New contributions to neutrinoless double-beta decay in supersymmetric theories

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In supersymmetric theories with R-parity violation, there are new contributions to neutrinoless double-beta $[(\beta\beta)_{0\nu}]$ decay that do not involve the exchange of Majorana neutrinos. Experimental information on $(\beta\beta)_{0\nu}$ decay can therefore be used to constrain parameters of supersymmetric theories. We also discuss neutrino mass in these theories. We then discuss how R-parity-violating interactions can be induced in the low-energy sector of a theory that respects R-parity conservation prior to symmetry breaking.

I. INTRODUCTION

Neutrinoless double-beta decay[1] $[(\beta\beta)_{0\nu}$ decay] is believed to be a sensitive probe of physics beyond the standard model. In the standard $SU(2)_L \times U(1)$ model since $B-L$ conservation is automatic, $(\beta\beta)_{0\nu}$ vanishes to all orders in the weak interactions. However, once the possibility of $B-L$ violation (either spontaneous or explicit) is accepted, the above process, which violates the $B-L$ quantum number by two units, can occur in a variety of ways. The simplest and the most well-known way[2] is via the exchange of Majorana neutrinos or the presence of $(B-L)$-violating right-handed currents, and in fact, present limits on the lifetime for $(\beta\beta)_{0\nu}$ decay are generally interpreted to give an upper limit on the Majorana mass of the neutrino. It has since been pointed out that, in left-right-symmetric theories with Majorana neutrinos,[3] there exist new contributions to $(\beta\beta)_{0\nu}$ decay that involve the exchange of heavy right-handed neutrinos[3,4] as well as the exchange of doubly charged Higgs bosons.[5] Experimental information on $(\beta\beta)_{0\nu}$ decay can, therefore, be used to constrain parameters of the left-right-symmetric models. In this paper, we point out that general supersymmetric extensions of the standard model give rise to lepton-number-violating interactions,[6] which, in turn, can lead to completely new contributions to $(\beta\beta)_{0\nu}$ decay that involve exchange of gluinos and photinos (and not Majorana neutrinos). Existing low-energy data constrain the parameters that control the magnitude of the double-beta decay rate but are such that the new contribution can be large enough to be visible in on-going experiments. Turning the question around, experimental data on neutrinoless double-beta decay can be used to constrain the parameters of the supersymmetric model. We also study the neutrino masses in these models and then discuss the simplest extensions of the standard model where $B-L$ conservation is automatically restored. In such models, R parity can be broken by the vacuum leading to an effective theory at low energies that has induced R-parity-violating interactions. This kind of embedding may help to fix the strength of the R-parity-violating interactions.

II. DESCRIPTION OF THE MODEL

The minimal supersymmetric extension of the standard model[7] consists of the following superfields, with their $SU(2)_L \times U(1) \times SU(3)_c$ transformation indicated within parentheses next to the fields: quarks Q $(2,\frac{1}{3},3)$; U^c $(1,-\frac{4}{3},3^*)$; D^c $(1+\frac{2}{3},3^*)$; leptons L $(2,-1,1)$; E^c $(1,+2,1)$; Higgs bosons H_u $(2,1,0)$; H_d $(2,-1,0)$. We have suppressed the generation index for fermions and assume that supersymmetry (SUSY) breaking is dictated by $N=1$ supergravity.[7] The nongauge interactions of this model are specified by the following general form of the superpotential where we have imposed global baryon-number symmetry to avoid rapid proton decay:

$$W=h_u Q H_u U^c + h_d Q H_d D^c + h_e L H_d E^c + \mu_1 H_u H_d + \mu_2 H_u L + f_{pqr} Q_p L_q D_r^c + \widetilde{h}_{[p,q]r} L_p L_q E_r^c \;, \qquad (1)$$

where p,q,r stand for generation indices and $[p,q]$ implies antisymmetry in the two indices; h_u, h_d, and h_e are matrices in generation space. Supersymmetry-breaking terms in the Lagrangian dictated by a super-Higgs effect in an $N=1$ supergravity model have the form

$$\mathscr{L}_{\rm SB}=m_{3/2}\int d^2\theta\,\theta^2 W+\int d^2\theta\,\theta^2\mu_a W_a^\alpha W_{\alpha,a}+\text{H.c.}\;, \qquad (2)$$

where a goes over all the gaugino fields of the model such as the W gaugino ($\widetilde{W}$) Z gaugino ($\widetilde{Z}$), photino ($\widetilde{\gamma}$), and gluino ($\widetilde{g}$); W^α represents the supersymmetric representation of the gauge field tensor $F_{\mu\nu}$; and θ is the two-component fermionic coordinate.

It is clear from Eq. (1) that R parity[8] defined by $(-1)^{3B+L+2S}$ (where B, L, and S stand for baryon number, lepton number, and spin, respectively) and lepton-number symmetry are violated by the μ_2, f, and $\widetilde{h}$ terms in the superpotential. They will, therefore, lead to processes that violate lepton-number conservation. Many of these processes have already been studied previously.[9] Here, we wish to focus on the neutrino mass and neutrinoless double-beta decay. Before proceeding to this applica-

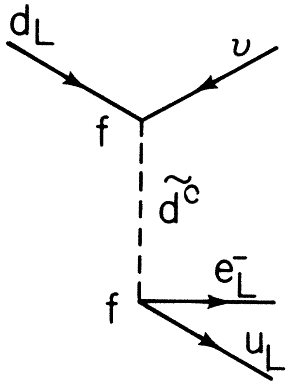

FIG. 1. A new contribution to neutron beta decay in a supersymmetric extension of the standard model.

tion, we would like to briefly comment on the allowed ranges of the couplings f and $\tilde{h}$. The most stringent constraints on the couplings come from $\mu \to 3e$ and $K_L \to \mu \bar{e}$ decay;[1] but both these constraints can be satisfied by constraining $\tilde{h}_{[\mu,e]e^c}$ and $\tilde{h}_{[\mu,e]\mu^c}$ to be tiny; but since these couplings do not play any role in neutrinoless double-beta decay, we will not concern ourselves with the details and simply set $\tilde{h}=0$. Also for simplicity, let us set $\mu_2=0$ and assume that f is diagonal in the quark indices. The part of the f coupling relevant in our discussion is given by

$$W_{\Delta L \neq 0} = f_{ude} Q_1 L_1 D_1^c + \text{H.c.} + \cdots \ , \tag{3}$$

where the subscript 1 stands for fermions of the first generation.

Let us now turn to the magnitude of the coupling. To study this, we note that Eq. (3) contributes to β decay of neutrino via the Feynman diagram in Fig. 1. The strength of this contribution is given by $f_{ude}^2/M_{\tilde{d}^c}^2$ and has pure $V-A$ form, as can be easily seen after Fierz reshuffling. As a result, it adds to the usual Fermi coupling. This contribution must therefore be less than 1% of the Fermi coupling from observations on β decay and $\pi \to e\bar{\nu}$ decay leading to the constraint

$$f_{ude}^2/M_{\tilde{d}^c}^2 \le 2\sqrt{2} G_F \times 10^{-2} \ . \tag{4}$$

III. MAGNITUDE FOR $(\beta\beta)_{0\nu}$ DECAY

Coupled with the Majorana masses for the Z gaugino, photino, and gluino, the R-parity-violating interaction in Eq. (3) will contribute to $(\beta\beta)_{0\nu}$ decay via the Feynman diagrams in Figs. 2(a)—2(c). The contribution of the gluino exchange diagram in Figs. 2(a) and 2(c) can be estimated to be

$$M^{(\tilde{g})} \simeq 2 f_{ude}^2 \left[\frac{1}{M_{\tilde{u}}^4} + \frac{1}{M_{\tilde{d}}^4 c} \right] \times \frac{M_{\tilde{g}} 4\pi\alpha_s}{p^2 + M_{\tilde{g}}^2} \bar{u}^{\alpha}_L d_{R\beta} \bar{u}^{\beta}_L d_{R,\alpha} e_L^{\dagger} C e_L^{*} \tag{5}$$

(α, β denote the color index), whereas those involving the exchange of $\tilde{\gamma}$, $\tilde{Z}$ from Figs. 2(a) and 2(b) are

$$M^{(\tilde{Z})} \simeq \frac{f_{ude}^2 g^2}{\cos^2\theta_W} \left[\frac{\frac{1}{2} - \frac{2}{3}\sin^2\theta_W}{M_{\tilde{u}}^4} + \frac{-\frac{1}{2} + \sin^2\theta_W}{M_{\tilde{e}}^4} + \frac{\frac{1}{3}\sin^2\theta_W}{M_{\tilde{d}^c}^4} \right] \bar{u}^{\alpha}_L d_{R\alpha} \bar{u}^{\beta}_L d_{R\beta} e_L^{\dagger} C e_L^{*} \frac{M_{\tilde{Z}}}{p^2 + M_{\tilde{Z}}^2} \ , \tag{6}$$

$$M^{(\tilde{\gamma})} \simeq f_{ude}^2 4\pi\alpha \left[\frac{2}{3M_{\tilde{u}}^4} - \frac{1}{M_{\tilde{e}}^4} + \frac{1}{3M_{\tilde{d}^c}^4} \right] \bar{u}^{\alpha}_L d_{R\alpha} \bar{u}^{\beta}_L d_{R\beta} e_L^{\dagger} C e_L^{*} \frac{M_{\tilde{\gamma}}}{p^2 + M_{\tilde{\gamma}}^2} \ . \tag{7}$$

In order to estimate these contributions we first consider the gluino-exchange diagram which has the strongest couplings, among all these graphs. Although it involves color-octet operators if the gluino is heavy, the net interaction is pointlike; as a result, we can do a Fierz reshuffling and make both the operators into color-singlet ones with an extra factor $\frac{4}{3}$. For further estimation, let us assume $M_{\tilde{u}} = M_{\tilde{d}^c}$; since the present collider result seems to indicate[10] that $M_{\tilde{u}} = M_{\tilde{d}^c} \equiv M_{\tilde{q}} \ge 60$ GeV, with $M_{\tilde{g}} > 100$ GeV; using this, we can predict the strength for neutrinoless double-beta transition decay to be

$$A^{(\tilde{g})}_{(\beta\beta)_{0\nu}} \simeq \frac{16}{3} \frac{f_{ude}^2}{M_{\tilde{u}}^4} 4\pi\alpha_s M_{\tilde{g}} A_{\text{nuc}}(M_{\tilde{g}}) \ , \tag{8}$$

where $A_{\text{nuc}}(M_{\tilde{g}})$ represents the nuclear matrix element of the hadronic part of M.

However, we wish to emphasize that, crucial to our conclusion is the assumption that we can do Fierz reshuffling leading to color-singlet operators which is, of course, valid in the strict point-interaction limit. If this is not allowed, then, the matrix element will be extremely suppressed since we will need colored intermediate states. These objections do not apply to $\tilde{\gamma}$ or $\tilde{Z}$ exchange, which we consider now. Again, for simplicity, we assume $M_{\tilde{u}} = M_{\tilde{e}} = M_{\tilde{d}^c} \gtrsim 60$ GeV for the $\tilde{Z}$ exchange and we get for the strength of the $(\beta\beta)_{0\nu}$ amplitude

$$A^{(\tilde{Z})}_{(\beta\beta)_{0\nu}} \simeq \frac{f_{ude}^2 7 \times 10^{-2}}{M_{\tilde{q}}^4} M_{\tilde{Z}} A_{\text{nuc}}(M_{\tilde{Z}}) \ . \tag{9}$$

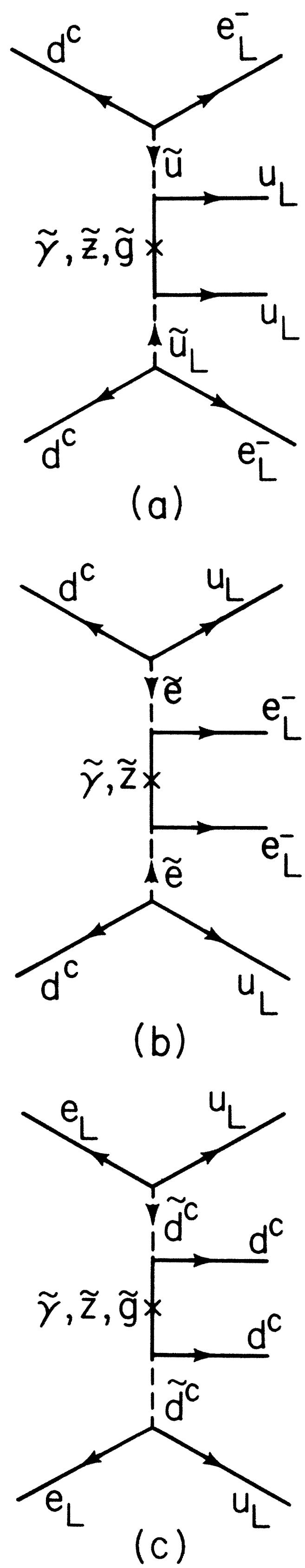

FIG. 2. New contributions to neutrinoless double-beta decay in supersymmetric extension of the standard model mediated by (a) photino ($\tilde{\gamma}$), (b) Z-gaugino ($\tilde{Z}$), and (c) gluino ($\tilde{g}$) exchange.

As far as the photino contribution goes if we make the same assumption as above, the amplitude vanishes; however, experimentally, the lower bound on $M_{\tilde{e}}$ is about 50 GeV, whereas as mentioned before collider results[10] put a lower bound on $M_{\tilde{q}} \geq 60$ to 75 GeV. As a result, we expect the cancellation to be imperfect and if we define $\Delta m^2 = M_{\tilde{q}}^2 - M_{\tilde{e}}^2$, we find

$$A^{(\tilde{\gamma})}_{(\beta\beta)_{0\nu}} \simeq \frac{f_{ude}^{\,2}}{M_{\tilde{q}}^{4}} \left[\frac{\Delta m^2}{M_{\tilde{e}}^2} \right] 9\times 10^{-2} M_{\tilde{\gamma}} A_{\rm nuc}(M_{\tilde{\gamma}}) \left[1 + \frac{m_{\tilde{q}}^2}{m_{\tilde{e}}^2} \right] . \tag{10}$$

The photino and Z-gaugino exchange diagrams could be comparable, because while the present lower limit on $M_{\tilde{\gamma}}$ is about[10] 5 GeV and that on $M_{\tilde{Z}}$ could be of order[11] 33 GeV, we assume the factor $(1 + m_{\tilde{q}}^2/m_{\tilde{e}}^2)\Delta m^2/M_{\tilde{e}}^2 = \frac{1}{5}$ to $\frac{1}{10}$. From now on we will assume the effective $\tilde{Z}$ strength to be comparable.

Let us now turn to discussion of the nuclear matrix element in order to put bounds on scalar-quark and scalar-lepton masses from experimental limits on double-beta decay. For this purpose, we will consider the gluino diagram separately from the photino and Z-gaugino diagram. In the case of the former, we assume that we can make a Fierz transform.

The nuclear matrix elements here are Fermi type and, as has been noted,[12] they are expected to be about 1% of the corresponding Gamow-Teller (GT) matrix elements, due to the fact that there are no nuclear states in the same isomultiplet as the ones that undergo double-beta transition. Using a finite-size approximation for the nucleus[13] and following the discussion of Haxton and Stephenson[1] (their Fig. 14b) and using the suppression of the Fermi-type matrix element, $M_F = 10^{-2} M_{\rm GT}$, we estimate that the constraint of $(\beta\beta)_{0\nu}$ decay of ^{76}Ge implies (for a typical gluino, photino, or Z-gaugino mass of 100 GeV)

$$\frac{f_{ude}}{M_{\tilde{u}}^{4}} a_x \lesssim G_F^2 \times 10^{-2.5} , \tag{11}$$

TABLE I. We summarize the lower bounds on the scalar-quark masses for various values of the gaugino mass coming from a lower bound on ^{76}Ge 0ν half-life of 2.5×10^{23} yr. The gluino case and the $\tilde{Z}+\tilde{\gamma}$ case are given separately.

Gaugino mass	Lower bound on $m_{\tilde{u}}$ for the gluino case	Lower bound on $m_{\tilde{u}}$ for the $\tilde{Z}+\tilde{\gamma}$ case
100 GeV	3 TeV	440 GeV
10^3 GeV	950 GeV	150 GeV
10^4 GeV	300 GeV	40 GeV
10^5 GeV	105 GeV	

where

$$a_{\tilde{g}} \approx \tfrac{16}{3}$$

or

$$a_{\tilde{Z}+\tilde{\gamma}} = 1.6 \times 10^{-1} , \tag{12}$$

where we have assumed

$$M_{\tilde{Z}} A_{\rm nuc}(M_{\tilde{Z}}) \approx (\Delta m^2 / M_{\tilde{e}}^2) M_{\tilde{\gamma}} A_{\rm nuc}(M_{\tilde{\gamma}}) .$$

Clearly, using the maximum value in the constraint in Eq. (4), we find, for the gluino case

$$M_{\tilde{q}} \geq 2 \text{ TeV} \tag{13}$$

and for the photino, Z-gaugino case

$$M_{\tilde{q}} \geq 380 \text{ GeV} .$$

In Table I, we give the constraints for various values of the gaugino masses using Fig. 14 of Haxton *et al.*[1] as rescaled by Caldwell *et al.* using the ^{76}Ge 0ν half-life[14] of $\geq 2.5 \times 10^{23}$ yr.

IV. NEUTRINO MASS

The dominant contribution to neutrino masses in this class of theories will come from Feynman diagrams depicted in Fig. 3. They involve two electroweak symmetry-breaking doublet vacuum expectation values (VEV's) and also a supersymmetry-breaking parameter $m_{3/2}$. If we ignore the contribution of higher generations to the electron neutrino mass, then we estimate that the induced Majorana mass for the neutrino from Fig. 3 is given by

$$m_\nu \simeq \frac{f_{ude}^{\,2}}{16\pi^2} \frac{m_d^{\,2}}{m_{3/2}} . \tag{14}$$

For a scalar-quark mass of 500 GeV, $f_{ude}^{\,2} \leq 7.5 \times 10^{-2}$; choosing the SUSY-breaking scale $m_{3/2} \simeq 100-1000$ GeV, we find $m_\nu \leq 0.5-0.05$ eV. This is interesting since it implies that contribution of the neutrino mass to $(\beta\beta)_{0\nu}$ decay is about 10–100 times smaller than the contributions of direct R-parity-violating interactions. As a result, detailed characteristics of $(\beta\beta)_{0\nu}$ decay must be studied before one can conclude a value for the Majorana mass for the neutrino from the $(\beta\beta)_{0\nu}$ lifetime, once it is observed.

V. R-PARITY CONSERVATION IN EXTENDED TO ELECTROWEAK MODELS

We now wish to point out that the existence of an R-parity- (and L-) violating interaction depends on the kind of electroweak gauge group one works with. As an illustrative example, consider the SUSY electroweak group to be $SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$. In this case, there exist no interactions[15] in the superpotential that violate lepton number. To see this, we assign the quarks and leptons to the gauge group as follows:[15]

$$\text{for quarks,}\quad Q(2,0,\tfrac{1}{3}),\ U^c(1,-\tfrac{1}{2},-\tfrac{1}{3}),\ D^c(1,+\tfrac{1}{2},-\tfrac{1}{3}) ;$$

$$\text{for leptons,}\quad L(2,0,-1),\ E^c(1,+\tfrac{1}{2},+1),\ \nu^c(1,-\tfrac{1}{2},+1) ;$$

$$\text{for Higgs bosons,}\quad H_u(2,\tfrac{1}{2},0),\ H_d(2,-\tfrac{1}{2},0) .$$

The most general gauge-invariant superpotential in this model can be written as

$$W = h_u Q H_u U^c + h_d Q H_d D^c + h_e L H_d E^c + h_\nu L H_u \nu^c + \mu H_u H_d . \tag{15}$$

We see that all terms in Eq. (15) conserve the lepton number (and R parity). Of course, the h_ν term in Eq. (15) may induce a large Dirac mass for the neutrino, which in turn may require R-parity violation[15] or lepton-number violation via new Higgs bosons to understand the small neutrino mass. R-parity- (and lepton-number) violation is not intrinsic to the theory (for instance, we could set $h_\nu = 0$ and, we will not need any lepton-number violation at all). Thus, it may be more natural to work with this kind of extended gauge models to understand the origin of lepton-number violation should it be observed. In order to obtain R-parity-violating interactions from this picture, we can give a nonzero VEV to the scalar partner of the right-handed neutrino,[15,16] i.e., $\langle \tilde{\nu}^c \rangle \neq 0$. Then, via the Feynman diagram shown in Fig. 4, interactions of type QLD^c and LLE^c are introduced. The strengths of these interactions is estimated to be the following:

$$\begin{aligned} &\text{for } QLD^c,\ \ h_d h_\nu \langle \tilde{\nu}^c \rangle / \mu ; \\ &\text{for } LLE^c,\ \ h_e h_\nu \langle \tilde{\nu}^c \rangle / \mu . \end{aligned} \tag{16}$$

Since h_ν contributes to Dirac masses of neutrinos, we expect $h_\nu \simeq (m_D / M_W)$, which for the first-generation fermions is $\approx 10^{-4}$ or so. Similarly, we expect $h_d \approx 10^{-4}$; as

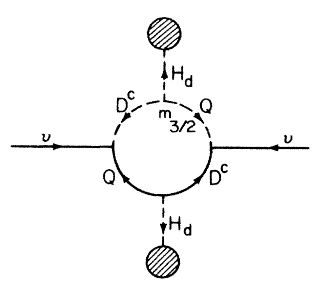

FIG. 3. One-loop diagram for neutrino mass in the presence of R-parity-violating interactions.

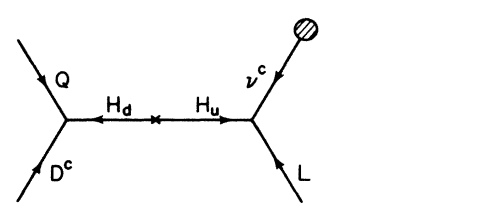

FIG. 4. Feynman diagram that induces R-parity violation at low energies.

a result, even if $\langle \tilde{\nu}^c \rangle \approx \mu$, the strength of effective R-parity-violating interactions is very weak. It is, however, easy to extend this theory further by duplicating the Higgs superfields to (H_u^i, H_d^i), $i=1,2$ with the same quantum numbers as H_u and H_d such that the strength of induced R-parity-violating coupling becomes much bigger. To see this, we write the new superpotential as

$$W=\sum_{i=1,2}(h_u^i Q H_u^i U^c + h_d^i Q H_d^i D^c + h_e^i L H_d^i E^c + h_\nu^i L H_u^i \nu^c) \ . \tag{17}$$

It is now phenomenologically acceptable for the Higgs fields ${H_u}^2$ and ${H_d}^2$ to have no VEV; the Yukawa couplings ${h_u}^2$, ${h_d}^2$, ${h_e}^2$, and ${h_\nu}^2$ can then be of order 1 leading to large induced R-parity-violating interactions at low energies.

In conclusion, we point out that in the most general supersymmetric extension of the standard model, there exist new contributions to neutrinoless double-beta decay which may be used to provide constraints on the parameters of the model. We also comment on the possible extensions of the model where such contributions are absent prior to spontaneous symmetry breaking and can arise with arbitrary strength in the low-energy sector after symmetry breaking.

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