Lepton Number Violating Radiative $W$ Decay in Models with R-parity Violation

THOMAS G. RIZZO
High Energy Physics Division
Argonne National Laboratory
Argonne, IL 60439

Abstract

Models with explicit R-parity violation can induce new rare radiative decay modes of the $W$ boson into single supersymmetric particles which also violate lepton number. We examine the rate and signature for one such decay, $W \to \tilde{l}\gamma$, and find that such a mode will be very difficult to observe, due its small branching fraction, even if the lepton number violating coupling in the superpotential is comparable in strength to electromagnetism. This parallels a similar result obtained earlier by Hewett in the case of radiative $Z$ decays.

*Research supported by the U.S. Department of Energy, Division of High Energy Physics, Contract W-31-109-ENG-38.
In the simple Minimal version of the Supersymmetric Standard Model (MSSM), both baryon ($B$) and lepton ($L$) numbers are simultaneously conserved quantities due to the imposition of a discrete symmetry called R-parity. One can then assign to each Standard Model (SM) particle and its SUSY partner a multiplicatively conserved quantum number, given by $R = (-1)^{2S + 3B + L}$, where $S$ is just the particle’s spin. In addition to the gauge symmetries and the assumption of minimal particle content, R-parity conservation severely limits the possible interactions among the usual SM fermions and their superpartners as well as the properties of the SUSY partners themselves. The two most important phenomenological consequences \[1\] are well-known: (i) The SUSY partners of the conventional SM particles which carry negative R-parity can only be produced in pairs and (ii) the lightest supersymmetric particle (LSP), is an electrically neutral, color-singlet and is stable. This second property is the one which is conventionally used experimentally to search for SUSY particles, i.e., once they are pair-produced their decays involve final states which include the LSP that only appears as a missing energy signature in a collider detector. If R-parity were broken both these conclusions would be invalidated leading to an entirely new phenomenology.

Of course this minimal approach may not be that realized by nature. In particular, it is possible to construct phenomenologically viable models wherein R-parity is violated either spontaneously\[2\], through the acquisition of a vacuum expectation value (vev) by a sneutrino, or explicitly via the existence of additional terms in the superpotential\[3\] constructed out of the conventional superfields. These additional terms are possible as the gauge symmetries alone do not forbid their existence. If such new interactions are present they can lead not only to a destabilization of the LSP but also new production modes for SUSY partners not present in the MSSM and thus forces us to a re-evaluate of the traditional search techniques for these particles. Clearly, the breaking of R-parity implies that these new interactions will violate $L$ and $B$. In such a R-parity breaking scenario this more general form of the
superpotential, $W$, can be written as

$$ W = h_{ij} L^i_L H_2 \bar{E}^j_R + h'_{ij} Q^i_L H_2 \bar{D}^j_R + h''_{ij} Q^i_L H_1 \bar{U}^j_R $$

$$ + \lambda_{ijk} L^i_L L^j_L \bar{E}^k_R + \lambda'_{ijk} L^i_L Q^j_L \bar{D}^k_R $$

$$ + \lambda''_{ijk} \bar{U}^i_R \bar{D}^j_R \bar{D}^k_R, $$

where $ijk$ are generation indices, and the $h$'s and $\lambda$'s are a priori unknown Yukawa couplings. $Q_L, L_L, etc$ represent the usual left-handed chiral superfields of the MSSM. Whereas the $h$ couplings are the standard ones responsible for generating fermion masses when the scalar components of $H_1$ and $H_2$ acquire vev’s, the $\lambda, \lambda'$, and $\lambda''$ terms lead to generation-dependent $\Delta L$ and $\Delta B$ interactions. Of course, if all such terms are allowed simultaneously, proton decay proceeds unsuppressed and thus, both $\Delta L$ and $\Delta B$ terms cannot be present. For example, if $\lambda = \lambda' = 0$, only $B$ violation would occur. In particular, it has be recently shown \cite{4} that if one assumes only the particle content of the MSSM and the lack of rapid proton decay, then two unique discrete symmetries are possible: $R$ and $B$, the so-called baryon-parity. Thus the most reasonable scenario to consider for our purposes is one in which $B$ is conserved while both $R$ and $L$ are violated; this corresponds to setting all of the $\lambda''$ terms to zero in $W$. In principle, this class of models should be considered as to be just as likely a scenario for the realization of SUSY as is the more conventional MSSM with the same particle content. This is the scheme we consider below.

The interaction Lagrangians that result from $W$ can be written as

$$ \mathcal{L} = \lambda_{ijk} [\tilde{\nu}_L^i \tilde{e}_R^k e_L^j + \tilde{e}_L^j \tilde{e}_R^k \nu_L^i + (\tilde{e}_R^k)^* (\tilde{\nu}_L^i)^c e_L^j - (i \leftrightarrow j)] + h.c. $$

and

$$ \mathcal{L} = \lambda'_{ijk} [\tilde{\nu}_L^i \tilde{d}_R^k d_L^j + \tilde{d}_L^j \tilde{d}_R^k \nu_L^i + (\tilde{d}_R^k)^* (\tilde{\nu}_L^i)^c d_L^j $$

$$ - \tilde{\nu}_L^i \tilde{d}_R^k u_L^j - \tilde{u}_L^j \tilde{d}_R^k e_L^i - (\tilde{d}_R^k)^* (\tilde{\nu}_L^i)^c u_L^j] + h.c. $$

3
Once B-parity has been imposed, many of the remaining couplings can be constrained via various phenomenological considerations using these Lagrangians as has been done by several groups of authors[5]. Combining all of their results, one finds that the only potentially large couplings, \( i.e. \), ones that are capable of being of the same strength as electromagnetism, are[6]: \( \lambda_{131}, \lambda'_{3jk}, \lambda'_{121}, \lambda'_{222}, \lambda'_{223}, \lambda'_{232}, \) and \( \lambda'_{233}. \) Note particularly that interactions involving third generation particles are least likely to be constrained by existing data.

Although one can look for the direct influence of R-parity violating interactions, it may be possible to search for their indirect effects through, \( e.g. \), loop diagrams. One such possibility, recently examined by Hewett [7], is the loop-induced decay \( Z \to \tilde{\nu}\gamma, \) with a rate that might be only an order of magnitude smaller than that for \( Z \to H\gamma \) in the SM. In this paper we wish to examine the corresponding process in \( W^\pm \) decay, \( i.e. \), \( W \to \tilde{l}\gamma. \) Unfortunately, as we will see, the branching fraction for this process is found to be much smaller than in the \( Z \) case due to helicity supression and the form of the superpotential, \( W. \)

The diagrams responsible for the \( W \to \tilde{l}\gamma \) decay are shown in Fig. 1 and involve one R-parity violating vertex. For simplicity, we have assumed that only one of the two \( \Delta L \) couplings is non-zero, \( i.e. \), we ignore the possible contributions of the \( \lambda \) terms in the superpotential and concentrate on the \( \lambda' \) terms since more of them can be large. We note, however, that since \( \lambda'_{131} \) can be sizeable, a loop involving first generation leptons might yield a significant contribution. In fact, one finds that such contributions are suppressed due to the small masses of these particles. This being the case, it is the couplings \( \lambda'_{233} \) and \( \lambda'_{333} \) which are relevant here. In the former case, the \( \tilde{l} \) is a smuon whereas in the latter it is a stau. We, of course, do not know the size of this active \( \lambda' \) coupling so we simply scale it to the electromagnetic strength as is customary, \( i.e. \), we write the effective \( t\tilde{l} \) interaction as

\[
\mathcal{L} = -e\sqrt{F} \bar{l}_L \bar{b}(1 - \gamma_5)t
\]
where $F$ is an unknown parameter which may be as large as unity. With this normalization, the amplitude for the $W \rightarrow \tilde{l}\gamma$ decay process can be written as

$$\mathcal{A} = \left[ F^1 \left( \frac{q_\mu k_\nu - g_\mu \nu k \cdot q}{M_W^2} \right) + iF^2 \frac{\epsilon_\mu \nu \sigma \tau q^\sigma k^\tau}{M_W^2} \right] \epsilon^\nu_{\gamma W}$$

(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

$$\Gamma(W \rightarrow \tilde{l}\gamma) = \frac{M_W^3}{96\pi^2} (F^1 + F^2)^2 \left( 1 - \frac{m_{\tilde{l}}^2}{M_W^2} \right)^3$$

(6)

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

$$F^1 = -\frac{i\alpha g N_c m_b \sqrt{F}}{\sqrt{2\pi} \delta} [Q_u I_1 + Q_d I_2]$$

(7)

$$F^2 = -\frac{i\alpha g N_c m_b \sqrt{F}}{\sqrt{2\pi} \delta} [Q_u I_3 + Q_d I_4]$$

where $N_c = 3$ is the usual color factor, $m_b$ is the b-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:

$$I_1 = 1 + 2m_{\tilde{l}}^2 \delta^{-1} G_{-1}(m_t, m_b) + 2[\delta^{-1}(m_b^2 - m_{\tilde{l}}^2 - M_W^2) - 1/2]G_0(m_t, m_b)$$

$$+ 2M_W^2 \delta^{-1} G_1(m_t, m_b)$$

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

$$+ 2M_W^2 \delta^{-1} G_1(m_b, m_t)$$

(8)

$$I_3 = -G_0(m_t, m_b)$$

$$I_4 = G_{-1}(m_b, m_t) - G_0(m_b, m_t)$$
where the $G_n$ are given by

$$G_n(m_i, m_j) = \int_0^1 dz z^n \ln \left[ \frac{m_i^2 (1 - z) + m_j^2 z - z(1 - z) m_{\tilde{l}}^2}{m_i^2 (1 - z) + m_j^2 z - z(1 - z) M_W^2} \right]$$  \hspace{1cm} (9)$$

It is important to note that both $F_{1,2}$ are proportional to $m_b$ and not $m_t$. This comes about due to the fact that the $W$ charged current interactions are purely left-handed, as is the coupling in Eq.(4), and that unless a mass term from a propagator of an internal quark line is picked up, the result will vanish since a trace over an odd number of $\gamma$-matrices will then be taken. In the actual calculation only the b-quark mass term is picked up so that instead of increasing in magnitude as $m_t$ increases, the rate for $W \to \tilde{l}\gamma$ will decrease for large $m_t$ due to suppression from the top-quark propagator.

Using $m_b=5$ GeV, $\alpha^{-1}=127.9$, and $M_W=80.15$ GeV as numerical input, we can calculate the partial decay width as a function of $m_\tilde{l}$ for different values of $m_t$ and the parameter $F$. The result of this calculation, expressed as the branching fraction(B) for the $W \to \tilde{l}\gamma$ process, is shown in Fig.2. As advertised, even for $F=1$, we see that B is less than about $10^{-8}$ for all values of $m_t$ and $m_\tilde{l}$ and, as anticipated, falls with increasing $m_t$. This rate is sufficiently tiny that this reaction will be impossible to observe. If the helicity structure of the R-parity violating interaction had been opposite, the rate could have been larger by a factor of $m_t^2/m_b^2 \simeq 10^3$ and potentially observable since the overall numerical factor in the amplitude would then have been proportional to $m_t$ instead of $m_b$.

Unfortunately, our result implies that signals for R-parity violating interactions must be sought elsewhere than in radiative $W$ decays.
ACKNOWLEDGEMENTS

The author would like to thank J.L. Hewett for discussions related to this work. This research was supported in part by the U.S. Department of Energy under contract W-31-109-ENG-38.
References

[1] For an introduction and phenomenological overview of SUSY, see H.E. Haber and G.L. Kane, Phys. Rep. 117, 75 (1985).

[2] M.C. Gonzalez-Garcia, J.C. Romano, and J.W.F. Valle, University of Valencia report FTUV-91-42 (1991); J.C. Romao, C.A. Santos, and J.W.F. Valle, Phys. Lett. B288, 311 (1992); J.C. Romao and J.W.F. Valle, Nucl. Phys. B381, 87 (1992); A. Masiero and J.W.F. Valle, Phys. Lett. B251, 273 (1990); A. Santamaria and J.W.F. Valle, Phys. Rev. D39, 1780 (1989) and Phys. Lett. B195, 423 (1987); J.C. Romao, F. de Campos, and J.W.F. Valle, Phys. Lett. B292, 329 (1992).

[3] L.J. Hall and M. Suzuki, Nucl. Phys. B231, 419 (1984); S. Dawson, Nucl. Phys. B261, 297 (1985); S. Dimopolous and L.S. Hall, Phys. Lett. B207, 210 (1987); H. Dreiner and G. Ross, Nucl. Phys. B365, 597 (1991) and Oxford University report OUTP-92-08P (1992); R. Mohapatra, Phys. Rev. D34, 3457 (1989); E. Ma and D. Ng, Phys. Rev. D41, 1005 (1990); M. Doncheski and J.L. Hewett, Argonne National Laboratory report ANL-HEP-PR-92-28 (1992), to appear in Z. Phys. C.

[4] M.C. Bento, L.J. Hall, and G.G. Ross, Nucl. Phys. B292, 400 (1987); see also the paper by Dreiner and Ross in Ref. 2.

[5] See, for example, V. Barger, G.F. Giudice, and T. Han, Phys. Rev. D40, 2987 (1989); C.S. Aulakh and R.N. Mohapatra, Phys. Lett. B119, 316 (1982); L.J. Hall, Mod. Phys. Lett. A5, 467 (1990); S. Dimopoulos et al., Phys. Rev. D41, 2099 (1990); R. Arnowitt and P. Nath in Phenomenology of the Standard Model and Beyond, eds by D.P. Roy and P. Roy, World Scientific, Singapore, 1989.

[6] R.M. Godbole, P. Roy, and X. Tata, CERN report CERN-TH.6613/92 (1992).
[7] J.L. Hewett, Argonne National Laboratory report ANL-HEP-CP-92-23 (1992), and talk given at the Workshop on Photon Radiation From Quarks, Annecy, France, December 2-3, 1991.
Figure Captions

Figure 1. The Feynman diagrams responsible for $W \rightarrow \tilde{\ell}\gamma$ in R-parity violating models.

Figure 2. The branching fraction, $B$, for the decay $W \rightarrow \tilde{\ell}\gamma$ as a function of the slepton mass, $m_{\tilde{\ell}}$, assuming $F=1$ for various values of the top quark mass: $m_t=100$ (150, 200) GeV corresponds to the solid (dash-dotted, dashed) curve.