LEPTON AND FLAVOR VIOLATION IN SUSY MODELS

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Lepton and flavor violating processes resulting from neutral scalar–lepton mixing are examined in the context of supersymmetric models. Contributions arising at the one-loop level for $\mu \to e\gamma$, $\mu \to 3e$, $\mu\mu \rightleftarrows ee$ as well as for neutrinoless $\beta\beta$-decay are found to be suppressed for a general class of supersymmetry breaking parameters.

1. Introduction. Lepton number (L) and lepton flavour (F) violating processes have not yet been observed in nature [‡1]. Even though lepton and flavour non-conservation is automatically conserved in the standard model [5], there is no satisfactory explanation for the existence of such global symmetries.

F and L violating processes could arise in the context of the standard $SU_c(3) \times SU_L(2) \times U(1)$ gauge group or its minimal parity-conserving $SU_c(3) \times SU_R(2) \times SU_L(2) \times U(1)$ extension [6] with a suitably enlarged fermionic or scalar sector. One of the available possibilities is the extension of the fermionic sector by introducing a gauge singlet field N, the well-known right-handed neutrino [7]. An analogous extension of the bosonic sector could involve the introduction of flavour mixing isodoublets or couplings to other representations which are compatible with renormalizability and gauge symmetry [8]. In this letter we will examine the implications on L and F conservation of supersymmetric extensions of the standard model.

In the supersymmetric generalizations of the standard model [‡2] one writes the Yukawa couplings (for notation see table 1)

$$V = \mathbf{Q\bar{H}U^c} + \mathbf{LHE^c} + \mathbf{QHD^c}. \tag{1.1}$$

These, however, are not the most general trilinear couplings which are consistent with the $SU(3) \times SU(2) \times U(1)$ gauge symmetry. The following trilinear Yukawa couplings which violate lepton number are also possible

$$\mathbf{QLD} + \mathbf{LE^cL}. \tag{1.2}$$

The naive inclusion of these terms into the superpotential (1.1) will lead to disastrous consequences. More precisely, with such terms we have to face L violation [‡3] at the level of renormalizable interactions unless we impose symmetries by hand, to prevent them. The usual symmetries one imposes to avoid problems like those mentioned above, are the so-called R-symmetries [10]. These symmetries unlike the case with the usual global symmetries do not necessarily commute with the supersymmetry generators. Nevertheless, one could imagine special kinds of those symmetries that might allow explicit lepton violating terms which under special combinations of leptons and higgses would lead to L violation [11] in accordance with the experimental limits. Other possibilities related either with the spontaneous violation of R-symmetry [12], or with other mechanisms [13] [‡4], of course, could lead to lepton-number violation.

‡1 For limits on the branching ratios of the various L and F violating processes see ref. [1] for $\mu \to 3e$, $R < 2.4 \times 10^{-11}$, ref. [2] for $\mu \to e$, $R < 1.7 \times 10^{-12}$, ref. [3] for $\mu^- \to e^-$ conversion, $R < 1.6 \times 10^{-11}$, and ref. [4] for $\mu^- \to e^+$, $R < 9 \times 10^{-10}$.

‡2 For a review of supersymmetric models see ref. [9].

‡3 We may also have baryon-number violation if we include the term $\mathbf{U^cD^cD^c}$. Combining this term with the $\mathbf{LE^cL}$ we have fast proton decay [9].

‡4 For a review on L violation in SUSY see ref. [14].

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Table 1
The various particles which appear in the supersymmetric extension of the standard model.

Name	Spin	Symbol	Name	Spin	Symbol
gluons	1	g^i, $i=1,\dots,8$	gluinos	1/2	$\tilde{g}^i$, $i=1,\dots,8$
gauge bosons	1	$W^{\pm}, Z$	gauginos	1/2	$\tilde{W}^{\pm}, \tilde{Z}^0$
photon	1	γ	photino	1/2	$\tilde{\gamma}$
quarks	1/2	$Q=\begin{pmatrix} u^i \\ d^i \end{pmatrix}_L$	squarks	0	$\tilde{Q}=\begin{pmatrix} \tilde{u}^i \\ \tilde{d}^i \end{pmatrix}$
		u_L^{ic}, d_L^{ic} $i=1,2,3$			$\tilde{u}_i^c$, $\tilde{d}_i^c$
leptons	1/2	$L=\begin{pmatrix} \nu_i \\ e_i^- \end{pmatrix}_L$	sleptons	0	$\tilde{L}=\begin{pmatrix} \tilde{\nu}_i \\ \tilde{e}_i^- \end{pmatrix}$
		e_{iL}^c $i=1,2,3$			$\tilde{e}_i^c$
higgses	0	$H=\begin{pmatrix} h_i^+ \\ h_i^0 \end{pmatrix}$, $i=1,2$	higgsinos	1/2	$\tilde{H}=\begin{pmatrix} \tilde{h}_i^+ \\ \tilde{h}_i^0 \end{pmatrix}$, $i=1,2$
		$\bar{H}=\dots$			$\tilde{\bar{H}}=\dots$

In what follows we discuss the problem of L and F violation, resulting from neutral scalar–lepton mixing in the context of the supersymmetric version of the standard model. Using a suitable supersymmetry breaking potential we derive the slepton mass matrix and subsequently we calculate the relevant Feynman graphs for the various violating processes.

2. *SUSY breaking and the neutral scalar boson mass matrix.* In the supersymmetric version of the standard $SU(3)_c \times SU(2)_L \times U(1)$ model [‡5] the matter fields are contained in chiral supermultiplets composed of left-handed Weyl spinors and complex scalar bosons. The Yukawa lagrangian and the non-gauge part of the scalar potential are derivable from the superpotential

$$W = a_{ij}\mathbf{L}_i^{(0)}(\mathbf{E}_j^c)^{(0)}\mathbf{H} + b_{ij}\mathbf{L}_i^{(0)}\mathbf{N}_j^{(0)}\bar{\mathbf{H}} + \tfrac{1}{2}M_{ij}^{(0)}\mathbf{N}_i^{(0)}\mathbf{N}_j^{(0)} + \dots, \tag{2.1}$$

which is an analytic function of the chiral superfields

$$\mathbf{L} = (e_L, \tilde{e}; -(\nu_L, \tilde{\nu})), \quad \mathbf{E}^c = (e_L^c, \tilde{e}^c), \quad \mathbf{H} = (\tilde{H}_L, H), \quad \bar{\mathbf{H}} = (\tilde{\bar{H}}_L, \bar{H}).$$

The mass part of the Yukawa lagrangian can be derived from the effective superpotential

$$W = m_{ij}\mathbf{e}_i\mathbf{e}_j^c + m_{ij}^D \mathbf{\upsilon}_i \mathbf{N}_j + \tfrac{1}{2}M_{ij}\mathbf{N}_i\mathbf{N}_j. \tag{2.2}$$

The Yukawa lagrangian will be

$$L_Y = \sum_{i,j} m_{ij}\, e_{iL} e_{jL}^c + \tfrac{1}{2}(\nu_i, N_i)_L \begin{pmatrix} 0 & m_{ij}^{(D)} \\ m_{ij}^{(D)} & M_{ij} \end{pmatrix} \begin{pmatrix} \nu_{iL} \\ N_{iL} \end{pmatrix} + \text{h.c.}, \tag{2.3}$$

‡5 See table 1.

while the scalar potential [‡6]

$$V=\tilde{e}^*m^+m\tilde{e}+\tilde{e}^{*c}m^+m\tilde{e}^c+\tilde{N}^*m^{+(D)}m^{(D)}\tilde{N}+|m^{(D)}\tilde{\nu}+M\tilde{N}|^2. \tag{2.4}$$

$\tilde{e}, \tilde{e}^c, \tilde{N}, \tilde{\nu}$ etc. are vectors in 3D flavor space and m, m^D, M are 3×3 matrices. The charged boson mass matrix is m^+m while the neutral boson mass matrix is

$$\begin{pmatrix}0 & m^{(D)}\\ m^{(D)} & M\end{pmatrix}^+\begin{pmatrix}0 & m^{(D)}\\ m^{(D)} & M\end{pmatrix}=\begin{pmatrix}m^{+(D)}m^{(D)} & m^{+(D)}M\\ M^+m^{(D)} & m^{+(D)}m^{(D)}+M^+M\end{pmatrix},$$

as required by supersymmetry.

Supersymmetry breaking will modify the above picture. In a large class of models with soft supersymmetry breaking [15,16] emerging from spontaneously broken supergravity [17,18], the fermions do not receive supersymmetry breaking masses at the tree level while the bosons do. The supersymmetry breaking part of the scalar potential is in general of the form

$$\delta V=m_{3/2}^2\sum_i|\varphi_i|^2+\bar{A}m_{3/2}(W+W^*)+\bar{B}m_{3/2}\sum_i(\varphi_i\partial W/\partial\varphi_i+\text{h.c.}). \tag{2.5}$$

The mass scale $m_{3/2}$ is the characteristic supersymmetry-breaking scale identified with the gravitino mass. $\bar{A}$ and $\bar{B}$ are dimensionless parameters depending on the details of physics at very high energies ($O(M_{\text{Planck}})$) which are of order unity.

Alternative ways to break supersymmetry which cannot be cast in the form (2.5) exist but they are generally ad hoc and as a rule require the introduction of many new fields which lead to phenomenologically uncomfortable situations. In our case the most general supersymmetry-breaking potential which can result from supergravity, assuming that the supersymmetry breaking is family blind, is

$$\delta V=m_{3/2}^2(|\tilde{e}|^2+|\tilde{e}^c|^2+|\tilde{\nu}|^2+|\tilde{N}|^2)+\bar{A}m_{3/2}(\tilde{e}m\tilde{e}^c+\text{h.c.})+Am_{3/2}(\tilde{\nu}m^{(D)}\tilde{N}+\text{h.c.})+(B/2)m_{3/2}(\tilde{N}M\tilde{N}+\text{h.c.}). \tag{2.6}$$

The charged-boson mass-matrix-squared in the basis $\tilde{e}, \tilde{e}^c, \tilde{e}^*, \tilde{e}^{*c}$ is just the 6×6 matrix

$$\begin{pmatrix}m^+m+m_{3/2}^2 & A^*m_{3/2}m^+\\ Am_{3/2}m & m^+m+m_{2/3}^2\end{pmatrix}. \tag{2.7}$$

The neutral-boson mass-matrix-squared in the basis $\tilde{\nu}, \tilde{N}, \tilde{\nu}^*, \tilde{N}^*$, is the 12×12 matrix

$$\left(\begin{array}{cc|cc} \multicolumn{2}{c|}{\multirow{2}{*}{$\mathcal{M}^+\mathcal{M}+m_{3/2}^2$}} & 0 & A^*m_{3/2}m^{+(D)}\\ & & A^*m_{3/2}m^{+(D)} & (B^*/2)M^+m_{3/2}\\ \hline 0 & Am_{3/2}m^{(D)} & \multicolumn{2}{c}{\multirow{2}{*}{$\mathcal{M}^+\mathcal{M}+m_{3/2}^2$}}\\ Am_{3/2}m^{(D)} & (B/2)Mm_{3/2} & & \end{array}\right),$$

[‡6] L_γ and V are constructed from (2.2) according to: $L_\gamma=\Sigma_{i,j}(\partial W/\partial\varphi_i\partial\varphi_j)\psi_{iL}\psi_{jL}+\text{h.c.}$, and $V=\Sigma_i|\partial W/\partial\varphi_i|^2$, where the sum runs over all chiral superfields $(\psi_L,\varphi)_i$.

with $\mathcal{M}$ the fermion mass matrix,

$$\mathcal{M} = \begin{pmatrix} 0 & m^{(\mathrm{D})} \\ m^{(\mathrm{D})} & M \end{pmatrix} .$$

3. L and F violation in SUSY models. The new particles present in the supersymmetric extension of the standard model contribute to L and F violating processes via diagrams such as shown in figs. 1–3. The diagrams of fig. 1 involve slepton mixings of the type $\varphi^*\varphi$ which are not influenced by the supersymmetry breaking. In contrast, the diagrams of fig. 2 which involve neutral slepton mixing of the type $\varphi\varphi$, do contain a supersymmetry breaking contribution at lowest level. Thus, neutral slepton mixings that break supersymmetry arise from the structure of the slepton mass matrix (2.8). Although the structure of (2.8) is in general complicated, the scales $m^{(\mathrm{D})}$ and $m_{3/2}$ are vastly different from M. Denoting with m_i the eigenvalues of light neutrinos and with M_i the eigenvalues of the heavy ones we can assume $m_i < m_{3/2} \ll M_i$. Focusing on the light s neutrino sector, we can see that the fields $\nu^{(0)}, \nu^{(0)*}$ will remain essentially uncoupled with eigenvalues

$$m_i^2 \approx m_{3/2}^2(1+\lambda_i), \quad m_i'^2 \approx m_{3/2}^2(1+\lambda_i'), \quad \lambda_i, \lambda_i' \ll 1. \tag{3.1}$$

In lepton flavour changing reactions we encounter the propagator $\tilde{\nu}_e \tilde{\nu}_\mu^*$ which is

$$\sum_i \tilde{U}_{ei}^* \tilde{U}_{\mu i} \frac{1}{p^2 - m_i^2} . \tag{3.2}$$

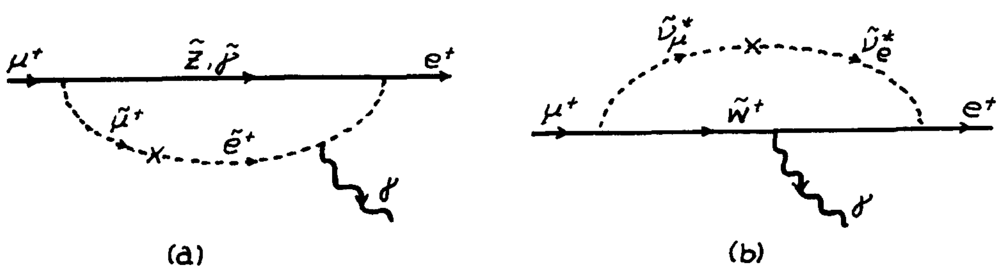


Fig. 1. Typical photonic diagrams which may lead to lepton flavor violating processes in supersymmetric theories. These mechanisms can also lead to $\mu \to e\, e^+e^-$ and (μ, e) conversion, etc., in which case photons are virtual.

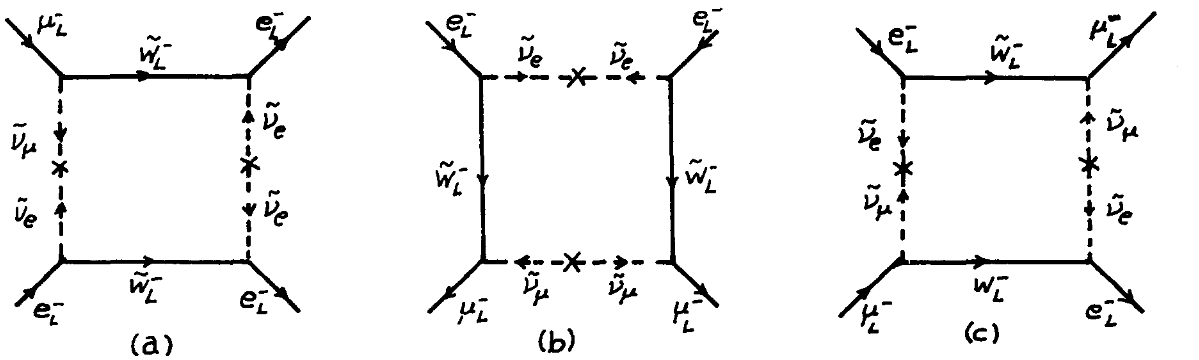


Fig. 2. Diagram (a) leads to $\mu \to 3e$ F violating process while diagrams (b) and (c) lead to $M-\overline{M}$ oscillation. The contribution of (c) is much smaller than that of (b).

With $\widetilde{U}$ we have denoted the matrix $S_e^+ S_{\widetilde{\nu}}$ where S_e is the mixing matrix for the charged leptons while $S_{\widetilde{\nu}}$ is the corresponding one for the light sneutrinos. For $\lambda_i \ll 1$ we can approximate the $\widetilde{\nu}_e \widetilde{\nu}_\mu^*$ supersymmetry-breaking contribution with

$$\xi_{\nu_e^* \nu_\mu} m_{3/2}^2/(p^2 - m_{3/2}^2)^2 , \tag{3.3}$$

where $\xi_{\nu_e^* \nu_\mu}$ is a suitably defined lepton violating parameter.

$$\xi_{\nu_e^* \nu_\mu} = \sum \widetilde{U}_{ei}^* \widetilde{U}_{\mu i} \lambda_i , \tag{3.4}$$

which is expected to be much smaller than unity. In order to see this more clearly, we make the further assumption that the matrix M is degenerate. Then [‡7] $\widetilde{U} \approx U^{(11)}$ and $\lambda_i \approx m_{\nu_i}^2/m_{3/2}^2$. Thus,

$$\xi_{\nu_e^* \nu_\mu} \approx \sum_i U_{ei}^{*(11)} U_{\mu i}^{(11)} (m_{\nu_i}^2/m_{3/2}^2). \tag{3.5a}$$

Notice that $\xi_{\nu_e^* \nu_\mu}$ here depends explicitly on the square of the neutrino mass, something which is well understood in ordinary models as well. For two generations one gets

$$\xi_{\nu_e^* \nu_\mu} \leqslant (m_{\nu_\mu}^2 - m_{\nu_e}^2)/m_{3/2}^2 \approx 10^{-20} \quad \text{for } m_{\nu_\mu} \approx 100\,\text{eV}. \tag{3.5b}$$

Such a mechanism would be unobservable (see eq. (4.4) below).

Let us consider the special case of lepton number violation for one generation. In this case, defining $\xi \equiv (m_{3/2}/M \ll 1$ and $n \equiv (m_D/m_{3/2}) \ll 1$, we find that the effective mass matrix in the light scalar neutrino will be

$$m_{3/2}^2 \begin{pmatrix} 1 & \xi\eta^2(A - B/2) \\ \xi\eta^2(A - B/2) & 1 \end{pmatrix}, \tag{3.6}$$

with eigenvalues

$$m_1^2 = m_{3/2}^2 [1 + \xi\eta^2(A - B/2)] , \quad m_2^2 = m_{3/2}^2 [1 + \xi\eta^2(A - B/2)] ,$$

and

$$\begin{pmatrix} \widetilde{\nu}^{(0)} \\ \widetilde{\nu}^{(0)*} \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & -1 \\ 1 & 1 \end{pmatrix} \begin{pmatrix} \widetilde{\nu}_1 \\ \widetilde{\nu}_2 \end{pmatrix}. \tag{3.7}$$

The $\widetilde{\nu}^{(0)} \widetilde{\nu}^{(0)*}$ propagator will be

$$\tfrac{1}{2} [1/(p^2 - m_1^2) - 1/(p^2 - m_2^2)] \simeq [\xi\eta^2(A - B/2)/(p^2 - m_{3/2}^2)^2] m_{3/2}^2 , \tag{3.8}$$

which can again be expressed in terms of the neutrino mass as

$$m_{3/2} m_\nu (A - B/2)/(p^2 - m_{3/2}^2).$$

The propagator suppression factor at low energies will this be

$$\xi_{\widetilde{\nu}\widetilde{\nu}^*} \approx m_\nu (A - B/2)/m_{3/2} . \tag{3.9a}$$

‡7 With $U^{(11)}$ we denote the matrix $S_e^+ S^{(11)}$, where $S^{(11)}$ stands for the mixing matrix of light neutrino eigenstates. For details see, however, ref. [19].

Substituting $m_{\nu_e} \approx 10$ eV and $m_{\nu_\mu} \approx 100$ eV we get

$$\xi_{\tilde{\nu}_e^* \tilde{\nu}_e} \approx 10^{-10}, \quad \xi_{\tilde{\nu}_\mu^* \tilde{\nu}_\mu} \simeq 10^{-9}. \tag{3.9b}$$

4. Examples of L and F violating processes. In order to be explicit we will examine some particular F and L violating processes in detail. Let us begin with the $\mu \to e\gamma$ process (see fig. 1). This is expected to be greatly suppressed either because the "neutral currents" (see fig. 1a) continue to be diagonal or the $\xi_{\nu_e \nu_\mu}$ violating parameter entering fig. 1b is very small (see eq. (3.5b)). We will therefore examine cases which involve lepton violating processes like $\mu \to 3e$ (fig. 2a) and M–$\bar{\text{M}}$ oscillations (figs. 2b and 2c) which are proportional to the neutrino masses. The corresponding quantities are $\xi_{\nu_e^* \nu_e}$, $\xi_{\nu_\mu^* \nu_\mu}$ and $\xi_{\nu_e^* \nu_\mu}$. The first two of these have been estimated in eq. (3.9b). The third cannot be accurately estimated since it involves in addition flavor mixing. Barring, however, some unusual circumstance we expect it to be of the same order as given by eq. (3.9b).

We will begin our discussion with the process of fig. 2a. We get

$$m = \tfrac{1}{4} g^4 \xi_{\nu_e \nu_\mu^*} \xi_{\nu_e^* \nu_e} (m_{3/2}^2 I_{\lambda\mu}/\sqrt{2})(1 - P_{12}) \bar{u}(p_1) \gamma^\lambda (1 + \gamma_5) u(p_\mu) \bar{u}(p_e) \gamma_\mu (1 - \gamma_5) u(p_2). \tag{4.1}$$

In the limit of small external momenta the loop integral $I_{\lambda\mu}$ is

$$I_{\lambda\mu} \approx \int \frac{d^4 k}{(2\pi)^4} \frac{k_\lambda k_\mu}{(k^2 - m_{\tilde{W}}^2)^2} \frac{1}{(k^2 - m_{3/2}^2)^3} = \frac{g_{\lambda\mu}}{128\pi^2} \frac{\alpha^2 + 4\alpha - 5 - 2(2\alpha + 1)\ln\alpha}{(1-\alpha)^4} = \frac{g_{\lambda\mu}}{128\pi^2} J(\alpha), \tag{4.2}$$

with $\alpha \equiv m_{3/2}^2/m^2$.

Thus the branching ratio is computed to be

$$R \simeq \{[J(\alpha)/8\pi^2] G_F m_W^4 m_{3/2}^2/m_{\tilde{W}}^4\}^2 |\xi_{\nu_e^* \nu_\mu} \xi_{\nu_e^* \nu_e}|^2. \tag{4.3}$$

For $m_{\tilde{W}} \approx m_{3/2} \approx O(m_W)$ (4.3) becomes

$$R \simeq [(1/24\pi^2) G_F m_W^2]^2 |\xi_{\nu_e^* \nu_e} \xi_{\nu_e^* \nu_\mu}|^2 \simeq 10^{-7} |\xi_{\nu_e^* \nu_\mu} \xi_{\nu_e^* \nu_e}|^2. \tag{4.4}$$

Since $|\xi_{\nu_e^* \nu_\mu}|$ and $|\xi_{\nu_e^* \nu_e}| \ll 1$ as we can see from (3.5), these scalar neutrino contributions to $\mu \to 3e$ decay are irrelevant.

Next, let us examine muonium–antimuonium oscillations. These, have been previously considered by Halprin [10]. The relevant diagram is shown in fig. 2b. Proceeding as above we find that

$$m = (G_F/\sqrt{2}) C \bar{u}(p_1)(1 - \gamma_5) u(p_2) \bar{u}(q_1)(1 + \gamma_5) u(q_2),$$

with

$$C = (1/64\pi^2) G_F m_W^2 \xi_{\nu_e^* \nu_e} \xi_{\nu_\mu^* \nu_\mu} (m_{3/2}^2 m_W^2/m_{\tilde{W}}^4) \hat{J}. \tag{4.5}$$

The loop integral $\hat{J}$ is

$$\hat{J} = \tfrac{1}{2}(1-\alpha)^{-4} [1 + 9\alpha - 9\alpha^2 + 6\alpha(1+\alpha)\ln(\alpha)],$$

i.e. $\hat{J} = \frac{1}{2}$ for $\alpha \ll 1$, $\frac{1}{2}(\alpha \sim 1)$ and $1/2\alpha^5$ for $\alpha \gg 1$.

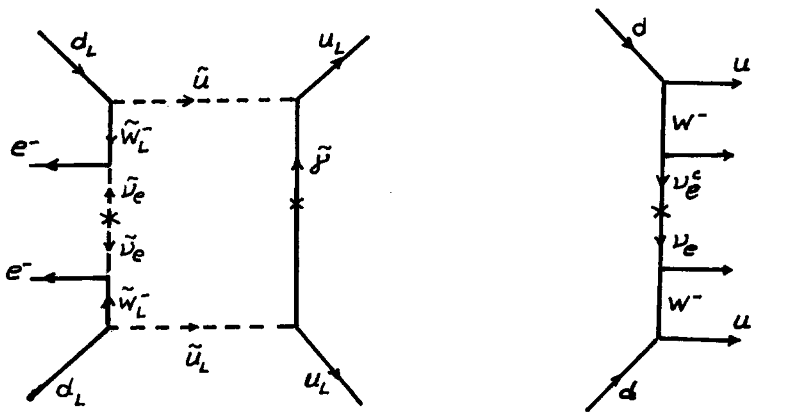
Taking $m_{3/2} \approx m_W \approx m_{\tilde{W}}$ we get

$$C \approx 10^{-4} \xi_{\nu_e^* \nu_e} \xi_{\nu_\mu^* \nu_\mu}, \tag{4.6}$$

which, with the estimates of eq. (3.9b) gives a very small value ($O(10^{-23})$) for C, which leads to an unobservable rate for this process. The contribution of fig. 2c is expected to be even smaller.

In fig. 3 we show typical diagrams leading to neutrinoless double beta decay. In the supersymmetric case the amplitude takes the form

$$\tfrac{1}{4} g^4 (4\pi\alpha) 2 m_{\tilde{\gamma}} [\xi_{\nu_e^* \nu_e} m_{3/2}^2/(m_{\tilde{W}})^4 (m_{\tilde{u}})^4] u(q_1)(1 + \gamma_5) e(q_2) \bar{u}(p_2') u(p_2) \bar{u}(p_1)(1 - \gamma_5) u(p_2). \tag{4.7}$$



(a) (b)

Fig. 3. (a) is a typical diagram at the quark level which can lead to neutrinoless double β-decay in supersymmetric theories. Note that one needs the coupling $\nu_e \nu_e^*$ which is the analog of the Majorana mass term of the usual 0ν $\beta\beta$-decay process. (b) is the corresponding diagram for the $\beta\beta$-decay in the non-supersymmetric case.

We have already made a Fierz transformation to write the amplitude in the usual form in which d-quarks annihilate and the two u-quarks are created at the same point. The loop integral is

$$I = \frac{m_{\tilde{W}}^4 m_{\tilde{u}}^4}{4} \int \frac{d^4 k}{(2\pi)^4} \frac{k^2}{(k^2 - m_{\tilde{W}}^2)} \frac{1}{(k^2 - m_{\tilde{u}}^2)^2} \frac{1}{(k^2 - m_{3/2}^2)^2} \frac{1}{k^2 - m_{\tilde{\gamma}}^2} . \tag{4.8}$$

The above amplitude, in coordinate space at the nucleon level, gives rise to a very short ranged operator analogous to that for heavy neutrino exchange i.e. [19]

$$m = \tfrac{1}{8} g^4 (\eta_N^L / m_\omega^4 m_p) \bar{u}(p_2')(1-\gamma_5) u(p_1') u(p_2)(1-\gamma_5) u(p_1) \bar{u}(q_1)(1-\gamma_5) u(q_2), \tag{4.9}$$

with

$$\eta_N^L = \sum_i U_{\mu i}^{*11} U_{ei}^{*11} e^{-i\varphi_i} m_p / M_i . \tag{4.10}$$

The quantity $\eta_N^{(L)}$ in (4.9) contains all the information about the mixing angles and neutrino mass and we call it lepton violating parameter. The φ_i are the phases which correspond to the *CP* eigenvalues of the respective mass eigenstates [19,20].

From the above we find that the corresponding lepton violating parameter in the supersymmetric case is

$$\eta_{\tilde{u}} = 8(4\pi\alpha)(m_W/m_{\tilde{W}})^4 \, [m_{3/2}^2 m_{\tilde{\gamma}} m_p / (m_{\tilde{u}}^2)^2] \, \xi_{\nu_e^* \nu_e} I . \tag{4.11}$$

To estimate the integral I we assume that $m_{\tilde{u}} \approx m_{\tilde{W}} \approx m_{3/2}$. Also taking, $m_{\tilde{\gamma}} \approx 5$ GeV [‡8] $\ll m_{\tilde{W}}$ we find $\eta_{\tilde{u}} \approx 5 \times 10^{-9}\, \xi_{\nu_e^* \nu_e}$ which must be compared with the limit $\eta_N < 10^{-7}$ [20]. Thus, 0ν $\beta\beta$-decay is unobservable for reasons of course different from those given by Halprin in ref. [13] (the s quarks need not be final particles). It appears, therefore that the standard supersymmetry breaking which is favored in the present framework of the theory, leaves no observable remnants in the leptonic sector. Our arguments apply also to $K_L \to \mu e$ and $K^+ \to \eta^+ \mu e$ decays not specifically examined in this work. The reason for such suppression is the fact that the s leptons which can reasonably admix with each other, remain almost degenerate, even though they become much heavier than

[‡8] For reasons well known [21] the lightest sparticle must be electrically neutral. Here, we adopt this value for $m_{\tilde{\gamma}}$ from ref. [21].

their corresponding leptonic partners. It is this degeneracy [22] and not the large value of their mass which is responsible for the suppression.

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