass $M_{\chi^0} \lesssim 160$–280 GeV and chargino mass $M_{\chi^+} \lesssim 180$–320 GeV. The lower mass corresponds to an integrated luminosity of $L = 10^4$ pb^{-1} and the upper to $L = 10^5$ pb^{-1}. The corresponding masses for SSC are $M_{\chi^0} \lesssim 200$–360 GeV and $M_{\chi^+} \lesssim 220$–400 GeV.

6. Chargino and neutralino decays

In models with R-parity conservation all supersymmetric particles have cascade decays finishing in the LSP. However if R-parity is broken new decay channels are open and the supersymmetric particles can decay both *directly* to the standard states breaking R-parity or through R-parity conserving cascade decays that will finish in the lightest massive SUSY particle which has to decay to standard states breaking R-parity.

For simplicity we are going to study the decays of the lightest neutralino and the lightest chargino, which on the other hand, one expects would be the earliest-produced supersymmetric particles. Heavier states would have cascade decays that we are not going to consider here.

The lightest neutralino has to decay always to standard states breaking R-parity. If its mass is lower than the mass of the gauge bosons it decays to the three-body final states

$$\chi^0 \to \nu_j \mathrm{f}\bar{\mathrm{f}} \quad \text{with width} \quad \Gamma^0_{13bj} = 8\left(v_\mathrm{f}^2 + a_\mathrm{f}^2\right)\Gamma^{3b}\left(M_{\chi^0}, 0, M_\mathrm{Z}, O''_{\mathrm{L}4j}, O''_{\mathrm{R}4j}\right),$$

$$\chi^0 \to \ell_j \mathrm{f}_\mathrm{u}\bar{\mathrm{f}}_\mathrm{d} \quad \text{with width} \quad \Gamma^0_{23bj} = \Gamma^{3b}\left(M_{\chi^0}, 0, M_\mathrm{W}, K_{\mathrm{L}j4}, K_{\mathrm{R}j4}\right). \tag{46}$$

For the decay $\chi^0 \to \nu_j \ell_j^+ \ell_j^-$ there is interference between the charged and neutral current and the width is given by

$$\Gamma^0_{33bj} = \Gamma^{3b\prime}(M_{\chi^0},\, K_{\mathrm{L}j4},\, K_{\mathrm{R}j4},\, O''_{\mathrm{L}4j},\, O''_{\mathrm{R}4j}). \tag{47}$$

On the other hand if the neutralino is heavier than the gauge bosons the two-body decays

$$\chi^0 \to \ell_j \mathrm{W} \quad \text{with width} \quad \Gamma^0_{\mathrm{W}j} = \tfrac{1}{2}\Gamma^{2b}(M_{\chi^0},\, 0,\, M_{\mathrm{W}},\, K_{\mathrm{L}j4},\, K_{\mathrm{R}j4}),$$

$$\chi^0 \to \nu_j \mathrm{Z} \quad \text{with width} \quad \Gamma^0_{\mathrm{Z}j} = \Gamma^{2b}(M_{\chi^0},\, 0,\, M_{\mathrm{Z}},\, O''_{\mathrm{L}4j},\, O''_{\mathrm{R}4j}) \tag{48}$$

are dominant over the three-body decays. The explicit expressions for the widths Γ^{3b}, $\Gamma^{3b\prime}$, and Γ^{2b} are given in appendix A. The existence of the majoron implies that in spontaneously broken R-parity, the neutralino can always decay invisibly to

$$\chi^0 \to \nu_j \mathrm{J} \tag{49}$$

with a decay width

$$\Gamma^0_{\mathrm{J}j} = \frac{1}{32\pi} M_{\chi^0}\left(C^2_{\mathrm{L}4j} + C^2_{\mathrm{R}4j}\right),$$

$$C_{\mathrm{L}ij} = -\epsilon_i \epsilon_j C_{\mathrm{R}ij} = \sum_{k=1}^{3} \left(\epsilon_j N_{ik} N_{j4} + \epsilon_i N_{jk} N_{i4}\right) h_{\nu k3} \frac{v_{\mathrm{R}}}{\sqrt{2}\, V}. \tag{50}$$

The lightest chargino can decay both directly to the standard particles breaking R-parity or to the lightest neutralino conserving R-parity. If lighter than the gauge bosons, the neutralino decays to the three-body final states

$$\chi^+ \to \ell_j^+ \mathrm{f}\bar{\mathrm{f}} \quad \text{with width} \quad \Gamma^+_{13bj} = 8\left(v_{\mathrm{f}}^2 + a_{\mathrm{f}}^2\right)\Gamma^{3b}(M_{\chi^+},\, 0,\, M_{\mathrm{Z}},\, O'_{\mathrm{L}4j},\, O'_{\mathrm{R}4j}),$$

$$\chi^+ \to \bar{\nu}_j \mathrm{f_u}\bar{\mathrm{f}}_{\mathrm{d}} \quad \text{with width} \quad \Gamma^+_{23bj} = \Gamma^{3b}(M_{\chi^+},\, 0,\, M_{\mathrm{W}},\, K_{\mathrm{L}4j},\, K_{\mathrm{R}4j}),$$

$$\chi^+ \to \bar{\nu}_j \nu_j \ell_j^+ \quad \text{with width} \quad \Gamma^+_{33bj} = \Gamma^{3b\prime}(M_{\chi^+},\, K_{\mathrm{L}4j},\, K_{\mathrm{R}4j},\, O'_{\mathrm{L}4j},\, O'_{\mathrm{R}4j}) \tag{51}$$

breaking R-parity or to the R-odd three-body final state

$$\chi^+ \to \chi^0 \mathrm{f_u}\bar{\mathrm{f}}_{\mathrm{d}} \quad \text{with width} \quad \Gamma^+_{3b0} = \Gamma^{3b}(M_{\chi^+},\, M_{\chi^0},\, M_{\mathrm{W}},\, K_{\mathrm{L}44},\, K_{\mathrm{R}44}). \tag{52}$$

However when the breaking of R-parity is large enough to have large cross section for the single production of the chargino, this decay can become less important than the previous ones. When heavier than the gauge bosons but lighter than $M_W + M_{\chi^0}$ the chargino decays to

$$\chi^+ \to \nu_j W^+ \quad \text{with width} \quad \Gamma^+_{Wj} = \tfrac{1}{2}\Gamma^{2b}(M_{\chi^+}, 0, M_W, K_{L4j}, K_{R4j}),$$
$$\chi^+ \to \ell^+_j Z \quad \text{with width} \quad \Gamma^+_{Zj} = \Gamma^{2b}(M_{\chi^+}, 0, M_Z, O'_{L4j}, O'_{R4j}) \tag{53}$$

together with the R-parity conserving three-body decay in eq. (52). When $M_{\chi^+} \geqslant M_W + M_{\chi^0}$ it can also decay to

$$\chi^+ \to \chi^0 W^+ \quad \text{with width} \quad \Gamma^+_{Wj} = \tfrac{1}{2}\Gamma^{2b}(M_{\chi^+}, M_{\chi^0}, M_W, K_{L44}, K_{R44}). \tag{54}$$

This decay is kinematically possible when the chargino is relatively heavy and the breaking of R-parity cannot be large enough to make it unimportant. Again the chargino can always decay in the majoron

$$\chi^+ \to \ell_j J \tag{55}$$

with width given in eq. (50) where now

$$C_{Lj4} = \frac{v_R}{\sqrt{2}V}\sum_{k=1}^{3} h_{\nu k3}\eta_4 U_{4k} V_{j4}, \qquad C_{Rj4} = -\frac{v_R}{\sqrt{2}V}\sum_{k=1}^{3} h_{\nu k3}\eta_j U_{jk} V_{44}. \tag{56}$$

Both the neutralino and the chargino decay widths are larger for the lepton j belonging to the third family.

7. Signals

We now detail the signals of the first kinematically accessible supersymmetric states, which include the single production of the lightest neutralino, the single production of the lightest chargino, and the R-parity conserving pair production of the lightest neutralino, all of them followed by the subsequent decay of the SUSY particle.

The possible final states for the *single production of the lightest neutralino* are summarized in table 3 for the two production mechanisms through neutral and charged current and for different mass ranges of the neutralino. In all ranges the neutralino may decay invisibly to majoron and this may dilute any possible signal. However, a numerical inspection of the different branching ratios shows that, in some regions of the parameter space, the decay rate for this channel can be smaller or of the same order than the other visible channels. One particularly promising channel would be the *single production of the neutralino via charged current* together with a τ-lepton, followed by the decay of the neutralino in a τ of

TABLE 3
Possible final states following from the single production of the lightest neutralino

Production mechanism	$M_{\chi^0} \leqslant M_W$		$M_{\chi^0} \geqslant M_Z$	
	Decay mode	Signal	Decay mode	Signal
$pp \to \chi^0 \nu_\tau X$	$\chi^0 \to 3$ body	$4\,\not{p}_T$ $2\,\not{p}_T + (\ell^+\ell^-)$ $2\,\not{p}_T + (\tau\ell)$ $\not{p}_T + [\tau, \not{p}_T] + 2$ Jets	$\chi^0 \to W\tau$ $\chi^0 \to Z\nu_\tau$ $\chi^0 \to J\nu_\tau$	$\not{p}_T + (\tau W)$ $2\,\not{p}_T + Z$ $3\,\not{p}_T$
	$\chi^0 \to \nu J$	$3\,\not{p}_T$		
$pp \to \chi^0 \tau X$	$\chi^0 \to 3$ body	$3\,\not{p}_T + \tau$ $\not{p}_T + (\ell^+\ell^-) + \tau$ $\not{p}_T + (\tau\ell) + \tau$ $\tau + [\tau, \not{p}_T] + 2$ Jets	$\chi^0 \to W\tau$ $\chi^0 \to Z\nu_\tau$ $\chi^0 \to J\nu_\tau$	$\tau + (\tau W)$ $\tau + \not{p}_T + Z$ $\tau + 2\,\not{p}_T$
	$\chi^0 \to \nu J$	$\tau + 2\,\not{p}_T$		

The total charge of the particles between parenthesis must be zero. The charged particles without charge assignment can be either positive or negative. The number preceding the missing transverse momentum signal counts the number of neutral particles contributing to that missing momentum.

the *same charge* plus an electron or a muon and a neutrino. This signal is possible because the neutralino is a Majorana particle. The branching ratio for this decay in the region where the single production is big is of order 10^{-2}. This signal has the advantage that there is no background from two gauge boson production. The one source of background for this signal would be 3W production either direct or through $t\bar{t}W$, with the top decaying into Wb. The W's can then decay in 2τ's with the same sign and a lepton with a different sign. The cross sections for these processes have been evaluated in ref. [32] and they are of order 4×10^{-4} pb at LHC and 10^{-3} pb at SSC. Therefore before including any cut the background is already lower than the signal for all values of $M_{\chi^0} \lesssim 90$ GeV, as can be seen from fig. 3 for LHC and fig. 4 for SSC. However we must point out that a 90 GeV neutralino would lead to one single event per year at SSC and LHC (assuming an integrated luminosity of 10^4 pb^{-1}), while 10 events per year would be expected only if the neutralino is lighter than 80 GeV. On the other hand 100 events per year would require $M_{\chi^0} \lesssim 70$ GeV.

In table 4 we give the possible final states for the single production of the lightest chargino for the different mass ranges. Here the situation is a bit more complicated because one has to include the possible intermediate decay to the lightest neutralino in eqs. (52) and (53). Again here the decay to the majoron would dilute the signal but as in the previous case its width can be small in the range of masses where the single production cross section is big. We have also here final states with multi-leptons which can be highly visible.

For the case of the pair production of the lightest neutralino we give the possible final states in table 5. Here R-parity breaking converts the $\not{p}_T$ signal characteristic of the models with R-parity conservation in visible signals including

TABLE 4
Same as in table 3 for the lightest chargino

Production mechanism	$M_{\chi^+} \leqslant M_W$		$M_W + M_{\chi^0} \geqslant M_{\chi^\pm} \geqslant M_Z$		$M_{\chi^+} \geqslant M_W + M_{\chi^0}$	
	Decay mode	Signal	Decay mode	Signal	Decay mode	Signal
$pp \to (\chi^\pm \tau)X$	$\chi^\pm \to 3$ body	$\left\{\begin{matrix}(\tau^+\tau^-)+(\ell^+\ell^-)\\ 2\not p_T+(\tau\ell)\\ (\tau^+\tau^-)+2\text{ Jets}\\ \not p_T+\tau+2\text{ Jets}\end{matrix}\right.$	$\chi^\pm \to W\nu_\tau$ $\chi^\pm \to Z\tau$ $\chi^\pm \to J\tau$	$\not p_T+(W\tau)$ $Z+(\tau^+\tau^-)$ $\not p_T+(\tau^+\tau^-)$	$\chi^\pm \to W\nu_\tau$ $\chi^\pm \to Z\tau$ $\chi^\pm \to J\tau$	$\not p_T+(W\tau)$ $Z+(\tau^+\tau^-)$ $\not p_T+(\tau^+\tau^-)$
	$\chi^\pm \to \tau J$	$\not p_T+(\tau^+\tau^-)$				
	$\chi^\pm \to \chi^0 f\bar{f}'$	$\left\{\begin{matrix}(\ell\tau)+\not p_T\\ \tau+2\text{ Jets}\end{matrix}\right\} \otimes \left\{\begin{matrix}[2,3]\ \not p_T\\ \not p_T+(\ell^+\ell^-)\\ \not p_T+(\tau\ell)\\ [\tau,\not p_T]+2\text{ Jets}\\ (W\tau)\\ \not p_T+Z\end{matrix}\right.$	$\chi^\pm \to \chi^0 f\bar{f}'$	$\left\{\begin{matrix}(\ell\tau)+\not p_T\\ \tau+2\text{ Jets}\end{matrix}\right\} \otimes \left\{\begin{matrix}[2,3]\ \not p_T\\ \not p_T+(\ell^+\ell^-)\\ \not p_T+(\tau\ell)\\ [\tau,\not p_T]+2\text{ Jets}\\ (W\tau)\\ \not p_T+Z\end{matrix}\right.$	$\chi^\pm \to \chi^0 W$	$(\tau W)+\left\{\begin{matrix}[2,3]\ \not p_T\\ \not p_T+(\ell^+\ell^-)\\ \not p_T+(\tau\ell)\\ [\tau,\not p_T]+2\text{ Jets}\\ (W\tau)\\ \not p_T+Z\end{matrix}\right.$
$pp \to \chi^\pm \nu_\tau X$	$\chi^\pm \to 3$ body	$\left\{\begin{matrix}\not p_T+\tau+(\ell^+\ell^-)\\ 3\not p_T+\ell\\ \not p_T+[\tau,\not p_T]+2\text{ Jets}\end{matrix}\right.$	$\chi^\pm \to W\nu_\tau$ $\chi^\pm \to Z\tau$ $\chi^\pm \to J\tau$	$2\not p_T+W$ $\not p_T+Z+\tau$ $2\not p_T+\tau$	$\chi^\pm \to W\nu_\tau$ $\chi^\pm \to Z\tau$ $\chi^\pm \to J\tau$	$2\not p_T+W$ $\not p_T+Z+\tau$ $2\not p_T+\tau$
	$\chi^\pm \to \tau J$	$2\not p_T+\tau$				
	$\chi^\pm \to \chi^0 f\bar{f}'$	$\left\{\begin{matrix}2\not p_T+\ell\\ \not p_T+2\text{ Jets}\end{matrix}\right\} \otimes \left\{\begin{matrix}[2,3]\ \not p_T\\ \not p_T+(\ell^+\ell^-)\\ \not p_T+(\tau\ell)\\ [\tau,\not p]+2\text{ Jets}\\ (W\tau)\\ \not p_T+Z\end{matrix}\right.$	$\chi^\pm \to \chi^0 f\bar{f}'$	$\left\{\begin{matrix}2\not p_T+\ell\\ \not p_T+2\text{ Jets}\end{matrix}\right\} \otimes \left\{\begin{matrix}[2,3]\ \not p_T\\ \not p_T+(\ell^+\ell^-)\\ \not p_T+(\tau\ell)\\ [\pi,\not p_T]+2\text{ Jets}\\ (W\tau)\\ \not p_T+Z\end{matrix}\right.$	$\chi^\pm \to \chi^0 W$	$\not p_T W+\left\{\begin{matrix}[2,3]\ \not p_T\\ \not p_T+(\ell^+\ell^-)\\ \not p_T+(\tau\ell)\\ [\tau,\not p_T]+2\text{ Jets}\\ (W\tau)\\ \not p_T+Z\end{matrix}\right.$

TABLE 5
Possible final states following from the pair production of the lightest neutralino

Production mechanism	$M_{\chi^0} \leqslant M_W$		$M_{\chi^0} \geqslant M_Z$	
	Decay mode	Signal	Decay mode	Signal
$pp \to \chi^0\chi^0 X$	$\chi^0 \to 3$ body $\chi^0 \to 3$ body	$6\not p_T$ $4\not p_T + (\ell^+\ell^-)$ $4\not p_T + (\tau\ell)$ $2\not p_T + (\ell^+\ell^-) + (\ell'^+\ell'^-)$ $2\not p_T + (\ell^+\ell^-) + (\tau\ell)$ $2\not p_T + (\tau\ell) + (\tau\ell')$ $3\not p_T + [\tau, \not p_T] + 2$ Jets $\not p_T + (\ell^+\ell^-) + [\tau, \not p_T] + 2$ Jets $\not p_T + (\tau\ell) + [\tau, \not p_T] + 2$ Jets $[\tau, \not p_T] \otimes [\tau, \not p_T] + 4$ Jets	$\chi^0 \to W\tau$, $\chi^0 \to W\tau$ $\chi^0 \to W\tau$, $\chi^0 \to J\nu$ $\chi^0 \to W\tau$, $\chi^0 \to Z\nu$ $\chi^0 \to Z\nu$, $\chi^0 \to Z\nu$	$2(W\tau)$ $2\not p_T + (W\tau)$ $(W\tau) + Z + \not p_T$ $2Z + 2\not p_T$
	$\chi^0 \to 3$ body $\chi^0 \to J\nu$	$5\not p_T$ $3\not p_T + (\ell^+\ell^-)$ $3\not p_T + (\tau\ell)$ $2\not p_T + [\tau, \not p_T] + 2$ Jets	$\chi^0 \to Z\nu$, $\chi^0 \to J\nu$ $\chi^0 \to J\nu$, $\chi^0 \to J\nu$	$Z + 3\not p_T$ $4\not p_T$
	$\chi^0 \to J\nu$ $\chi^0 \to J\nu$	$4\not p_T$		

multi-leptons, due to the decay of the lightest neutralino. However if the decay into majoron happens to be dominant there would be no signal.

In the three tables we have explicitly denoted when the charged lepton has to be a τ because the R-parity breaking is bigger in the 3rd family. Any other lepton denoted by ℓ can belong to any generation. We have also the gauge bosons in the final states because it gives more information about the kinematics of the fermions produced in their decays.

When going to chargino–neutralino masses higher than those discussed here one has to deal with the cascade decays what would lead to a list of possible final states too large to consider now. In general we can say that if we are in the sector of the theory where the R-parity breaking is small to be detected in the single production of the SUSY states, one would have the double production followed by the standard cascade decays of the MSSM ending in the lightest neutralino which now would decay. Therefore R-parity breaking would show up only in the last step of the chain by converting the $\not p_T$ signal into visible states. When the breaking of R-parity is big enough to have both single and double production the chain of possible final states is enormous. One can then select a final state with promising characteristic signal and follow the decay chain that would lead to that final state.

In the study of chargino and neutralino production at hadron colliders in the MSSM, special emphasis has been put in the signal coming from multilepton decays of the charginos and neutralinos [31], mainly those with 3 and 5 charged leptons in the final state. The main sources of background for these signals are the production and decay of multiple gauge bosons and the top quark pair production followed by semileptonic chain decays of the top. It is not clear whether the signal can be above background in the MSSM. As we have seen in the case of R-parity breaking models *new sources of multilepton events are possible* that can only enhance the signal with respect to that in the MSSM and in consequence its detection becomes more hopeful.

8. Conclusion

We have shown that, in spontaneously broken R-parity, the rates for single production of the lightest SUSY fermions at hadron supercolliders can be sizeable without conflicting any laboratory, cosmological or astrophysical observation. Moreover, the reach of LHC and SSC are similar, especially for the most optimistic integrated LHC luminosity corresponding to $L \sim 10^5$ pb^{-1} in a year. This is seen from figs. 3–6. For example at LHC it is possible to have more than 10 neutralino events for masses in the range $M_{\chi^0} \lesssim 160$–280 GeV and chargino mass $M_{\chi^+} \lesssim 180$–320 GeV. The lower masses correspond to $L = 10^4$ pb^{-1} and the upper to $L = 10^5$ pb^{-1}. The corresponding masses that can be explored at SSC are in the range $M_{\chi^0} \lesssim 200$–360 GeV to $M_{\chi^+} \lesssim 220$–400 GeV. We have also analysed all possible decay channels of the lowest lying SUSY fermions in spontaneously broken R-parity, as given in the tables. These define the complete map of the expected signals.

The possibility of having sizeable SUSY production rates in these models more than justifies the interest in performing detailed background studies and simulations of individual signals and this will be taken up elsewhere. It should, however, be clear from our discussion that the SUSY signatures can be more easily visible experimentally if R-parity is broken than if not. As a specific example of a realistic signature we suggested the production of two like-sign taus plus an opposite sign electron or muon. This process can arise from single neutralino production, followed by decay. The standard model background for this is negligible before applying any cuts, for neutralino mass below 90 GeV. From figs. 3 and 4 one sees that a 10 event signal could be possible if $M_{\chi^0} \lesssim 80$ GeV. However we want to remark that these results have still to be folded with the appropriate detection efficiencies. Only when these are known one will be able to conclude definitively whether such signal will be visible or not.

We thank Vernon Barger for suggesting the like-sign tau signal discussed in the text, and for pointing out the work of ref. [32].

Appendix A

In sect. 6 we have defined the following decay widths:

$$\Gamma^{2b}(m_i, m_j, m_{\rm B}, d_{\rm L}, d_{\rm R}) = \frac{G_{\rm F} m_i^3}{4\sqrt{2}\pi}\sqrt{\left[1-(x_j+x_{\rm B})^2\right]\left[1-(x_j-x_{\rm B})^2\right]}$$
$$\times\left\{(d_{\rm L}^2+d_{\rm R}^2)\left[(1-x_j^2)^2+x_{\rm B}^2(1+x_j^2)-2x_{\rm B}^4\right]\right.$$
$$\left.-12 d_{\rm L} d_{\rm R} x_j x_{\rm B}^2\right\}, \tag{A.1}$$

with $x_j = m_j/m_i$ and $x_{\rm B} = m_{\rm B}/m_i$, and

$$\Gamma^{3b}(m_i, m_j, m_{\rm B}, d_{\rm L}, d_{\rm R})$$
$$= \frac{G_{\rm F}^2 m_i^5}{8 x_{\rm B}^8 \pi^4}\int_0^{x_{\rm B}^2 - x_j^2} {\rm d}y\, y\left\{d_{\rm R}^2 y(x_{\rm B}^2 - x_j^2 - y) I_{11,11}(x_{\rm B}, x_j, y)\right.$$
$$-d_{\rm L}^2\left[\left(y^2 - y(2+x_{\rm B}^2 - x_j^2) + 1 - x_j^2 + x_{\rm B}^2\right) I_{11,11}(x_{\rm B}, x_j, y)\right.$$
$$\left.+(2y - 2 - x_{\rm B}^2 + x_j^2) I_{11} + I\right] + 2 d_{\rm L} d_{\rm R} x_j x_{\rm B}\left[-I_{11,11}(x_{\rm B}, x_j, y)\right.$$
$$\left.\left.+I_{11}(x_{\rm B}, x_j, y)\right]\right\}, \tag{A.2}$$

where now $x_{\rm B} = m_i/m_{\rm B}$. We have followed the notation of ref. [30]. For completeness we give here the value of the functions

$$I = \frac{\pi}{2}\gamma,$$
$$I_{11} = -\frac{\pi}{2y}\ln[1-\gamma y],$$
$$I_{11,11} = \frac{\pi}{2}\frac{\gamma}{1-\gamma y},$$
$$\gamma = 1 - \frac{x_j^2}{x_{\rm B}^2 - y}. \tag{A.3}$$

In the limit where $x_j = 0$:

$$\Gamma^{3b}(m_i, 0, m_{\rm B}, d_{\rm L}, d_{\rm R}) = \frac{G_{\rm F}^2 m_i^5}{192\pi^3}(d_{\rm L}^2 + d_{\rm R}^2) f(x_{\rm B}), \tag{A.4}$$

with

$$f(x_{\rm B}) = \frac{12}{x_{\rm B}^4}\left[-\tfrac{1}{6}x_{\rm B}^2 - \tfrac{1}{2} + \frac{1}{x_{\rm B}^2} + \frac{1-x_{\rm B}^2}{x_{\rm B}^4}\ln(1-x_{\rm B}^2)\right]. \tag{A.5}$$

For the case with interference between the charge and neutral current

$$\Gamma^{3b\prime}(m_i, c_{\rm L}, c_{\rm R}, d_{\rm L}, d_{\rm R})$$

$$= \frac{G_{\rm F}^2 m_i^5}{192\pi^3}\left[8(d_{\rm L}^2 + d_{\rm R}^2)(v_{\rm f}^2 + a_{\rm f}^2) f\left(\frac{m_i}{M_{\rm Z}}\right) + (c_{\rm L}^2 + c_{\rm R}^2) f\left(\frac{m_i}{M_{\rm W}}\right)\right.$$

$$\left. + 4(a_{\rm f} - v_{\rm f}) c_{\rm L} d_{\rm L}\, g\left(\frac{m_i}{M_{\rm Z}}, \frac{m_i}{M_{\rm W}}\right)\right], \tag{A.6}$$

where

$$g(x, y) = \frac{12}{y^6 x^6}\left\{-\tfrac{1}{2}x^2 y^2(-5x^2 - 5y^2 + 3x^2 y^2)\right.$$

$$-\tfrac{1}{2}(-3x^2 - 2y^2 + x^2 y^2) x^2 (1-y^2)\ln(1-y^2)$$

$$-\tfrac{1}{2}(-2x^2 - 3y^2 + x^2 y^2) y^2 (1-x^2)\ln(1-x^2)$$

$$+(x^2 + y^2)(x^2 y^2 - x^2 - y^2)$$

$$\times\left\{\ln(1-x^2)\ln\left(\frac{y^2 + x^2 - x^2 y^2}{x^2}\right) - \mathrm{Li}_2\left(\frac{y^2 - x^2 y^2}{y^2 + x^2 - x^2 y^2}\right)\right.$$

$$\left.\left. + \mathrm{Li}_2\left(\frac{y^2}{y^2 + x^2 - x^2 y^2}\right)\right\}\right\}. \tag{A.7}$$

References

[1] H. Haber and G. Kane, Phys. Rep. 117 (1985) 75;
H.P. Nilles, Phys. Rep. 110 (1984) 1

[2] L. Hall and M. Suzuki, Nucl. Phys. B231 (1984) 419;
S. Dimopoulos and L.J. Hall, Phys. Lett. B207 (1988) 210;
E. Ma and D. Ng, Phys. Rev. D41 (1990) 1005

[3] V. Barger, G.F. Giudice and T. Han, Phys. Rev. D40 (1989) 2987

[4] H. Dreiner and G.G. Ross, Oxford preprint OUTP-91-15P

[5] B.A. Campbell, S. Davison, J. Ellis and K. Olive, Phys. Lett. B256 (1991) 457;
W. Fischler, G.F. Giudice, R.G. Leigh and S. Paban, Phys. Lett. B258 (1991) 45

[6] V.A. Kuzmin, V.A. Rubakov and M.E. Shaposhnikov, Phys. Lett. B155 (1985) 36
[7] H. Dreiner and G.G. Ross, private communication
[8] R. Mohapatra and J.W.F. Valle, Phys. Lett. B186 (1987) 303;
S. Dimopoulos and L.J. Hall, Phys. Lett. B196 (1987) 135;
J. Cline and S. Raby, Ohio State preprint DOE/ER/01545-444 (1990)
[9] C. Aulakh and R. Mohapatra, Phys. Lett. B119 (1983) 136;
A. Santamaria and J.W.F. Valle, Phys. Lett. B195 (1987) 423; Phys. Rev. Lett. 60 (1988) 397; Phys. Rev. D39 (1989) 1780
[10] A. Masiero and J.W.F. Valle, Phys. Lett. B251 (1990) 273
[11] P. Nogueira, J.C. Romão and J.W.F. Valle, Phys. Lett. B251 (1990) 142
[12] J.W.F. Valle, Phys. Lett. B196 (1987) 157;
M.C. Gonzalez-Garcia and J.W.F. Valle, Nucl. Phys. B355 (1991) 330
[13] J. Schechter and J.W.F. Valle, Phys. Rev. D25 (1982) 774;
Y. Chikashige, R. Mohapatra and R. Peccei, Phys. Rev. Lett. 45 (1980) 1926
[14] J. Carter, talk at Joint Int. Lepton-Photon Symposium and Europhysics Conference on High Energy Physics, Geneva, July 1991, to appear in the proceedings
[15] M.C. Gonzalez-Garcia and Y. Nir, Phys. Lett. B232 (1990) 383;
P. Nogueira and J.C. Romão, Phys. Lett. B234 (1990) 371
[16] R. Mohapatra and J.W.F. Valle, Phys. Rev. D34 (1986) 1642
[17] J.W.F. Valle, Nucl. Phys. B (Proc. Suppl.) 11 (1989) 118
[18] J. Bernabeu, A. Santamaria, J. Vidal, A. Mendez and J.W.F. Valle, Phys. Lett. B187 (1987) 303;
G.C. Branco, M.N. Rebelo and J.W.F. Valle, Phys. Lett. B225 (1989) 385;
N. Rius and J.W.F. Valle, Phys. Lett. B246 (1990) 249
[19] M. Dittmar, M.C. Gonzalez-Garcia, A. Santamaria and J.W.F. Valle, Nucl. Phys. B332 (1990) 1;
M.C. Gonzalez-Garcia, A. Santamaria and J.W.F. Valle, Nucl. Phys. B342 (1990) 108
[20] J.W.F. Valle, *in* Neutrinos Beyond the Standard Model, Proc. WEIN-89, ed. P. Depommier (Editions Frontieres, Gif-sur-Yvette) p. 277; Nucl. Phys. B (Proc. Suppl.) 13 (1990) 195, 520
[21] J.C. Romão, C.A. Santos and J.W.F. Valle, Phys. Lett. B288 (1992) 311
[22] J.E. Kim, Phys. Rep. 150 (1987) 1
[23] J.W.F. Valle, Proc. Meeting on Tau-Charm Factory Detector and Machine, Seville, May 1991, Valencia preprint FTUV/91-35
[24] J.C. Romão and J.W.F. Valle, Phys. Lett. B272 (1991) 436; Nucl. Phys. B381 (1992) 87
[25] J.C. Romão, N. Rius and J.W.F. Valle, Nucl. Phys. B363 (1991) 369
[26] J. Schechter and J.W.F. Valle, Phys. Rev. D22 (1980) 2227
[27] L. Wolfenstein, Nucl. Phys. B186 (1981) 147;
J.W.F. Valle, Phys. Rev. D27 (1983) 1672;
J.W.F. Valle and M. Singer, Phys. Rev. D28 (1983) 540
[28] J.W.F. Valle, Gauge theories and the physics of neutrino mass, ed. A. Faessler, Prog. Part. Nucl. Phys. 26 (1991) 91 and references therein
[29] H. Dreiner and R. Phillips, Rutherford Laboratory preprint RAL-91-045
[30] L. Bento, J.C. Romão and A. Barroso, Phys. Rev. D33 (1986) 1488
[31] R. Barbieri et al., Nucl. Phys. B367 (1991) 28
[32] B. Barger, A.L. Stange and R.J.N. Phillips, MAD/PH/664, to appear in Phys. Rev. D