Radiative $W$ and $Z$ Decays and Spontaneous R-Parity Violation

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Abstract

We point out that in a class of supersymmetric models where R-parity violation is induced by the spontaneous breaking of local $B - L$ symmetry, the R-parity violating $W$ decay $W \to \tilde{\ell} \gamma$ and $Z$ decay $Z \to \tilde{\nu} \gamma$, forbidden in the minimal supersymmetric standard model (MSSM), occur at an enhanced rate compared to other models with R-parity breaking. We find that the branching fractions for these modes can be of order $10^{-5}$. 
I. Introduction

Supersymmetrization of the standard model brings along with it the unpleasant feature that baryon and lepton numbers \((B\) and \(L\)) are no longer automatic symmetries of the Lagrangian. It is therefore customary to impose these symmetries on the model in order to avoid rapid proton decay or lepton number violation which are not yet observed in nature. Both these symmetries are however simultaneously obeyed if a discrete R-parity symmetry defined as \((-1)^{2S+3B+L}\) is imposed on the Lagrangian. The additional assumption of R-parity conservation severely limits the possible interactions among the fermions and their superpartners in the minimal supersymmetric standard model (MSSM). It not only implies that all superpartners of standard model particles must be produced in pairs but also that the lightest superparticle (the neutralino) must be stable. As mentioned above there is no a priori theoretical reason for either of these commonly made assumptions to hold.

Supersymmetric theories without R-parity conservation were introduced nearly ten years ago\([2,3]\) in order to examine the experimental constraints on the extent of departure from exact R-conservation. Two classes of theories were considered: one, where the R-parity violation is spontaneous\([2]\) and a second, where it is explicit in the original superpotential\([3]\). This latter possibility arises since the symmetries of the conventional gauge interactions of the MSSM do not \textit{a priori} forbid such terms in the superpotential. Many implications and tests of these two ideas have been subsequently analyzed in literature\([4]\).

It has recently been noted\([5,6]\) that if the sleptons are lighter than the \(W\) and \(Z\) bosons then R-parity violation implies a new decay channel for the latter, \(W \rightarrow \tilde{l}\gamma\) and \(Z \rightarrow \tilde{\nu}\gamma\), with the sleptons possibly decaying subsequently to two quarks via the mediation of R-violating interactions. These new decay modes of \(W\) and \(Z\) would be detectable in collider experiments provided the corresponding branching fractions are significantly large. It
turns out, however [5,6], that in most models with explicit R-parity breaking these branching fractions are found to be of order $10^{-7}$ or $10^{-8}$ keeping them beyond the reach of present experiments.

In this note, we focus our attention on a different class of models for R-violation. It was noted some time ago [7] that in extensions of the supersymmetric standard model where the gauge symmetry contains local $B - L$ as an explicit subgroup, R-parity invariance is automatic. Its violation therefore can emerge if the spontaneous breaking of $B - L$ is caused by a non-zero vev for the right-handed sneutrino (i.e., $< \tilde{\nu}^c > \neq 0$). An additional advantage of this model is that while lepton number violating terms are induced after spontaneous breaking, baryon number remains an exact symmetry, thereby avoiding any danger of rapid proton decay. It is the purpose of this note to point out that the single photon radiative $W$ and $Z$ decays are enhanced by two to three orders of magnitude in these models compared to the explicit R-parity violating scenarios and there is therefore a chance that they may be observable in collider experiments.

II. Details of the Model

The simplest gauge group which contains the standard model as well as a $U(1)_{B-L}$ factor is $SU(2)_L \times U(1)_{I_{SR}} \times U(1)_{B-L}$ which is itself in turn a subgroup of the left-right symmetric gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. We will illustrate our point using the first gauge group which is simpler to analyze although our results hold as well for the left-right symmetric SUSY model. The matter spectrum of the model consists of three generations of quarks and leptons as in the standard model plus three right-handed neutrinos (denoted by $\nu^c$). Their assignments under the extended gauge group are given in Table I.
The superpotential for this model can be written as follows:

\[ W = h_l LH_d e^c + h_\nu LH_u \nu^c + h_u QH_u u^c + h_d QH_d d^c + \mu H_u H_d \]  

(1)

In this equation, we have suppressed the generation indices. Before discussing the R-parity violation, let us comment briefly on the status of neutrino masses in this model. It is easy to see that the \( h_\nu \) terms in the superpotential will induce Dirac masses for the neutrinos which could \textit{a priori} be of the order of the charged lepton masses. The simplest way to cure this problem is to introduce two \( B - L \) carrying weak isosinglet fields, \( \Delta(1, +1, -2) \) \( \text{ and } \bar{\Delta}(1, -1, +2) \), and let them acquire nonzero vev's of the order a TeV or more. Then a new coupling in the superpotential of the form \( \nu^c \nu^c \Delta \) will give rise to the familiar seesaw mechanism for the neutrino making their masses small. Without effecting our final conclusions, we adopt a simpler approach without the extra \( \Delta \) fields. In order to understand the small neutrino masses in this simplified model, we will set all the elements of the coupling matrix \( h_\nu \) to zero except \( h_{\nu_33} \), so that both \( \nu_e \) and \( \nu_\mu \) have zero Dirac masses. The Dirac mass induced for the tau neutrino is then of order of the tau lepton mass itself; the smallness
of the $\nu_\tau$ mass can then be understood by a 3x3 see-saw mechanism as introduced in Ref.[8] after $\nu_{33}^c$ acquires a nonzero vev. To see this, let us assume that the $B - L$ symmetry breaking is implemented by $< \nu^c > = v_R$ and the rest of the symmetry breaking is caused by $< H_u > = \kappa_u$ and $< H_d > = \kappa_d$. It is then easily seen that the $\nu_{\tau}^c$ mixes with the linear combination of the gauginos corresponding to the generators $I_{3R}$ and $B - L$. Let us call this gaugino $\lambda_{\nu_{\tau}}$. Then the 3x3 mass matrix corresponding to the $\nu_\tau, \nu_\tau^c, \lambda_{\nu_{\tau}}$ basis appears as follows:

$$
\begin{pmatrix}
0 & h_\nu \kappa_u & 0 \\
h_\nu \kappa_u & 0 & \tilde{g} v_R \\
0 & \tilde{g} v_R & m_\lambda
\end{pmatrix}
$$

The 33-entry in the above matrix represents the SUSY breaking gaugino mass term. As was already remarked in Ref.8, this leads to a tiny mass for the left handed neutrino given by $m_\nu \simeq \left( \frac{m_D^2 m_\lambda}{\tilde{g}^2 v_R^2} \right)$, where $m_D = h_\nu \kappa_u$ and $\tilde{g}^2 = (g_R^2 + g_{BL}^2)/4$. This double see-saw result shows that in the expression for the neutrino mass there is an additional suppression coming from the Majorana mass of the gaugino. In supersymmetric theories, if this Majorana mass of the gaugino is set to zero at the tree level it can only arise at the two loop level and therefore can be less than a GeV. We then find that for $m_D \simeq 1 GeV$, $m_\lambda \simeq 1 GeV$ and $v_R \simeq 10 TeV$ we obtain a value for $m_{\nu_\tau} \leq 10 eV$ which is acceptable from both laboratory and cosmological considerations. Note further that, as $m_\lambda \to 0$, the tau neutrino mass vanishes.

It is also worth pointing out that if the other entries in the coupling matrix $h_\nu$ (entries other than the 33 entry) are nonzero, one could either invoke the usual see-saw mechanism by introducing the $\Delta$ and $\bar{\Delta}$ as mentioned before, or introduce three gauged $U(1)_{B-L}$ symmetries, one per generation, and give vevs to all three $\tilde{\nu}^c$'s. In the latter case we will have a generalized see-saw mechanism operating separately within each generation.
We will not elaborate on this possibility here. In either case, the main result of our paper remains unchanged.

When $\tilde{\nu}^c$ acquires a nonzero vev, it leads to R-parity violating interactions in the Lagrangian below the scale $v_R$. It is easily seen that such terms only induce lepton violating terms in the low energy Lagrangian. By making the off diagonal terms of $h_l$ arbitrarily small, we can make the theory consistent with the observed bounds on lepton number conservation. The main goal of our paper, i.e., to demonstrate that in this class of models the radiative $W$ and $Z$ decays are enhanced compared to other R-parity breaking models, follows directly from this approach.

Breaking local $B - L$ symmetry by the mechanism employed in this paper however implies other constraints on the model which have relevance to the strength of the radiative decays considered. To see this, note that a nonzero $h_\nu$ coupling combined with nonzero $v_R$ and $\kappa_u$ imply a nonzero vev for the tau sneutrino, i.e., $<\tilde{\nu}_\tau> = v_L$. In models with global $B - L$ symmetry breaking by $\tilde{\nu}^c$, there exists a massless Majoron and astrophysical constraints on the Majoron’s properties then imply that $v_L < 10 - 100 MeV$. On the other hand, in our model there is no massless Majoron. Thus the only constraint on $v_L$ comes from the fact that it leads to mixing between the tau neutrino and the $B - L$ gaugino as well as between the tau lepton and the wino. This mixing can cause departures from the apparent universality of leptonic decays of the muon and the tau leptons. These universality constraints are however weaker than those which apply in the case of the Majoron and are easily satisfied for $v_L$ less than a few GeV. As a conservative upper limit for $v_L$, we will choose a value of 1 GeV. In this case, the tau-gaugino mixing angle is less than 0.01. To see what constraints are implied by this, we note that

$$v_L \simeq h_\nu v_R \left( \frac{\mu \kappa_d + m_3 \kappa_u}{M^2} \right)$$  \hspace{1cm} (2)
For $v_R \simeq 10T eV$, $M \simeq 100GeV$, $\kappa_u, \kappa_d \simeq 100GeV$, $m_\pm \simeq 100GeV$ and $h_\nu \simeq 10^{-2}$, the above bound on $v_L$ is seen to be easily satisfied. Clearly as the constraints on $v_L$ improve, some of the parameters on the right-hand side of Eq.2 will become smaller. We will see that the strength of the radiative $W$ and $Z$ decays will depend on the values of these same parameters.

Let us briefly examine the charged and neutral scalar bosons in this theory. In the unitary gauge, there are the sleptons of the electron and muon type, which we assume are unmixed with the usual Higgs bosons. Due to the fact that $<\tilde{\nu}_\tau^c> \neq 0$, the tau slepton mixes with both $H_u$ and $H_d$ with the result being that we have two charged physical scalar bosons each of which will contain an admixture of the $\tilde{\tau}^+$. In the MSSM, where the tau slepton does not mix with the Higgs bosons, the physical charged Higgs boson is almost always heavier than the $W$ boson so that the $W$ cannot decay into it. In our model, there are two new ways in which the charged Higgs boson sector is different. First, the usual charged scalar can mix with the $\tilde{\tau}^+$ as already mentioned above and, secondly, the mass of $H_u^+$ gets additional contributions from the nonzero vev of $\tilde{\nu}^c$. Because of this, the physical eigenstate with the larger mass is predominantly the MSSM Higgs boson while the other eigenstate is predominantly the tau slepton which can be lighter than the $W$.

As we will see below, this lighter eigenstate will couple predominantly to the $t$ and $b$ quarks. Similarly, there will also be a neutral Higgs boson which will be predominantly the tau sneutrino and it will also have significant couplings to the $t$ and $b$ quarks. We will assume that its mass is also less than the mass of the $Z$ boson. These two particles can then appear in the radiative decay of the $W$ and $Z$.

III. Calculation of The Single Photon Radiative Decay of $W$ and $Z$

Now we are ready to discuss the induced lepton number violating term that is respon-
sible for the single photon radiative decays of the massive gauge bosons. The tree diagram shown in Fig.1 which arises after supersymmetry breaking leads after $B-L$ breaking to the effective $tb\bar{\nu}$ interaction of the form

$$\mathcal{L} = f [ \bar{t}_L b(1 + \gamma_5) t + \bar{b}_L \bar{t}(1 + \gamma_5) t ] + h.c. \quad (3)$$

where $f$ is given in our model by

$$f = \left( \frac{h_t h_v m_{\tilde{\nu} R}}{M_{H_u}^2} \right) \quad (4)$$

As mentioned before, we have a final state scalar that will be a linear combination of the charged Higgs fields which is predominantly the tau slepton.

If we choose $M_{H_u} \approx 100 GeV$, and all other parameters as discussed above, we see that $f$ can easily be of order one to three. We will assume this in making the predictions for the radiative branching ratios given below. The coupling in Eq.3 induces the decay amplitude for $W \to \tilde{\tau}\gamma$ via the one loop diagram in Fig.2. With this normalization, the amplitude for the $W \to \tilde{\tau}\gamma$ decay process can be written as

$$\mathcal{A} = \left[ F_1 \left( \frac{q_\nu k_\mu - g_{\nu\mu} k \cdot q}{M_W^2} \right) + iF_2 \epsilon_{\mu\nu\sigma\tau} q^\sigma k^\tau \right] \epsilon^\mu \epsilon^\nu \quad (5)$$

with $q(k)$ being the momentum of the photon($W$). In terms of the form factors $F_{1,2}$, the decay width is given by

$$\Gamma(W \to \tilde{\tau}\gamma) = \frac{M_W^3}{96\pi^2} (F_1^2 + F_2^2) \left(1 - \frac{m_\tilde{\tau}^2}{M_W^2}\right)^3 \quad (6)$$

with $m_\tilde{\tau}$ being the slepton mass. Defining the mass difference, $\delta = m_\tilde{\tau}^2 - M_W^2$, we find that $F_{1,2}$ can be written as

$$F_1 = \frac{-ieqN_c m_\nu L}{4\sqrt{2}\pi^2 \delta} [Q_u I_1 + Q_d I_2] \quad (7)$$
\[
F_2 = \frac{-i e g N_c m_t}{4 \sqrt{2} \pi^2 \delta} [Q_u I_3 + Q_d I_4]
\]

where \(N_c = 3\) is the usual color factor, \(m_t\) is the t-quark mass, \(g\) is the conventional weak coupling constant, \(Q_{u,d}\) are the electric charges of the up- and down-quarks, and \(I_i\) can be expressed as sums of parameter integrals:

\[
I_1 = 1 + (2m_t^2 \delta^{-1} - 1)G_{-1}(m_t, m_b) + 2[\delta^{-1}(m_b^2 - m_t^2 - M_W^2) + 1/2]G_0(m_t, m_b)
+ 2M_W^2 \delta^{-1} G_1(m_t, m_b)
\]
\[
I_2 = 1 + 2m_b^2 \delta^{-1}G_{-1}(m_b, m_t) + 2[\delta^{-1}(m_t^2 - m_b^2 - M_W^2) - 1/2]G_0(m_b, m_t)
+ 2M_W^2 \delta^{-1} G_1(m_b, m_t)
\]
\[
I_3 = G_{-1}(m_t, m_b) - G_0(m_t, m_b)
\]
\[
I_4 = -G_0(m_b, m_t)
\]

where the \(G_n\) are given by

\[
G_n(m_i, m_j) = \int_0^1 dzz^n \ln \left[ \frac{m_i^2 (1 - z) + m_j^2 z - z(1 - z)m_i^2}{m_i^2 (1 - z) + m_j^2 z - z(1 - z)M_W^2} \right]
\]

It is important to note that both \(F_{1,2}\) are proportional to \(m_t\).

Similar considerations apply to the Z decay to sneutrino plus photon. As in the previous case, here also we have an enhancement arising from the presence of the top-quark mass although a suppression also occurs due the heavy top-quark propagators in the loop. The decay rate for the radiative Z decay can be written as:

\[
\Gamma(Z \rightarrow \bar{\nu} \gamma) = \frac{\alpha^2 M_Z^3}{192\pi^3 m_t^2} f^2 \left(1 - \frac{m_t^2}{M_Z^2}\right)^3 \frac{(1 - 8/3\sin^2 \theta_W)^2 I}{\sin^2 \theta_W \cos^2 \theta_W} \tag{10}
\]

with

\[
I = |I_1(x, y) - I_2(x, y)|^2 + |I_2(x, y)|^2 \tag{11}
\]
where the $I'_i$'s are complicated functions of $x = (4m^2_i/m^2_{\tilde{\nu}})$ and $y = (4m^2_i/M^2_Z)$ given in Ref.11.

The predictions for the $W \rightarrow \tilde{l}\gamma$ and $Z \rightarrow \tilde{\nu}\gamma$ branching fractions as functions of the $\tilde{\nu}$ and $\tilde{l}$ masses are given in Figs.3 and 4 assuming $f=1$. We see that they can be as large as $10^{-5}$.

Let us now discuss the experimental signatures for these processes. Due to R-parity breaking, the slepton or the sneutrino which appears as the decay product of the $W$ or $Z$ will itself decay predominantly into quarks. In the case of the slepton (i.e. $\tilde{\tau}^+$), an analysis of the same diagram as in Fig.1 shows that it decays predominantly to the charm and strange quarks if the sum of the masses for the chargino and neutralino is larger than that of the slepton as is usually expected in most models. Turning now to the sneutrino final state, it will predominantly decay into $c\bar{c}$ and $b\bar{b}$ modes. Thus the overall signature for R-parity violating radiative $W$ and $Z$ decays in this model will be two jets plus a photon with the mass of the two-jets reconstructing to that of the corresponding slepton. These two jets will having leading heavy flavor components. As far as feasibility of detection of the above processes is concerned, the radiative $W$ decay suggested in this paper is probably beyond the reach of LEPII but might be observable at a high luminosity $e^+e^-$ collider with a larger center of mass energy. At hadron supercolliders such as the SSC or LHC, the rather large luminosity available implies[12] that we should expect more than $10^3$ events per year. The main difficulty at such machines will be the backgrounds from other sources. The radiative $Z$ decay should be observable at LEPI by sitting on the $Z$ peak provided sufficient integrated luminosity is accumulated.

It is worth pointing out that in a non-supersymmetric two Higgs model, if the charged Higgs boson is lighter than the $W$ boson, a similar appearing decay can arise. It is however likely to be smaller due to the presence of the vacuum mixing parameter $\tan\beta$.

**IV.Neutralino Decay Modes**
In this section, we briefly mention some other distinct experimental signatures of the model that arise in the neutralino sector. This sector can be split into two parts in our scenario: one heavy and the other light, the former consisting of the $\nu_\tau$ and the heavy Zino and the latter consisting of the Higgsinos and the lighter Zinos of the MSSM. The heavier neutralinos can decay into $\tau^+\tau^-\nu_\tau, \tau^-c\bar{s}$ etc., due to spontaneous R-parity breaking induced interactions. In the light neutralino sector, if we ignore the small vev of $\tilde{\nu}$, then there are only four mass eigenstates. While our comments below apply to all the eigenstates which have a significant photino component contained within them, we will give the generic name photino to only one of the neutralinos in what follows. A specific prediction of our model, in contrast to other R-parity breaking models[13], is that the photino can decay only to $\tau^-c\bar{s}$ as opposed to decays like $\tau^-\mu^+\nu_e$ etc. Such decay modes have recently been searched for by the OPAL collaboration[14] with a null result. The present model on the other hand would lead only to signatures of type $e^+e^- \to \tilde{\gamma}\tilde{\gamma}$ with the photinos subsequently decaying to $\tau^-c\bar{s}$. This is a very different signature than previously considered.

V. Conclusion

We have isolated a class of models with spontaneous R-parity violation where the single photon radiative decays of $W$ and $Z$ are significant in contrast with the MSSM with explicit R-parity violation wherein such decays are highly suppressed. Any evidence for such decays could indicate the existence of new gauge symmetries beyond the standard model. The attractive aspect of the model we consider here is that in the limit of an exact extended gauge symmetry, R-parity remains unbroken[7,15] so that any manifestation of R-parity breaking at low energy appears only in the form of lepton number violation and baryon number remains an exactly conserved symmetry. Such models can appear naturally in the low energy limit of some superstring theories. This should provide a strong motivation to look for such decays of the $W$ and $Z$ at LEP and at other colliders in the future. We also
point out some new R-parity breaking decays of the neutralinos that can be sought at both LEPI and LEPII.

Acknowledgement

One of us (R.N.M.) would like to thank K.S. Babu and A. Jawahery for discussions. T.G.R. would like to thank J.L. Hewett for discussions. The work of R.N.M. was supported by a grant from the National Science Foundation and the work of T.G.R. was supported by a grant from the Department of Energy.
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Figure Caption

Figure 1. The Feynman diagram that induces the R-violating slepton-t-b coupling after B-L breakdown.

Figure 2. The one-loop diagram that induces the process $W \to \tilde{\tau}_\gamma$ decay amplitude.

Figure 3. $W \to \tilde{\tau}_\gamma$ branching fraction predicted in our model as a function of the slepton mass assuming $f=1$. The solid (dash-dotted, dashed) curve corresponds to a top-quark mass of 100 (150, 200)GeV.

Figure 4. $Z \to \tilde{\nu}_\gamma$ branching fraction as a function of the sneutrino mass assuming $f=1$. The curves are for the same top-quark masses as in Fig.3.