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Supersymmetry with Spontaneous R Parity Breaking in Z^0 Decays: the case of an Additional Z

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Abstract

Single production of SUSY particles in the decays of the Z^0 may proceed with large rates in models with spontaneously broken R parity. We focus on the case where there is a lepton number symmetry as part of the gauge group. In the simplest of such models there is a single additional neutral gauge boson and the strength of R_p violating interactions is related to that of the new gauge force. We study the phenomenological implications of the model for Z^0 decays, including the study of the rates for single chargino production in Z^0 decays, *i.e.*, $Z^0 \to \chi^{\pm}\tau^{\pm}$, as well as for the so-called *zen events*, and find that they may be measurable at LEP. The first process, characteristic of spontaneously broken R parity models, may proceed with a branching ratio as large as $few \times 10^{-5}$ and should lead to a very clear signature. *Zen* events have a different origin than in the minimal SUSY standard model (MSSM). Their rates may also be enhanced with respect to MSSM. Our estimates of these rates take into account all of the relevant observational constraints, such as those that follow from neutrino physics, from precise Z^0 property measurements, as well as SUSY particle searches at LEP.

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1 Introduction

So far most discussions of supersymmetric (SUSY) phenomenology [1] including the searches for SUSY particles, conducted both at hadron [2] as well as e^+e^- colliders [3] have always *assumed* the existence of an exactly conserved R parity. This symmetry, imposed upon the theory, is assigned in such a way that all standard model particles (including the Higgs scalars) are R-even while their SUSY partners are R-odd. It acts like a selection rule that dictates that all SUSY particles can only be pair-produced, and the lightest must be stable. This then leads to a characteristic missing energy signature, that has so far been taken as the basis of SUSY searches. The assumption of R parity conservation may, in addition, have important cosmological implications.

R parity may well be violated, at some level, and it is therefore of great importance to pursue the phenomenological consequences of alternative scenarios. Because of its relation to lepton number, R parity violation is always, in a way or another, related or constrained by neutrino physics considerations [4].

The most predictive way of breaking R parity is spontaneously [5], through nonzero vacuum expectation values (VEVS) for scalar neutrinos, such as

$$v_L = \langle \tilde{\nu}_{L\tau} \rangle \tag{1}$$

$$v_R = \langle \tilde{\nu}_{R\tau} \rangle \tag{2}$$

The resulting R parity violating effects may be systematically investigated, since there are relatively few extra parameters.

If spontaneous R parity violation occurs in the absence of any additional gauge symmetry beyond $SU(2) \otimes U(1)$ [6], it leads to the existence of a physical massless Nambu-Goldstone boson - a Majoron [7,8], which is then *the lightest SUSY particle*. There are in this case important observational constraints that come from astrophysics [9], as well as from the recent precision measurement of the Z^0 width and SUSY searches at LEP [10,3]. The astrophysical constraint arises due to the fact that, if Majorons are produced in a stellar environment, they freely escape due to their tiny coupling to matter. These new mechanisms of stellar energy loss lead to a severe constraint on the Majoron coupling to the electron $g_{ee} \lesssim 2 \times 10^{-13}$ and, consequently, on the sneutrino VEV, $v_L \lesssim 30 KeV$. As a result, this model leads to a new decay mode for the neutral gauge boson

$$Z^0 \rightarrow \rho + J \tag{3}$$

where ρ is a light scalar, related to the Majoron, with mass of order v_L, the scale characterizing spontaneous R parity breaking. This gives rise to an increase in the invisible width of the Z^0 by $\Delta\Gamma_{Z^0}^{inv} \sim 85\ MeV$, a value now excluded by the recent precision measurement of the Z^0 width at LEP [10].

There is in this case just one way out, *i.e.*, to extend the minimal $SU(2) \otimes U(1)$ supersymmetric model, so that R parity breaking is driven by *isosinglet* lepton VEVS, such as eq. (2). Unlike the minimal broken R parity model, in this case the astrophysical bound can be obeyed

without the need of fine-tuning [11]. In addition, the decay in eq. (3) does not take place, since the Majoron, being mostly an isosinglet, does not couple to the Z^0. This avoids conflict with the LEP measurement of $\Gamma_{Z^0}^{inv}$.

There is, however, a *second* possible scenario for the idea of spontaneous R parity violation, and that is the one we investigate in this paper. It corresponds to the situation where there is a lepton number symmetry *contained in the gauge group*. The simplest model of this type is the one where there is just a single additional $U(1)$ symmetry. In this case there is no Majoron, since the would-be Goldstone boson is *eaten* and becomes the longitudinal mode of the corresponding Z' gauge boson. In this model the scale of R parity violation is the same that characterizes the new gauge interaction, presumably around the TeV scale. Consequently, its effects are potentially large.

Here we analyse the constraints on R parity breaking in this second model, as well as the resulting signals in Z^0 decays. The constraints include laboratory and cosmological neutrino physics constraints, as well as those from collider searches for SUSY particles, and precision Z^0 measurements. The signals in Z^0 decays include experimentally measurable rates for the decay

$$Z^0 \rightarrow \tau^+\chi^- + \tau^-\chi^+ \tag{4}$$

as well as for zen-events. These events are different in origin from those of the standard R parity conserving case [12], and may occur with enhanced rates.

2 The Model and Mass Matrices

We consider, for definiteness, the simplest model that illustrates our suggested scenario for spontaneous R parity violation. The model was considered in ref [13,14], and is characterized by the hypercharge quantum numbers given in table 1. We take the quark sector to be completely canonical, but the lepton sector is affected, in an important way, by R parity breaking. We assume the case where the isosinglet mirror quarks g and g^c decouple at a large intermediate scale. The resulting low energy superpotential is given as

$$h_u Q H_u u^c + h_d Q H_d d^c + h_\nu \ell H_u \nu^c + \lambda_{10} n H_u H_d + h_e H_d \ell e^c \tag{5}$$

where h_u, h_d, h_e, h_ν are all arbitrary matrices in generation space. For simplicity, we assume CP conservation, so these matrices are taken to be real.

Symmetry breaking proceeds in two steps. First, the $U(1)_\chi$ symmetry breaks, at a scale $v_R = O(few\ TeV)$, in eq. (2). The presence of the h_ν coupling in the superpotential, eq. (5), implies that lepton number as well as R parity are spontaneously broken by the nonzero VEV v_R in eq. (2). This way the scale of R parity violation is the same that characterizes the new gauge interaction.

The second stage of symmetry breaking is that of electroweak breaking, and occurs through

$$\langle H_u \rangle = v_u \tag{6}$$

$$\langle H_d \rangle = v_d \tag{7}$$

with $v^2 = v_u^2 + v_d^2$ fixed by the W mass. This also generates quark and lepton mass terms.

The mass matrix of the charged leptons, arising from D-terms and F superpotential couplings, is given by

$$\begin{array}{c|ccc} & e_j^+ & \tilde{H}_u^+ & -i\tilde{W}^+ \\ \hline e_i & h_{eij}v_d & -h_{\nu ij}v_{Rj} & \sqrt{2}g_2 v_{Li} \\ \tilde{H}_d^- & -h_{eij}v_{Li} & \mu & \sqrt{2}g_2 v_d \\ -i\tilde{W}^- & 0 & \sqrt{2}g_2 v_u & M_2 \end{array} \tag{8}$$

where $g_{1,2}$ are the $SU(2)\otimes U(1)$ gauge couplings divided by $\sqrt{2}$ and $M_{1,2}$ denote the supersymmetry breaking gaugino mass parameters, related by

$$M_1/M_2 = \frac{5}{3}tan^2\theta_W. \tag{9}$$

It is always possible, if CP is conserved in this sector, to choose $M_2 > 0$, while μ may have either sign. We have also chosen the basis so that h_e takes a diagonal form, and, for simplicity, taken $v_L \to 0$.

Similarly, the neutral lepton mass matrix is [‡]

$$\begin{array}{c|ccccccc} & \nu_j & \nu_j^c & \tilde{H}_u & \tilde{H}_d & -i\tilde{W}_3 & -i\tilde{B}_Y & -i\tilde{B}_{Y_R} \\ \hline \nu_i & 0 & h_{\nu ij}v_u & h_{\nu ik}v_{Rk} & 0 & g_2 v_{Li} & -g_1 v_{Li} & g_R v_{Li} X_L^i \\ \nu_i^c & h_{\nu ji}v_u & 0 & h_{\nu ik}v_{Lk} & 0 & 0 & 0 & g_R v_{Ri} X_R^i \\ \tilde{H}_u & h_{\nu jk}v_{Rk} & h_{\nu jk}v_{Lk} & 0 & -\mu & -g_2 v_u & g_1 v_u & g_R v_u X_u \\ \tilde{H}_d & 0 & 0 & -\mu & 0 & g_2 v_d & -g_1 v_d & g_R v_d X_d \\ -i\tilde{W}_3 & g_2 v_{Lj} & 0 & -g_2 v_u & g_2 v_d & M_2 & 0 & 0 \\ -i\tilde{B}_Y & -g_1 v_{Lj} & 0 & g_1 v_u & -g_1 v_d & 0 & M_1 & 0 \\ -i\tilde{B}_{Y_R} & g_R v_{Lj} X_L^j & g_R v_{Rj} X_R^j & g_R v_u X_u & g_R v_d X_d & 0 & 0 & M_R \end{array} \tag{10}$$

where the X_i are obtained from table 1. The form of these matrices determines not only the masses of the physical supersymmetric *inos*, but also their couplings and decay patterns.

For convenience we define the parameter β as

$$\tan\beta = \frac{v_u}{v_d} \tag{11}$$

In our analysis we fix a characteristic value for $\tan\beta$. In addition, we will make the following reasonable simplifying assumptions about the R parity violating VEVS

$$\begin{aligned} v_{Li} &= \langle \tilde{\nu}_i \rangle = 0 \quad \forall i \\ v_{Ri} &= \langle \tilde{\nu}_i^c \rangle = 0 \quad i = 1,2 \\ v_{R3} &= v_R \end{aligned} \tag{12}$$

[‡]Under our assumptions, the singlet state $\tilde{n}$ is massless and decouples from the others.

In addition, we will approximate the mass of the extra neutral gauge boson Z' by

$$M_{Z'} \simeq \frac{10}{3}\frac{e^2}{c_w^2}[\frac{1}{10}(v_u^2 + v_d^2) + \frac{5}{8}v_R^2] \tag{13}$$

which amounts to neglecting the effect due to the mixing with the standard Z^0 *i.e.*, $Z' \simeq Z_X$. This is a good approximation in this model, in view of the constraints that follow both from laboratory [15] as well as nucleosynthesis [16].

Finally, for illustration, we choose to work with a specific ansatz [17] for the h_ν matrix, given by

$$h_\nu = \begin{pmatrix} 0 & h_{13} & 0 \\ h_{13} & 0 & h_{23} \\ 0 & h_{23} & h_{33} \end{pmatrix} \tag{14}$$

In order to have an acceptably small value for the ν_e mass we will take $h_{13} \simeq 0$.

The diagonalizing matrices are defined through

$$\psi^0_{weak} = N^\dagger \psi^0_{mass} \qquad \psi^-_{weak} = U_L \psi^-_{mass} \qquad \psi^+_{weak} = U_R \psi^+_{mass} \tag{15}$$

in the notation of Weyl spinors. With this we get in the spectrum of neutral fermion states, 6 heavy states, plus 5 light states. These include:

1. Two massless states $\nu_1 \simeq \nu_e$ and $\nu_2 \simeq \nu_e^c$ (this follows from taking $h_{13} \simeq 0$).
2. Two very ligtht states $\nu_3 \simeq \nu_\mu$ and $\nu_4 \simeq \nu_\mu^c$ with $m_{\nu_3} < m_{\nu_4}$. We impose the constraints

$$m_{\nu_3} \lesssim 100\, eV \qquad m_{\nu_4} \lesssim 1\, KeV \tag{16}$$

which follow from requiring an acceptably small cosmological density of relic neutrinos [18,19].

3. One massive state $\nu_5 \simeq \nu_\tau$. For reasonable values of the SUSY parameters, the ν_τ mass can easily violate the cosmological limit on stable neutrinos [18]. Although there are no modes of invisible neutrino decay involving Majoron emission [20,21], it is possible for the corresponding decay lifetimes due to

$$\nu_\tau \to 3\nu \tag{17}$$

[22,8] or, possibly, $\nu_\tau \to \nu e^+ e^-$, to be shorter than required by cosmology so as to efficiently suppress the contribution of relic ν_τ decay products, for a wide range of experimentally allowed ν_τ mass values. We will study below the possible ν_τ decay channels and the cosmological constraints on the lifetime. For the moment we only need to impose the ARGUS limit [23]

$$m_{\nu_\tau} \leq 35\, MeV \tag{18}$$

As a result of large m_{ν_τ} values being allowed, the effects associated with R parity violation may be correspondingly enhanced.

3 Leptonic Z^0 and W decays

The *charged lepton* decays of the Z^0 include decays to both standard and supersymmetric fermions, as well as mixed, R parity violating, decays. The generic form of the Z^0 decay width to charged fermions is

$$\Gamma(Z^0 \to \chi_i^- \chi_j^+) = \frac{e^2 M_{Z^0}}{24\pi c_w^2 s_w^2}\{[1-(x_i+x_j)^2][1-(x_i-x_j)^2]\}^{\frac{1}{2}} \\ \left\{(O_{Lij}^2 + O_{Rij}^2)\left[1-\tfrac{1}{2}(x_i^2+x_j^2)-\tfrac{1}{2}(x_i^2-x_j^2)^2\right] + 6O_{Lij}O_{Rij}x_i x_j \eta_i \eta_j\right\} \tag{19}$$

where we have set

$$x_i = \frac{|m_i|}{M_{Z^0}} \qquad \eta_i = sign(m_i) \tag{20}$$

The η_i's are signs that are required in our choice of notation in which the (CP conserving) diagonalizing matrices are real, but the mass eigenvalues may be either positive and negative. The $O_{L,R}$ coefficients are given by

$$O_{Lij} = (-\tfrac{1}{2} + s_w^2)\delta_{ij} - \tfrac{1}{2}U_{Li5}^\dagger U_{L5j} \\ O_{Rij} = s_w^2 \delta_{ij} - \tfrac{1}{2}U_{Ri4}^\dagger U_{R4j} - U_{Ri5}^\dagger U_{R5j} \tag{21}$$

As a result of R parity breaking, there are new Z^0 decay channels, given in eq. (4), where we have denoted by $\chi = \chi_4$ the lightest chargino. Such decays arise due to mixing of the supersymmetric fermions (partners of gauge and Higgs particles) with the weak-eigenstate leptons implied by eq. (8).

Similarly, the *neutral lepton* decays of the Z^0 also include decays to both standard and supersymmetric fermions, as well as mixed decays. The generic form of the Z^0 decay width to neutral fermions is

$$\Gamma(Z^0 \to \chi_i^0 \chi_j^0) = \frac{(2-\delta_{ij})e^2 M_{Z^0}}{24\pi c_w^2 s_w^2}(O'_{Lij})^2\{[1-(x_i+x_j)^2][1-(x_i-x_j)^2]\}^{\frac{1}{2}} \\ \left\{1-\tfrac{1}{2}(x_i^2+x_j^2)-\tfrac{1}{2}(x_i^2-x_j^2)^2 - 3x_i x_j \eta_i \eta_j\right\} \tag{22}$$

with

$$O'_{Lij} = \frac{1}{2}P^{ij} = \frac{1}{2}(N_{i1}N_{1j}^\dagger + N_{i2}N_{2j}^\dagger + N_{31}N_{3j}^\dagger + N_{i8}N_{8j}^\dagger - N_{i7}N_{7j}^\dagger) \tag{23}$$

In this expression we have taken into account the Majorana properties of the neutralinos. In the range of the supersymmetric parameters we are restricting ourselves to, the kinematically open channels are just those corresponding to $i,j \leq 7$.

For the W we have the following nonstandard decays ($i \geq 4$ and $j \geq 6$)

$$\Gamma(W^- \to \chi_i^- \chi_j^0) = \frac{e^2 M_W}{12\pi s_w^2}\left\{\left[1-(x_i+x_j)^2\right]\left[1-(x_i-x_j)^2\right]\right\}^{\frac{1}{2}} \tag{24}$$

$$\left\{\frac{1}{2}((O''_{Lij})^2 + (O_{Rij}'')^2)\left[1-\frac{1}{2}(x_i^2+x_j^2)-\frac{1}{2}(x_i^2-x_j^2)^2\right] + 3O''_{Lij}O''_{Rij}x_i x_j \eta_i \eta_j\right\} \tag{25}$$

with

$$O''_{Lki} = -N_{k9}U_{L5i} - \frac{1}{\sqrt{2}}N_{k8}U_{L4i} - \frac{1}{\sqrt{2}}\sum_{m=1}^{3} N_{km}U_{Lmi} \\ O''_{Rki} = -N_{k9}^* U_{R5i}^* + \frac{1}{\sqrt{2}}N_{k8}^* U_{R4i}^* \tag{26}$$

and here $x_i = \frac{|m_i|}{M_W}$.

4 Neutrino and neutralino decays

In this section as is the previous one we will start denoting generically by χ_i^0 *all* of the neutral lepton states, including the 5 light states that we will also call neutrinos (although two of them have only a superweak interaction). In general for neutral leptons producible in Z^0 decays there are two possible decay channels [§]:

- $\chi_i^0 \to \chi_j^0 \chi_k^0 \chi_k^0$ $i = 2, 11,\ j = 1, i-1$ and $k = 1, 5$ with width

$$\Gamma(\chi_i^0 \to \chi_j^0 \chi_k^0 \chi_k^0) = \frac{G_F^2 m_i^5}{3 \times 2^6 \pi^3 (2-\delta_{jk})} \mid P_{ij} \mid^2 \mid P_{kk} \mid^2 F(x_i, x_j) \tag{27}$$

where we have taken the $m_k \ll M_Z^0$. $F(x_i, x_j)$ is a function that comes when integrating over the nontrivial phase space when the propagator of the Z^0 is taken into account, and its limit is

$$\lim_{x_i, x_j \to 0} F(x_i, x_j) = 1 \tag{28}$$

- $\chi_i^0 \to \chi_j^0 f \bar{f}$ $i = 2, 11,\ j = 1, i-1$

$$\Gamma(\chi_i^0 \to \chi_j^0 f \bar{f}) = \frac{G_F^2 m_i^5}{3 \times 2^4 \pi^3}({g_V^f}^2 + {g_A^f}^2) \mid P_{ij} \mid^2 F(x_i, x_j) \tag{29}$$

where f denotes any quark or charged lepton ($g_V^f = \frac{T_3^f}{2} - Q^f s_w^2$ and $g_A^f = \frac{T_3^f}{2}$).

Applying these general formulae for the case of the $\nu_\tau \simeq \chi_5^0$, and taking into account that $P_{22} \simeq P_{44} \simeq 0$, $P_{11} \simeq P_{33} \simeq 1$, $P_{15} = P_{25} = 0$ and $P_{35} \ll P_{45}$ then we have that if $m_5 \geq 1\,MeV$ there are two possible decay channels

$$\Gamma(\nu_5 \to 3\nu) = \frac{G_F^2 m_5^5}{192\pi^3} \mid P_{45} \mid^2 \\ \Gamma(\nu_5 \to \nu e^+ e^-) = \frac{G_F^2 m_5^5}{48\pi^3}({g_V^e}^2 + {g_A^e}^2) \mid P_{45} \mid^2 \\ \simeq \frac{1}{4}\frac{G_F^2 m_5^5}{192\pi^3} \mid P_{45} \mid^2 \tag{30}$$

[§]For simplicity we neglect the possible effect of scalar exchange in neutralino decays.

If $m_5 < 1\,MeV$ only the invisible channel is possible.

With all that we have

$$Br^{inv}(5) = \frac{\Gamma(\nu_5 \to 3\nu)}{\Gamma(\nu_5 \to \nu e^+ e^-) + \Gamma(\nu_5 \to 3\nu)} = \begin{cases} 1 & m_5 < 1\,MeV \\ \frac{4}{5} & m_5 \geq 1\,MeV \end{cases} \tag{31}$$

However, as far as the LEP phenomenology goes, for all intents and purposes, we can set $Br^{inv}(5) = 1$, since the ν_τ lifetimes are so large that no decays occur in the detector.

Note that in spontaneously broken R parity models the lightest neutralino is unstable, and has both *visible*

$$\chi^0 \to \nu_\tau + \bar{f} f \tag{32}$$

as well as *invisible* decays of the type

$$\chi^0 \to 3\nu \tag{33}$$

Here however, the fact the decays in eq. (32) and eq. (33) do occur normally inside the detector, implies that they are very important to take into account, and are a sizeable source of zen events, see below.

5 Theoretical Rates for Zen Events

In the minimal supersymmetric standard model these events arise from [12]

$$Z^0 \to \chi^0 + \chi^{0'} \tag{34}$$

where the lightest stable χ^0 escapes detection and $\chi^{0'}$ decays visibly as

$$\chi^{0'} \to \chi^0 + \bar{f} f \tag{35}$$

[¶] In most of these events, all the visible Z^0 decay products are contained in one hemisphere.

In the the present broken R parity model, the width of the Z^0 into zen events is given by

$$\Gamma(zen) = \sum_{i=1}^{11} \sum_{j=i}^{11} \Gamma(Z^0 \to \chi_i^0 \chi_j^0) \{ Br^{inv}(i)[1 - Br^{inv}(j)] + Br^{inv}(j)[1 - Br^{inv}(i)] \} \tag{36}$$

Note that here *the missing energy is always carried by neutrinos* and not by the "photino" as in the MSSM. As we have seen above

$$Br^{inv}(1) = Br^{inv}(2) = Br^{inv}(3) = Br^{inv}(4) = Br^{inv}(5) = 1 \tag{37}$$

since ν_τ decays only occur *outside* the detector.

[¶]Here we neglect the possible presence of radiative decays $\chi^{0'} \to \chi^0 + \gamma$.

For $i \geq 6$ one has

$$Br^{inv}(i) = \frac{\sum_{j=1}^{i-1} \sum_{k=1}^{5} \Gamma(\chi_i^0 \to \chi_j^0 \chi_k^0 \chi_k^0) Br^{inv}(j)}{\sum_{j=1}^{i-1} \sum_{k=1}^{5} \Gamma(\chi_i^0 \to \chi_j^0 \chi_k^0 \chi_k^0) + \sum_{j=1}^{i-1} \sum_{f\,charged} \Gamma(\chi_i^0 \to \chi_j^0 f \bar{f})} \tag{38}$$

From the expressions in the previous section we now define

$$\begin{aligned} f(j) &= \frac{\sum_{k=1}^{5} \Gamma(\chi_i^0 \to \chi_j^0 \chi_k^0 \chi_k^0)}{\Gamma(\chi_i^0 \to \chi_j^0 e \bar{e})} \\ &= \frac{3}{8(g_V^{e\,2}+g_A^{e\,2})}(1 - \delta_{j1} - \delta_{j3} - \delta_{j5}) + \frac{1}{2(g_V^{e\,2}+g_A^{e\,2})}(\delta_{j1} + \delta_{j3} + \delta_{j5}) \\ &= \frac{1}{8(g_V^{e\,2}+g_A^{e\,2})}(3 + \delta_{j1} + \delta_{j3} + \delta_{j5}) \end{aligned} \tag{39}$$

and

$$\frac{\Gamma(\chi_i^0 \to \chi_j^0 e \bar{e})}{\sum_{f\,charged} \Gamma(\chi_i^0 \to \chi_j^0 f \bar{f})} = \bar{\kappa} = 4.33 \times 10^{-2} \tag{40}$$

where the top quark is not included.

With all this we get

$$Br^{inv}(i) = \frac{\sum_{j=1}^{i-1} f(j) \Gamma(\chi_i^0 \to \chi_j^0 e \bar{e}) Br^{inv}(j)}{\sum_{j=1}^{i-1} [f(j) + \frac{1}{\bar{\kappa}}] \Gamma(\chi_i^0 \to \chi_j^0 e \bar{e})} \tag{41}$$

6 Constraints on the Parameters

The recent LEP data on Z^0 decays, including the limits recently extracted on the masses of some of the charged supersymmetric particles [3], as well as $\bar{p}p$ collider limits on $W^\pm, Z^0$ and gluino production constrain the parameters of the minimal supersymmetric extension of the standard model, with conserved R parity [24]. Similarly, they may be used to constrain spontaneously broken R parity models, as discussed in ref [25] in the case where spontaneous R parity violation is realized within the $SU(2) \otimes U(1)$ gauge structure.

Here we perform the corresponding analysis for the present case of interest. It differs, both qualitatively and quantitatively, from the analysis presented in ref [24] for the MSSM, as well as from that of ref [25] due to the absence of the Majoron, and the presence of the Z_χ gauge boson. In addition to collider constraints, we have also implemented those constraints characteristic of broken R parity models. They are related to neutrino mass considerations. We now list the constraints used in our present analysis:

1. The heavy sector of the charged matrix in eq. (8) leads to two heavy *charginos i.e.*, (charged supersymmetric partners of gauge and Higgs bosons). The lightest of these states, $\chi_4 \equiv \chi^{\pm}$, has not been pair produced in Z^0 decays at LEP [3]. If $m_{\tilde{\ell}} > m_{\chi^{\pm}}$ the lightest chargino would decay leptonically, via the charged current, as

$$\chi^{\pm} \rightarrow \chi^0 \ell^{\pm} \nu_\tau \tag{42}$$

or semileptonically, *i.e.*, $\chi^{\pm} \rightarrow \chi^0 \bar{q} q$, when the virtual W decays hadronically. This leads to chargino mass limits

$$m_{\chi_4} \geq 45\, GeV \tag{43}$$

Note however, that there is a range in mass where the scalar leptons would not be pair produced in Z^0 decays, but would still be light enough to be produced in the decays of charginos or neutralinos, produced in Z^0 decays, *i.e.*, in the mass range between eq. (43) and

$$m_{\chi_4} \leq 90\, GeV \tag{44}$$

This is an interesting range of the broken R parity model that can be probed in Z^0 decays at LEP.

2. $p\bar{p}$ collider limits on the R ratio, $\sigma_{W^{\pm}} B(W^{\pm} \rightarrow e^{\pm}\nu)/\sigma_{Z^0} B(Z^0 \rightarrow e^+e^-)$ [26].

$$0.825 \leq \frac{R}{R_{SM}} \leq 1.091 \tag{45}$$

in the present model

$$\frac{R}{R_{SM}} = \frac{\Gamma_{Z^0}}{\Gamma_W}\frac{\Gamma_W^{SM}}{\Gamma_{Z^0}^{SM}} = \frac{\Gamma_W^{SM}}{\Gamma_{Z^0}^{SM}} \frac{\Gamma_{Z^0}^{SM} + \sum_{i,j=4,5} \Gamma(Z^0 \rightarrow \chi_i^+ \chi_j^-) + \sum_{i=6}^{11}\sum_{j=i}^{11} \Gamma(Z^0 \rightarrow \chi_i^0 \chi_j^0)}{\Gamma_W^{SM} + \sum_{i=4}^{5}\sum_{j=6}^{11} \Gamma(W \rightarrow \chi_i^+ \chi_j^0)} \tag{46}$$

3. LEP limits on the total Z^0 width [10].

$$\Delta\Gamma_{Z^0}^{total} \leq 64\, MeV \tag{47}$$

in this model

$$\Delta\Gamma_{Z^0}^{total} = \sum_{i,j=4,5} \Gamma(Z^0 \rightarrow \chi_i^+ \chi_j^-) + \sum_{i=6}^{11}\sum_{j=i}^{11} \Gamma(Z^0 \rightarrow \chi_i^0 \chi_j^0) \tag{48}$$

Here we made the reasonable assumption that all supersymmetric scalar particles are too heavy to be pair-produced in Z^0 decay. We already know that this holds for the charged ones, at least.

4. LEP limits on the invisible Z^0 width [10].

$$\Gamma_{Z^0}^{inv} = 496 \pm 18\, MeV \Longrightarrow \Gamma_{Z^0}^{inv} < 525\, MeV (95\%) \tag{49}$$

In this model

$$\Gamma_{Z^0}^{inv} = \sum_{i=1}^{11}\sum_{j=i}^{11} \Gamma(Z^0 \rightarrow \chi_i^0 \chi_j^0) Br^{inv}(i) Br^{inv}(j) \tag{50}$$

5. The CDF lower limit on the gluino mass $m_{\tilde{g}}$ [27] which restricts soft supersymmetry breaking electroweak gaugino mass parameter

$$M_2 > 20\, GeV \tag{51}$$

6. LEP constraints on the peak cross section versus total Z^0 width, which lead to an allowed ellipsis $(\sigma_0^{had}, \Gamma_{Z^0}^{total})$ [28]. If we know $\sigma_0^{had} \pm E\sigma_0^{had}$, $\Gamma_{Z^0}^{total} \pm E\Gamma_{Z^0}^{total}$ from the experiment and the correlation, ρ, between them, then the condition on any theoretical value $\sigma_0^{had^T}$ $\Gamma_{Z^0}^{total^T}$ to be in the 90% CL ellipsis is

$$\frac{(\sigma_0^{had^T} - \sigma_0^{had})^2}{E\sigma_0^{had^2}} - 2\rho \frac{(\sigma_0^{had^T} - \sigma_0^{had})(\Gamma_{Z^0}^{total^T} - \Gamma_{Z^0}^{total})}{E\sigma_0^{had} E\Gamma_{Z^0}^{total}} + \frac{(\Gamma_{Z^0}^{total^T} - \Gamma_{Z^0}^{total})^2}{E\Gamma_{Z^0}^{total^2}} \leq 4.61 \tag{52}$$

In the present model the modifications on $\Gamma_{Z^0}^{total}$ have been given above and σ_0^{had} is given by

$$\sigma_0^{had^T} \simeq 3.88 \times 10^5 \frac{12\pi}{M_{Z^0}^2} \Gamma_{ee} \frac{\Gamma_{had}}{(\Gamma^{total})^2} (nb) \tag{53}$$

where

$$\Gamma_{had} = \Gamma_{had}^{SM} + \sum_{i=6}^{11}\sum_{j=i}^{11} \Gamma(Z^0 \rightarrow \chi_i^0 \chi_j^0) \left\{ Br^{had}(i) + [1 - Br^{had}(i)] Br^{had}(j) \right\} \tag{54}$$

$$Br^{had}(i) = \frac{\sum_{j=1}^{i-1} \Gamma(\chi_i^0 \rightarrow \chi_j^0 q\bar{q})[1 - Br^{had}(j)] + \sum_{j=1}^{i-1} \Gamma(\chi_i^0 \rightarrow \chi_j^0\, all) Br^{had}(k)}{\sum_{j=1}^{i-1} \Gamma(\chi_i^0 \rightarrow \chi_j^0\, all)} \tag{55}$$

Taking into account what we have seen in the previous section and

$$\frac{\Gamma(\chi_i^0 \rightarrow \chi_j^0 q\bar{q})}{\Gamma(\chi_i^0 \rightarrow \chi_j^0 e\bar{e})} \simeq 20 \tag{56}$$

$$\frac{\Gamma(\chi_i^0 \rightarrow \chi_j^0\, all)}{\Gamma(\chi_i^0 \rightarrow \chi_j^0 e\bar{e})} = [f(j) + \frac{1}{\kappa}] \tag{57}$$

then

$$Br^{had}(i) = \frac{\sum_{j=1}^{i-1}\left\{20[1 - Br^{had}(j)] + [f(j) + \frac{1}{\bar{\kappa}}]Br^{had}(j)\right\}\Gamma(\chi_i^0 \to \chi_j^0 e\bar{e})}{\sum_{j=1}^{i-1}[f(j) + \frac{1}{\bar{\kappa}}]\Gamma(\chi_i^0 \to \chi_j^0 e\bar{e})} \qquad (58)$$

7. The existence of fermion states that cannot be kinematically produced in low energy weak decays changes the relative rates of various such processes, such as beta or μ decays [29] and leads to universality violations. The resulting constraints, that we adapt from ref [30], may be given as follows

$$\sum_{i=6}^{11}(N\,U_L)_{i1}(U_L^\dagger N^\dagger)_{1i} \lesssim 0.043$$
$$\sum_{i=6}^{11}(N\,U_L)_{i2}(U_L^\dagger N^\dagger)_{2i} \lesssim 0.008 \qquad (59)$$
$$\sum_{i=6}^{11}(N\,U_L)_{i3}(U_L^\dagger N^\dagger)_{3i} \lesssim 0.1$$

8. Neutrino oscillation data and direct searches for anomalous peaks on the energy distribution of the electrons and muons coming from the decays $\pi \to e\nu$ [31] [32] [33], $\pi \to \mu\nu$ [34] and $K \to \mu\nu$[35] lead to constraints on the mixing among the light neutral states.

9. The ARGUS limit on the τ neutrino mass

$$m_{\nu_\tau} \leq 35\, MeV \qquad (60)$$

10. The cosmological ν_τ lifetime limit [19]:

$$\tau \lesssim 1.5 \times 10^7 yr \left(\frac{m_5}{KeV}\right)^{-2} \qquad (61)$$

which gives

$$|\, P_{45}\,|^2 \left(\frac{m_5}{KeV}\right)^3 \gtrsim 6 \times 10^4 \qquad (62)$$

for $m_{\nu_\tau} \lesssim few\, MeV$. For larger ν_τ masses the limit is weaker, due to the Boltzmann suppression. The typical minimal value of the ν_τ mass for which the 3ν decay mode is efficient is around $\sim 300\, KeV$.

We have determined the values of the broken R parity model parameters that are consistent with all of the constraints discussed above. For definiteness we fix the values $M_{Z'} = 1\, TeV$ and take $v_L \to 0$. In fig 1,2,3 we show the regions in the M_2, μ plane allowed by all of the relevant constraints, for different values of $\tan\beta$ and $h_{\nu_\tau} \equiv h_{\nu 33}$, in each case. For a given β, the region is smaller for larger h_{ν_τ}, illustrating the important role by the ARGUS constraint, eq. (60).

7 Signatures in Z^0 Decays

Having mapped out the relevant allowed parameter space, we now study some of the phenomenological implications of the spontaneously broken R parity model. We focus on two signatures of supersymmetry in these models that could be experimentally measurable at LEP.

7.1 The decay $Z^0 \to \tau^+\chi^- + \tau^-\chi^+$

When R parity is spontaneously broken, SUSY fermions can be singly-produced in Z^0 decays, and this may be taken as a characteristic signal of these models.

For example, the study we have made in the present model, reveals that the new Z^0 decay given in eq. (4) may occur with a branching ratio that may be probed at LEP. Similar result was also found in the model of ref. [25].

We have plotted the regions where $BR(Z^0 \to \chi^\pm\tau^\pm) \geq 10^{-6}\,(A)10^{-5}\,(B)$ in fig 4, 5,6 for different values of of $\tan\beta$, with qualitatively very similar results. The effects of R parity breaking are then controlled by the magnitude of the τ neutrino mass, determined from eq. (10). To display this relationship we have let all parameters vary randomly within the region allowed by the observational constraints discussed above and determined the regions where $BR(Z^0 \to \chi^\pm\tau^\pm) \geq 10^{-6}(A)10^{-5}(B)$ directly in terms of the physical parameters m_χ and m_{ν_τ}. Our results are plotted in fig 7, corresponding to $\tan\beta = 3$. The results are very insensitive to the value of β. This figure synthesises the key elements of the broken R parity model and displays the direct correlation between the strength of R parity violating phenomena and the magnitude of the neutrino mass.

7.2 Zen Events

In the present model one expects many events with a substantial amount of missing energy, *carried by neutrinos*, and *not* by the lightest (*unstable*) neutralino. *Zen events*, similar to those suggested in ref. [12], are also a characteristic feature of the spontaneously broken R parity model. However, their main origin is now the process

$$Z^0 \to \chi^0\chi^0 \qquad (63)$$

with χ^0 decaying visibly in one hemisphere, and invisibly in the other. This contrasts with the origin of the zen events in the MSSM. Their corresponding theoretical rates were given above. They could also be *enhanced with respect to those of the MSSM*. The results of our study are shown in fig 8,9,10. We see that one may obtain a $BR(zen) \geq 10^{-4}$ for a wide range of the supersymmetric parameters. Detailed parameter assumptions are explained in the corresponding figure captions.

In summary, we suggested a variant of supersymmetry where R parity is a spontaneously broken symmetry and evaluated its predictions for Z^0 decays. The results given above are all consistent with the existing observational constraints. In addition, we also expect signals of R parity breaking at hadron colliders.

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	T_3	$\sqrt{\frac{5}{3}}Y$	$\sqrt{40}Y_X$
Q	$\begin{pmatrix}\frac{1}{2}\\-\frac{1}{2}\end{pmatrix}$	$\frac{1}{6}$	-1
u^c	0	$-\frac{2}{3}$	-1
e^c	0	1	-1
d^c	0	$\frac{1}{3}$	3
ℓ	$\begin{pmatrix}\frac{1}{2}\\-\frac{1}{2}\end{pmatrix}$	$-\frac{1}{2}$	3
H_d	$\begin{pmatrix}\frac{1}{2}\\-\frac{1}{2}\end{pmatrix}$	$-\frac{1}{2}$	-2
g^c	0	$\frac{1}{3}$	-2
H_u	$\begin{pmatrix}\frac{1}{2}\\-\frac{1}{2}\end{pmatrix}$	$\frac{1}{2}$	2
g	0	$-\frac{1}{3}$	2
ν^c	0	0	-5
n	0	0	0

Table 1: Quantum Numbers of the particles in the **27** of E_6 with respect to the gauge group $SU(2)_L \otimes U(1)_Y \otimes U(1)_X$.

Figure Captions

Fig 1:
Presently allowed regions of the supersymmetric parameters accessible to LEP searches in the spontaneously broken R parity model defined in table 1, for different values of h_{ν_τ}: (A) is for $h_{\nu_\tau} = 6.3 \times 10^{-4}$ and (B) is for $h_{\nu_\tau} = 3 \times 10^{-3}$. We have assumed as input values $M_{Z'} = 1\,TeV$, and $\tan\beta = 3$.

Fig 2:
Same as fig 1 for $\tan\beta = 10$ and h_{ν_τ} replaced by $h_{\nu_\tau} = 10^{-2}$ in case (B) and $h_{\nu_\tau} = 2 \times 10^{-3}$ in case (A).

Fig 3:
Same as fig 1 for $\tan\beta = 20$ and h_{ν_τ} replaced by and $h_{\nu_\tau} = 4\times10^{-3}$ in case (A) and $h_{\nu_\tau} = 2\times10^{-2}$ in case (B) .

Fig 4:
Presently allowed regions of the supersymmetric parameters leading to $BR(Z^0 \to \chi\tau) \geq 10^{-6}$ (A) and $BR(Z^0 \to \chi\tau) \geq 10^{-5}$ (B) We have assumed as input values $M_{Z'} = 1\,TeV$, and $\tan\beta = 3$.

Fig 5:

Fig 7:
Presently allowed regions of the supersymmetric parameters leading to $BR(Z^0 \to \chi\tau) \geq 10^{-6}$ (A) and $BR(Z^0 \to \chi\tau) \geq 10^{-5}$ (B), as a function of the physical mass of the lightest chargino and that of the τ neutrino. There is a direct correlation between the strength of R parity violating phenomena and the magnitude of the τ neutrino mass. We have assumed as input values $M_{Z'} = 1\,TeV$, and $\tan\beta = 3$.

Fig 8:
Presently allowed regions of the supersymmetric parameters leading to $BR(zen) \geq 10^{-4}$ for different values of h_{ν_τ}: (A) is for $h_{\nu_\tau} = 6.3 \times 10^{-4}$ and (B) is for $h_{\nu_\tau} = 3 \times 10^{-3}$. We have assumed as input values $M_{Z'} = 1\,TeV$, and $\tan\beta = 3$.

Fig 9:
Same as fig 8 for $\tan\beta = 10$ and h_{ν_τ} replaced by $h_{\nu_\tau} = 10^{-2}$ in case (B) and $h_{\nu_\tau} = 2 \times 10^{-3}$ in case (A).

Fig 10:
Same as fig 8 for $\tan\beta = 20$ and h_{ν_τ} replaced by and $h_{\nu_\tau} = 4\times10^{-3}$ in case (A) and $h_{\nu_\tau} = 2\times10^{-2}$ in case (B) .

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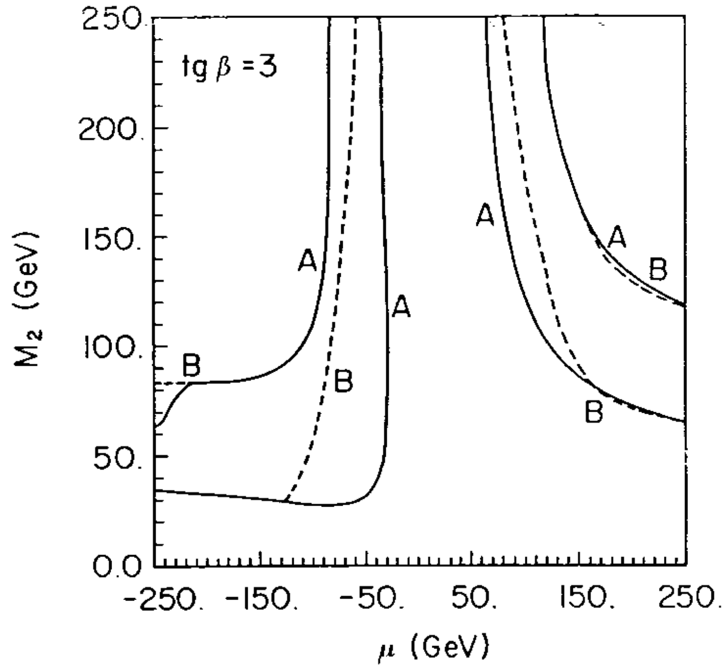
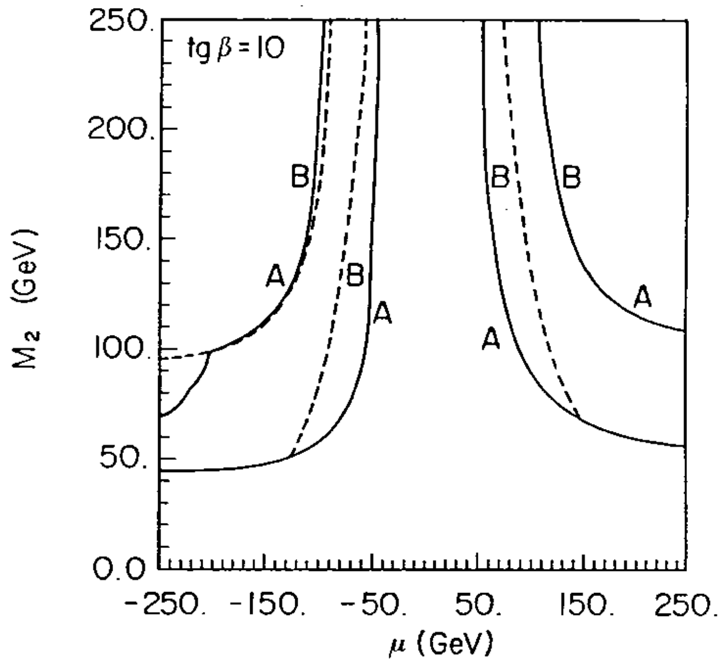


Fig. 1



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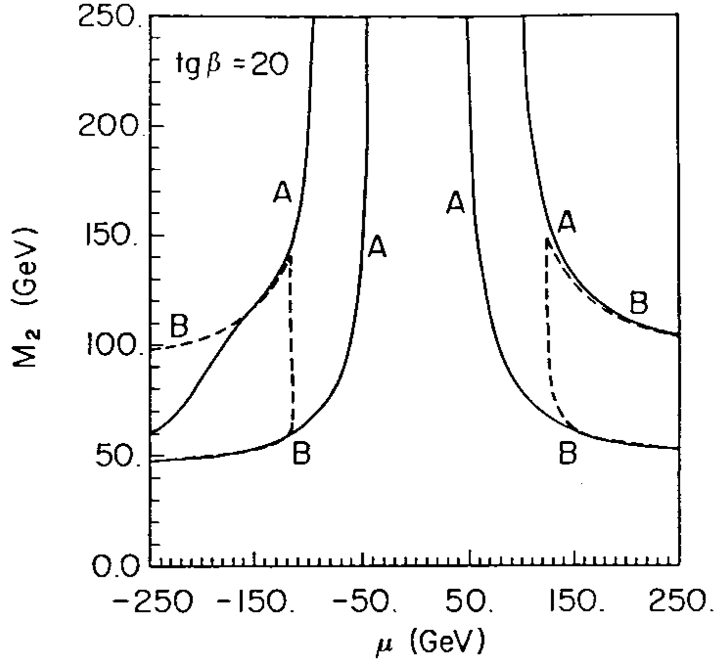

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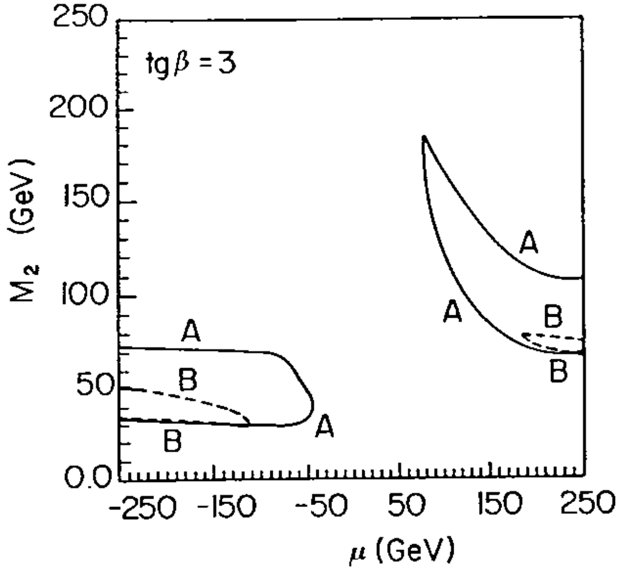

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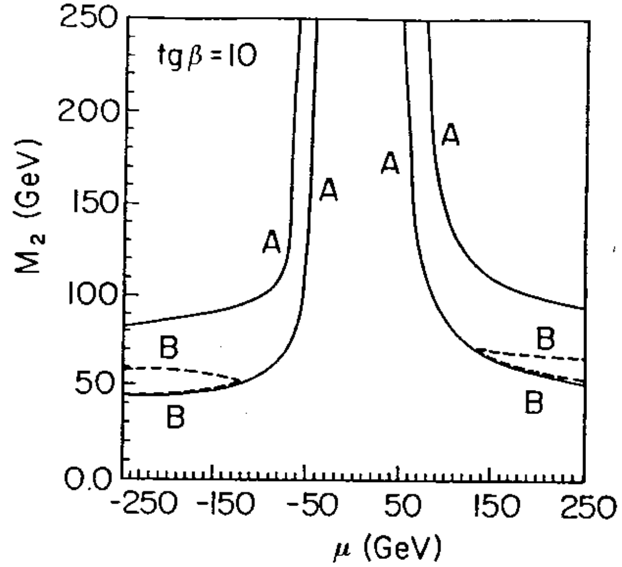

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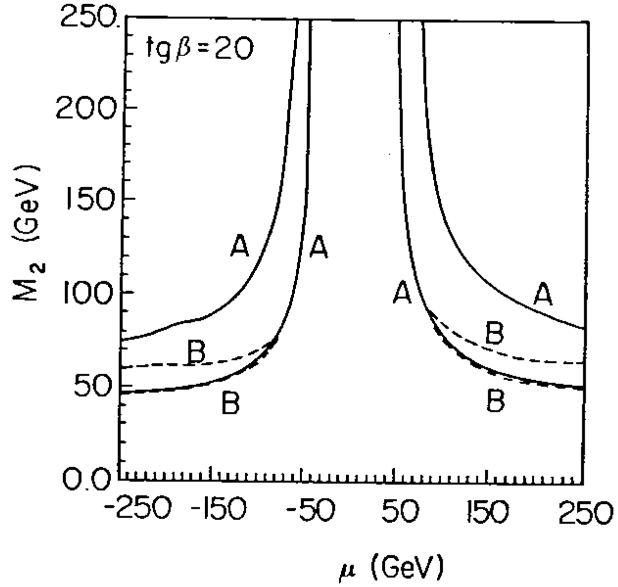

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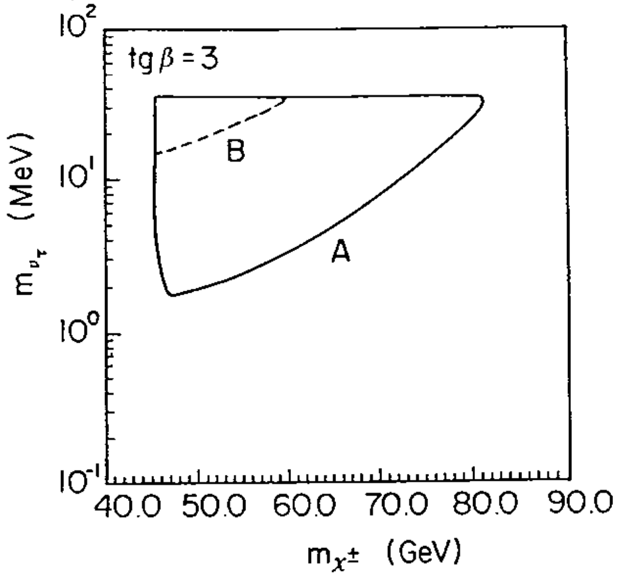

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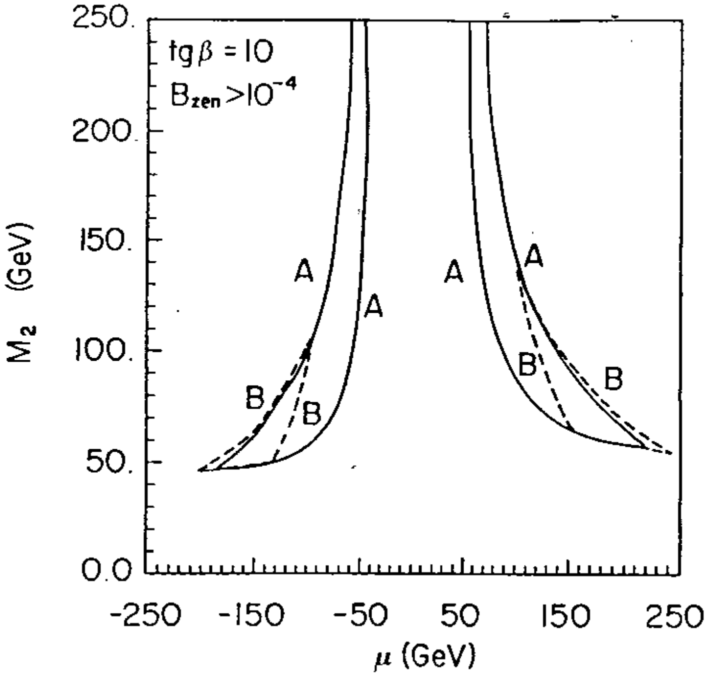

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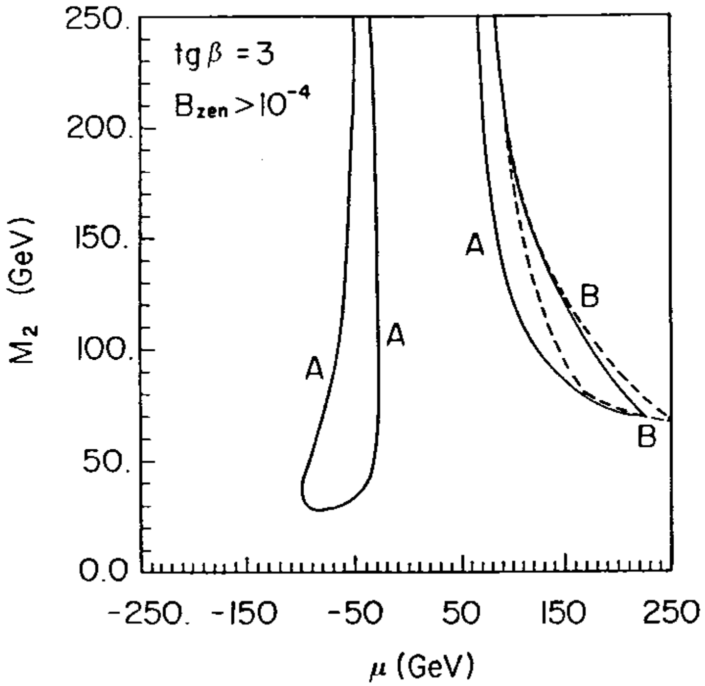

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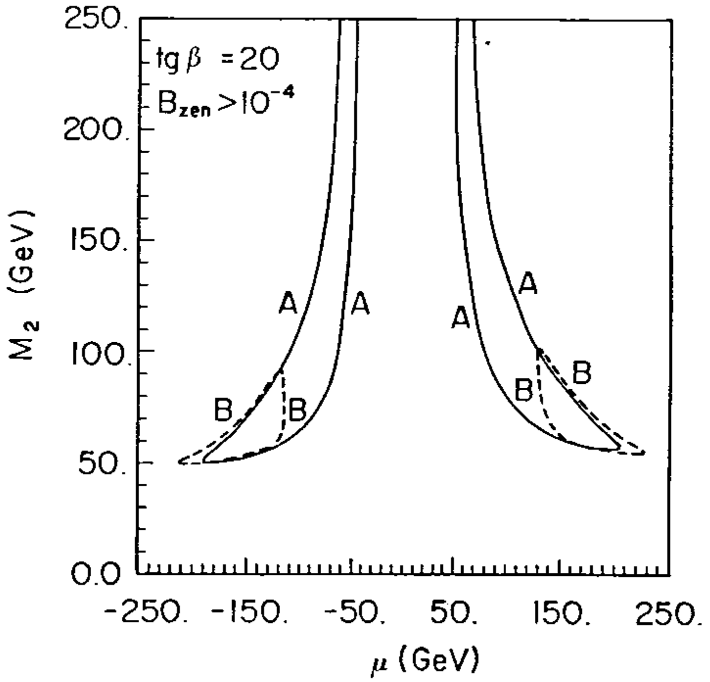

Fig. 10

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