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Assignment of lepton numbers in supersymmetry

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In a supersymmetric extension of the standard $SU(3)\times SU(2)\times U(1)$ gauge model of fundamental interactions, the assignment of lepton numbers to the color-singlet fermions is not unique. In addition to the conventional assignment, there are five other generically distinct scenarios, each with three possible variations. The existing experimental data constrain the values of various new couplings present in these models.

In the standard $SU(3)\times SU(2)\times U(1)$ gauge model of fundamental interactions, the only color-singlet fermions are the left-handed doublets $(\nu_e,e)_L$, $(\nu_\mu,\mu)_L$, $(\nu_\tau,\tau)_L$, and the right-handed singlets e_R, μ_R, τ_R. Hence the neutrinos are massless and the three lepton numbers L^e, L^μ, and L^τ are separately conserved. Experimentally, all available data are consistent with this model. On the other hand, small deviations are not ruled out. For example, if each neutrino has a right-handed singlet partner and acquires a Dirac mass, then mixing will occur in general among the three leptons and only one lepton number L $(=L^e=L^\mu=L^\tau)$ is conserved. Of course, if neutrinos have Majorana masses, then even this conservation can be violated.

Now consider a supersymmetric extension[1] of the standard model. Because each color-singlet fermion doublet has a scalar partner, new Yukawa interactions involving the latter are allowed as far as gauge invariance and supersymmetry are concerned. Since these interactions do not conserve lepton number as it is conventionally defined, they are usually forbidden by the imposition of a discrete symmetry. However, this is really a theoretical presumption because, as we show below, there are many other ways to assign lepton numbers in supersymmetry so that some of these terms are allowed without running into conflict with existing experimental data. Let Φ_1 and Φ_2 be the two Higgs superfields transforming under $SU(2)\times U(1)$ as $(2,-\frac{1}{2})$ and $(2,\frac{1}{2})$, respectively, then the conventional superpotential is given by

$$W^{(1)}=h_i\Phi_1L_i\bar{E}_i+h^u_{ij}\Phi_2Q_i\bar{U}_j+h^d_{ij}\Phi_1Q_i\bar{D}_j+\mu_0\Phi_1\Phi_2\,, \quad (1)$$

where L_i and $\bar{E}_i$ $(i=1,2,3)$ are the usual left-chiral color-singlet superfields transforming as $(2,-\frac{1}{2})$ and $(1,1)$; Q_i, $\bar{U}_i$, and $\bar{D}_i$ are the color-triplet superfields transforming as $(2,\frac{1}{6})$, $(1,-\frac{2}{3})$, and $(1,\frac{1}{3})$. As it is, $W^{(1)}$ conserves baryon number as well as L^e, L^μ, and L^τ separately. However, since $L_{1,2,3}$ transform exactly as Φ_1, one might contemplate the existence of the superpotentials

$$W^{(2)}=f_eL_3L_1\bar{E}_1+f_\mu L_3L_2\bar{E}_2+f_{ij}L_3Q_i\bar{D}_j+\mu_3L_3\Phi_2\,, \quad (2)$$

$$W^{(3)}=f_{e\mu\tau}L_1L_2\bar{E}_3\,, \quad (3)$$

$$W^{(4)}=f_{e\mu}L_3L_1\bar{E}_2+f_{\mu e}L_3L_2\bar{E}_1\,, \quad (4)$$

$$W^{(5)}=f'_eL_2L_1\bar{E}_1+f'_\tau L_2L_3\bar{E}_3+f'_{ij}L_2Q_i\bar{D}_j+\mu_2L_2\Phi_2\,. \quad (5)$$

Five different models, each with one or two conserved lepton numbers, can then be defined according to the choice of superpotentials as follows:

$$\text{model 1: } W=W^{(1)}+W^{(2)}\,, \quad (6)$$

$$\text{model 2: } W=W^{(1)}+W^{(3)}\,, \quad (7)$$

$$\text{model 3: } W=W^{(1)}+W^{(2)}+W^{(3)}\,, \quad (8)$$

$$\text{model 4: } W=W^{(1)}+W^{(2)}+W^{(4)}\,, \quad (9)$$

$$\text{model 5: } W=W^{(1)}+W^{(2)}+W^{(5)}\,. \quad (10)$$

In models 1 and 2, both L^e and L^μ are conserved, but $L^\tau=0$ in model 1 and $L^\tau=L^e+L^\mu$ in model 2. In models 3–5, there is only one conserved lepton number in each case, but the assignments for e, μ, τ are different, namely, $(1,-1,0)$ in model 3, $(1,1,0)$ in model 4, and $(1,0,0)$ in model 5. For easy reference, we summarize the above in Table I. It is also clear that for each model there are three possible variations resulting from the permutations of e, μ, and τ.

Remarkably enough, these models have all been neglected until recently when model 1 with $L^\tau=0$ was proposed.[2] Note that as long as there is some kind of a conserved lepton number, R parity [as defined by $R\equiv(-1)^{2j+3B+L}$] will still be conserved; the only difference is that some of the known particles, such as τ and ν_τ, will now have odd R parity instead. In the extreme case where all terms of the form $L_iL_j\bar{E}_k$, $L_iQ_j\bar{D}_k$, and $L_i\Phi_2$ are allowed, which is equivalent to assigning $L=0$ to all color-singlet superfields, the idea of a lepton number becomes somewhat arbitrary. What is usually done is to retain the conventional definition of lepton number and simply label the offending terms lepton-number nonconserving. With this terminology, R parity is explic-

TABLE I. Lepton-number assignments (L^e,L^μ) for models 1 and 2, and L for models 3–5.

Model	e	μ	τ
1	(1,0)	(0,1)	(0,0)
2	(1,0)	(0,1)	(1,1)
3	1	−1	0
4	1	1	0
5	1	0	0

itly broken and there have been many studies[3] of such a scenario. However, in going from conserving L^e, L^μ, and L^τ separately to not conserving any of them, five possible intermediate models, each with its three variations, have been bypassed. In the following we outline some of the salient features of these five models, describe their most important phenomenological constraints, and propose future experimental tests of their possible validity.

Model 1. This has been discussed in Ref. 2. Here we summarize the main results. Since Φ_1 and L_3 transform identically, we can define without loss of generality Φ_1 to be that linear combination which picks up a nonzero vacuum expectation value and L_3 to be the orthogonal combination which does not. The term $\mu_3 L_3 \Phi_2$ in Eq. (2) is thus generally present. If it is rotated away by a redefinition of Φ_1 and L_3, then both must have nonzero vacuum expectation values and the f couplings in Eq. (2) will contribute to the fermion masses. In other words, mixing of the L_3 fermions with the usual gauge and Higgs fermions (gauginos and Higgsinos) cannot be avoided. As a result, ν_τ acquires a "seesaw" Majorana mass given by[2]

$$m_{\nu_\tau} = \frac{\mu_3^2}{2\mu_0 \tan\beta}\left[1 - \frac{\mu_0 M_1 M_2}{M_Z^2(c^2 M_1 + s^2 M_2)\sin 2\beta}\right]^{-1}, \quad (11)$$

where $s = \sin\theta_W$, $c = \cos\theta_W$, $\tan\beta = v_2/v_1$ is the ratio of the vacuum expectation values of Φ_2 to Φ_1, and $M_{1,2}$ are soft supersymmetry-breaking Majorana mass terms which preserve the gauge symmetries U(1) and SU(2), respectively. If the cosmological bound[4] of 100 eV on m_{ν_τ} is used, then μ_3 is less than a few MeV and the mixing of ν_τ and τ with the gauginos and Higgsinos is at most of the order 10^{-4} which is negligible phenomenologically. On the other hand, the photino $\tilde{\gamma}$ can now decay into τ or ν_τ which becomes the lightest particle odd under R parity. The dominant mechanism is presumably the exchange of scalar particles such as $\tilde{e}_{L,R}$, $\tilde{\mu}_{L,R}$, and $\tilde{\tau}_L$, resulting in final states such as $e^-e^+\nu_\tau$, $\mu^-\mu^+\nu_\tau$, $\tau^-e^+\nu_e$, $\tau^-\mu^+\nu_\mu$, and $\tau^- u\bar{d}$, etc.

There are constraints on the various f couplings of $W^{(2)}$ from the $K_L - K_S$ mass difference, the $K_L \to \mu^+\mu^-$ rate, the standard-model relationship between M_W and G_F, as well as the $V-A$ nature of $\mu \to e\nu\bar{\nu}$, etc. However, they are not stringent enough to forbid substantial branching fractions for the exotic processes

$$\tilde{u}_L \to d\tau^+,\ \tilde{d}_L \to d\bar{\nu}_\tau,\ \tilde{d}_R \to \begin{cases} d\nu_\tau, \\ u\tau^-, \end{cases} \quad (12)$$

and

$$\tilde{\nu}_e \to e^-\tau^+,\ \tilde{e}_L \to e\bar{\nu}_\tau,\ \tilde{e}_R \to \begin{cases} e\nu_\tau, \\ \nu_e\tau^-. \end{cases} \quad (13)$$

Hence $p\bar{p}$ (or pp) $\to \tau^\pm\tau^\pm +$ two quark jets and $ep \to \tau^-\tau^- +$ one quark jet $+$ missing energy (p_T) are possible unique signatures of this model.

Model 2. This is the simplest new model with just one term in $W^{(3)}$, resulting in the Yukawa interactions

$$L_Y = f_{e\mu\tau}\left[\tilde{\nu}_e \bar{\tau}\left(\frac{1-\gamma_5}{2}\right)\mu - \tilde{\nu}_\mu \bar{\tau}\left(\frac{1-\gamma_5}{2}\right)e + \tilde{\mu}_L \bar{\tau}\left(\frac{1-\gamma_5}{2}\right)\nu_e - \tilde{e}_L \bar{\tau}\left(\frac{1-\gamma_5}{2}\right)\nu_\mu + \tilde{\tau}_R^* \bar{\nu}_e^c\left(\frac{1-\gamma_5}{2}\right)\mu - \tilde{\tau}_R^* \bar{\nu}_\mu^c\left(\frac{1-\gamma_5}{2}\right)e\right] + \text{H.c.} \quad (14)$$

Both L^e and L^μ are conserved, but instead of $L^\tau = 0$ as in model 1, we now have $L^\tau = L^e + L^\mu$. Hence τ and ν_τ will again be odd under R parity, but they do not mix with the gauginos and Higgsinos. Phenomenologically, the quark sector is completely untouched, and of all the possible effective four-fermion interactions which can be obtained from Eq. (14), only two are easily accessible experimentally: $\mu \to e\nu_\mu\bar{\nu}_e$ through $\tilde{\tau}_R$ exchange and $e^+e^- \to \tau^+\tau^-$ through $\tilde{\nu}_\mu$ exchange. The former amounts only to a redefinition of the Fermi coupling

$$\frac{G_F}{\sqrt{2}} = \frac{g^2}{8M_W^2} + \frac{|f_{e\mu\tau}|^2}{8m_{\tilde{\tau}_R}^2}, \quad (15)$$

which puts a limit

$$|f_{e\mu\tau}| < 0.04\left(\frac{m_{\tilde{\tau}_R}}{100\ \text{GeV}}\right), \quad (16)$$

as pointed out by Barger, Giudice, and Han in Ref. 3. The latter represents a minor correction to the standard process $e^+e^- \to \gamma(Z) \to \tau^+\tau^-$ and is not useful at present for limiting $f_{e\mu\tau}$ any further.

The photino $\tilde{\gamma}$ is again unstable in this model. It will decay into $\tau^-\mu^+\bar{\nu}_e$, $\tau^-e^+\bar{\nu}_\mu$, and their conjugates. More spectacularly, the $\tilde{W}^+$ gaugino will decay into $e^+\mu^+\tau^-$ as shown in Fig. 1, as well as into $e^+\bar{\nu}_\mu\nu_\tau$, $\mu^+\bar{\nu}_e\nu_\tau$, and $\tau^+\nu_e\nu_\mu$. The branching fraction of $\tilde{W}^+ \to e^+\mu^+\tau^-$ may well be greater than 10^{-3}.

Model 3. This is obtained by combining model 1 and model 2. Hence $L^e + L^\mu$ must be zero and there is only one conserved lepton number L with the assignments 1, -1, and 0, respectively, for e, μ, and τ. In addition to all the allowed processes of models 1 and 2, there are now some interesting new ones. For example, μ^- will decay into $e^-\bar{\nu}_e\bar{\nu}_e$ and $e^-\nu_\mu\nu_\mu$ through the exchange and mixing of $\tilde{\tau}_{L,R}$; τ^- will decay into $e^-\nu_\mu\bar{\nu}_\tau$ through $\tilde{e}_L$ exchange and $\mu^-\nu_e\bar{\nu}_\tau$ through $\tilde{\mu}_L$ exchange, etc. In μ^- decay, the measured parameters[5] ρ, η, δ, and ξ are not affected by these new interactions, whereas $\xi' = 1.00$

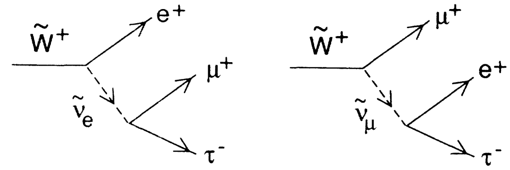

FIG. 1. Diagrams contributing to the decay $\tilde{W}^+ \to e^+\mu^+\tau^-$ in model 2.

± 0.04 implies

$$|f_{e\mu\tau}|^{1/2}(|f_e|^2+|f_\mu|^2)^{1/4}<0.3\left[\frac{m_{\rm eff}}{100\ \text{GeV}}\right], \quad (17)$$

where $m_{\rm eff}$ is the effective propagator mass from $\tilde{\tau}_{L,R}$ mixing and exchange. Similarly, the existing upper limits on $\pi^+\to\mu^+\bar{\nu}_e$ and $K^+\to\mu^+\bar{\nu}_e$ imply

$$|f_{e\mu\tau}|^{1/2}|f_{ud}|^{1/2}<0.1\left[\frac{m_{\rm eff}}{100\ \text{GeV}}\right] \quad (18)$$

and

$$|f_{e\mu\tau}|^{1/2}|f_{us}|^{1/2}<0.07\left[\frac{m_{\rm eff}}{100\ \text{GeV}}\right], \quad (19)$$

respectively. Another possible decay is $K^+\to\pi^-\mu^+e^+$. However, its amplitude is at most of the order $G_F m_K$ times that of $K^+\to\mu^+\bar{\nu}_e$; hence, its branching fraction is expected to be less than 10^{-13}, well below the present experimental upper limit of 7×10^{-9}. The same kind of suppression holds also for the conversion of μ^- into e^+ in nuclei.

Consider $\tilde{W}^+$ decay. In the conventional supersymmetric model, the final state is $\tilde{\gamma}$ + a real or virtual W^+ boson. If $e^-e^+\tau^+$ and $\mu^-\mu^+\tau^+$ are also observed, then model 1 is applicable. If $e^+\mu^+\tau^-$ is observed, then model 2 is applicable. If all three modes are observed, then model 3 is applicable. We note also that e, μ, and τ do not mix with one another in any of the three models discussed so far. Hence processes such as $\mu\to e\gamma$, $\mu\to eee$, and $\tau\to\mu ee$, etc., are all forbidden.

Model 4. This is an extension of model 1 with the constraint that $L^e=L^\mu$. Hence, there is again only one conserved lepton number L, but now with the assignments 1, 1, and 0, respectively, for e, μ, and τ. In this model, the new couplings $f_{e\mu}$ and $f_{\mu e}$ of Eq. (4) are constrained to be very small by processes such as $\mu\to eee$, μ^- conversion into e^- in nuclei, $\mu\to e\gamma$, $\mu^-\to e^-\nu_e\bar{\nu}_\mu$, $\pi^+\to\mu^+\nu_e$, $\pi^+\to\mu^-e^+e^+\nu$, $\pi^0\to e^\pm\mu^\mp$, $K^+\to\pi^+e^\pm\mu^\mp$, and $K_L\to e^\pm\mu^\mp$, etc. Consider, for example, $\mu\to eee$. The effective interaction is given by

$$H_{\rm int}=m_{\tilde{\nu}_\tau}^{-2}\left[f_{e\mu}^*f_e\bar{e}\left(\frac{1-\gamma_5}{2}\right)e\bar{e}\left(\frac{1+\gamma_5}{2}\right)\mu + f_{\mu e}f_e^*\bar{e}\left(\frac{1+\gamma_5}{2}\right)e\bar{e}\left(\frac{1-\gamma_5}{2}\right)\mu\right]. \quad (20)$$

Using $B(\mu\to eee)<1.0\times10^{-13}$, we find

$$|f_e|^{1/2}(|f_{e\mu}|^2+|f_{\mu e}|^2)^{1/4}<4.5\times10^{-4}\left[\frac{m_{\tilde{\nu}_\tau}}{100\ \text{GeV}}\right]. \quad (21)$$

Model 5. This is obtained by the lepton-number assignments 1, 0, and 0, respectively, for e, μ, and τ. The processes allowed in model 4 due to $f_{e\mu}$ and $f_{\mu e}$ are now forbidden. However, μ and ν_μ now join τ and ν_τ as nonleptonic superparticles. Whatever new interactions are allowed for τ and ν_τ in model 1 are now also allowed for μ and ν_μ. A linear combination of ν_μ and ν_τ will acquire a "seesaw" Majorana mass analogous to that of Eq. (11). The orthogonal combination will pick up a radiative mass through the exchange of two W bosons,[6] and becomes the lightest particle odd under R parity. Constraints on f_e' and f_{ij}' of Eq. (5) are similar to those on f_e and f_{ij} of Eq. (2). In addition, processes such as $\tau\to\mu\mu\mu$, $\tau\to\mu ee$, $\tau\to\mu\pi^0$, and $\tau\to\mu\gamma$, etc., constrain f_τ' and f_μ. For example, $B(\tau\to\mu\mu\mu)<2.9\times10^{-5}$ implies

$$|f_\mu|^{1/2}|f_\tau'|^{1/2}<0.09\left[\frac{m_{\tilde{\nu}_\tau}}{100\ \text{GeV}}\right], \quad (22)$$

and $B(\tau\to\mu\pi^0)<8\times10^{-4}$ implies

$$(|f_\mu|^2|f_{dd}'|^2+|f_\tau'|^2|f_{dd}|^2)^{1/4}<0.13 \quad (23)$$

for $m_{\tilde{\nu}_\mu}=m_{\tilde{\nu}_\tau}=100$ GeV.

In each of the above five models, there are clearly three distinct variations if we permute e, μ, and τ. Parallel discussions of these are straightforward and will be given elsewhere. What is important to note is that all experimental information on e, μ, τ and their neutrinos is helpful in deciding on the possible validity of these models. All rare processes are useful in setting limits on what kind of new physics will be observable beyond the standard model if supersymmetry plays a role. Whereas the standard model is very restrictive in how one may assign lepton numbers, the new particles and interactions of a supersymmetric extension allow for many possibilities and we have identified them in this paper accordingly.

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