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Supersymmetry, R-parity breaking and the neutrino magnetic moment

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Abstract

We show that supersymmetric theories with R-parity breaking naturally give rise to significant neutrino magnetic moments. In the minimal supersymmetric standard model the transition magnetic moment $\mu_{\nu_e \nu_\mu}$ can be as large as $10^{-13} \mu_B$ ($\mu_B =$ Bohr magneton). By requiring an $SU(2)_H$ horizontal symmetry of the Lagrangian in the Yukawa-less limit, $\mu_{\nu_e \nu_\mu}$ may be one order of magnitude larger. In this case the main naturalness constraint on the maximum $\mu_{\nu_e \nu_\mu}$ attainable is the electron mass value rather than the bound on the neutrino mass itself.

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1. Supersymmetry, with the introduction of quark and lepton scalars, makes it possible to write down renormalizable interactions which break lepton (L) and/or baryon (B) number. In fact, many of the phenomenological consequences of these theories, characterized by the breaking of the so-called R-parity, have already been studied in detail [1]. In this work we focus on a specific phenomenological aspect of R-parity broken supersymmetric theories with L-violating couplings: the neutrino magnetic moments.

The attention to a possible connection between the sizeable neutrino magnetic moment and R-parity breaking in supersymmetric theories has been drawn in a recent paper by Babu and Mohapatra [2]. Here we want to give a more complete analysis of the problem, in order to get the naturally attainable values of the neutrino magnetic moments in different versions of the theory. Applying a naturalness criterion, i.e. barring accidental cancellations among possible different contributions to the neutrino or the charged fermion masses, we come to the following conclusions. In the minimal supersymmetric standard model (MSSM) with R-parity broken in the superpotential, one can have neutrino magnetic moments of order $\mu \simeq 10^{-13} \mu_B$ ($\mu_B = \text{Bohr magneton}$). In this version of the theory, the main constraint on μ is set by the upper bound on the neutrino masses ($m_{\nu_e} \lesssim 10 \text{eV}$ from laboratory experiments, $m_{\nu_\mu}, m_{\nu_\tau} \lesssim 100 \text{eV}$ from cosmology). Although a value of $\mu \simeq 10^{-13} \mu_B$ is not close to the present laboratory bounds ($\mu \lesssim 10^{-10} - 10^{-9} \mu_B$ for the electron and muon neutrinos), it is far larger than the “Standard Model value” [3] of the neutrino magnetic moment $\mu/\mu_B \lesssim (10^{-19} - 10^{-18})(m_{\nu_e}/\text{eV})$. In other words, supersymmetric theories with L-breaking couplings naturally saturate the theoretical bound on the neutrino magnetic moment discussed in ref. [4].

To obtain $\mu \simeq 10^{-13} \mu_B$ in the MSSM does not require any particular symmetric structure of the Lagrangian. If such a symmetry is implemented, which requires departing from

the MSSM although still remaining with minimal particle content, an order of magnitude can be gained, leading to $\mu \lesssim 10^{-12} \mu_B$. This upper bound is obtained by applying the same naturalness criterion as before. In this case, the main constraint comes from the electron mass (or the d-quark mass) rather than the neutrino mass itself.

2. The MSSM [5] with explicit R-parity breaking via L-violation is described by the superpotential which, in addition to the standard Yukawa couplings, involves the $\Delta L \neq 0$ couplings

$$f^{\Delta L \neq 0} = \frac{1}{2} \lambda_{ijk} [L_i, L_j] e_k^c + \lambda'_{ijk} L_i Q_j d_k^c \quad (1)$$

where i, j, k are generation indices, L, Q are the lepton and quark left-handed doublets and e^c, d^c are (the charge conjugate of) the right-handed lepton and charge $-1/3$ quark singlets. The couplings induced by this superpotential are the ones that give rise to the neutrino magnetic moments, as well as to the neutrino mass. In particular, Fig. 1 shows all the one loop diagrams, arising from the purely leptonic couplings in (1), that contribute to the ν_e - ν_μ magnetic moment $\mu_{\nu_e \nu_\mu}$ after insertion of a photon vertex on any internal line. The diagrams of Fig. 1 (a),(b) are characterized by the fact that the couplings involved conserve the difference between the electron and muon lepton number $L_e - L_\mu$, whereas the diagrams of Fig. 1 (c),(d) respect L_τ . Analogous diagrams exist, originating from the couplings in (1) involving the quark fields, with the internal leptons (sleptons) replaced by quarks (squarks). In all of these diagrams, that contribute as well to the mass term $m_{\nu_e \nu_\mu}$ if no photon vertex is inserted, an helicity flip on the internal fermion line is necessary. As explicitly indicated, this also requires a mixing of the scalar leptons associated with the different chiralities, described by the scalar Lagrangian term $\Delta_{ij} \tilde{l}_i \tilde{l}_j^c$.

In a general softly broken supersymmetric lagrangian, the mixing square mass matrix

Δ_{ij} receives two distinct contributions: one from a supersymmetric coupling, $\mu(v_2/v_1)m_i\delta_{ij}$ in the standard notation [5] and in the physical basis for the leptons of mass m_i , and another from a soft supersymmetry-breaking term Δ_{ij}^{soft} . For the purposes of the present discussion, the MSSM is characterized by the property that also Δ_{ij}^{soft} is diagonal in the physical leptons basis and proportional to the lepton masses. One can therefore write

$$\Delta_{ij}^{\text{MSSM}} = \tilde{m}m_i\delta_{ij} \quad (2)$$

where $\tilde{m}$ is a typical supersymmetric mass parameter. Furthermore, in the limit of vanishing lepton masses, the sleptons are degenerate at a mass m_0 (the splitting between left and right sleptons is irrelevant to the present discussion).

Notice that, in general, many more diagrams can be drawn, contributing to the diagonal neutrino mass terms, whereas the corresponding contributions to the diagonal neutrino magnetic moments vanish by Lorentz symmetry.

Any of the diagrams of Fig. 1 contribute to $m_{\nu_e\nu_\mu}$ and $\mu_{\nu_e\nu_\mu}$ (after insertion of the photon line) respectively as

$$\delta m_{\nu_e\nu_\mu} \simeq \frac{\lambda_a\lambda_b}{16\pi^2}m\frac{\Delta}{m_0^2} \quad (3a)$$

$$|\delta\mu_{\nu_e\nu_\mu}| = \mu_B \frac{\lambda_a\lambda_b}{2\pi^2} \frac{m_em}{m_0^2} \left(\log\left(\frac{m_0}{m}\right) - 1 \right) \frac{\Delta}{m_0^2} \quad (3b)$$

where m is the mass of the internal charged lepton, λ_a and λ_b are the appropriate couplings, and Δ is the required mixing term. Of course $m/m_0 \ll 1$ and we are also making an expansion in Δ/m_0^2 . For any of the individual corresponding contributions in (3) we have

$$|\delta\mu_{\nu_e\nu_\mu}| = \mu_B \delta m_{\nu_e\nu_\mu} \frac{8m_e}{m_0^2} \left(\log\left(\frac{m_0}{m}\right) - 1 \right). \quad (4)$$

For each single product $\lambda_a\lambda_b$ there are two comparable graphs contributing to $\delta m_{\nu_e\nu_\mu}$ and to $|\delta\mu_{\nu_e\nu_\mu}|$. If we now require that each contribution to $\delta m_{\nu_e\nu_\mu}$ proportional to a fixed

product $\lambda_a \lambda_b$ be less than $10eV$, we get a bound on the corresponding $|\delta\mu_{\nu_e \nu_\mu}|$ with a logarithmic dependence on the masses m of the lepton exchanged. Numerically

$$|\delta\mu_{\nu_e \nu_\mu}| \lesssim \mu_B \cdot (0.2 - 0.4) \cdot 10^{-13} \left(\frac{100GeV}{m_0} \right)^2, \quad (5)$$

where the coefficient depends on which pair of graphs one is considering. From the negative experimental searches of sleptons, m_0 can be as low as $50GeV$, implying

$$|\delta\mu_{\nu_e \nu_\mu}| \leq O(10^{-13} \mu_B). \quad (6)$$

This bound can be saturated if and when the bound on the neutrino mass is saturated. In turn, for any given pair of diagrams (with definite lepton-slepton pair exchange), using (2) for the scalar mixing, this only depends on the values of the L-violating couplings λ_a, λ_b . Effects associated with the individual presence of any of these couplings have been studied in the literature [6] and the corresponding bounds obtained. The simultaneous presence of a pair of couplings required by the mass or the magnetic moment diagrams generally leads to stronger limits. They are:

i) for the couplings relevant to the diagrams of Fig. 1 (d), from the $\mu \rightarrow 3e$ process

$$\lambda_{122} \lambda_{121} \lesssim 10^{-6} \left(\frac{m_0}{100GeV} \right)^2 \quad (7)$$

ii) from the $\mu \rightarrow e\gamma$ transition, for the couplings occurring in Fig. 1 (c)

$$\lambda_{233} \lambda_{133} \lesssim 10^{-5} \left(\frac{m_0}{100GeV} \right)^2 \quad (8)$$

Both these bounds are stronger than the products of the individual limits [6] and do not allow the saturation of (5) by the corresponding diagrams. On the other hand, no limit exists on the products $\lambda_{123} \lambda_{232}$ of the couplings entering the first ones of the diagrams of

Fig. 1 (a),(b) stronger than the limit obtained requiring that this same contribution to $\delta m_{\nu_e \nu_\mu}$ be less than $10eV$, i.e.

$$\lambda_{123}\lambda_{232} \leq 4 \cdot 10^{-4} \left(\frac{m_0}{\tilde{m}} \right) \left(\frac{m_0}{100GeV} \right). \quad (9)$$

To get this bound, the form (2) for the scalar mixing term has been assumed. As a consequence, the related diagrams may give contributions to the magnetic moments of order $10^{-13} \mu_B$.

We neglected in our discussion the contribution of the second two diagrams of Fig. 1 (a),(b) since it cannot saturate the bound (6). In fact, these diagrams, which in the MSSM of eq. (2) are proportional to the electron mass, are mostly limited by the product $\lambda_{123}\lambda_{131} < 4 \cdot 10^{-3} (m_0/100GeV)^2$ of the individual bounds [6] on the relevant couplings. This ensures that their contribution to $\delta m_{\nu_e \nu_\mu}$ be an order of magnitude smaller than $10eV$.

Analogous considerations can be made for the diagrams involving the quark and squarks exchanges. Some of them may indeed give contributions comparable to the largest terms involving leptons (sleptons) only ($\delta\mu \simeq O(10^{-13} \mu_B)$). However one must be careful about the effects induced by the Cabibbo-Kobayashi-Maskawa mixings which result into the breaking, together with the R-parity breaking couplings, of all the individual lepton numbers.

3. We have seen that the MSSM may naturally give neutrino magnetic moments which far exceed the so-called “Standard Model value”, without invoking any particular symmetry structure of the Lagrangian. This suggests to ask whether a suitable departure from the MSSM may lead to even bigger neutrino magnetic moments. The role of appropriate approximate symmetries of the Lagrangian in connection with the neutrino problem has

been recently emphasized [2,7,8].

Along these lines we require that our Lagrangian, in the limit of vanishing Yukawa couplings, be symmetric under an $SU(2)_H$ horizontal symmetry acting on the first and the second generations. In connection with the neutrino magnetic moment, this same symmetry has already been invoked in a non supersymmetric context in ref. [8] and in the supersymmetric case in ref. [2]. In our view all of these models have problems with the $e - \mu$ mass splitting, which represents a violation of the $SU(2)_H$ -symmetry. The description of the $e - \mu$ mass difference by explicit different Yukawa couplings in a non-supersymmetric context looks unnatural because of the quadratic divergences induced by loops in asymmetric renormalizable operators. On the other hand, and perhaps more seriously, the breaking of the $SU(2)_H$ symmetry by soft terms [2] may account for the $e - \mu$ mass difference generated by radiative corrections only by an unnatural fine tuning, if possible at all.

Here we propose that the $SU(2)_H$ symmetry be respected in the Yukawa-less limit of the supersymmetric Lagrangian, including the soft breaking terms and the L-violating couplings, being just violated by the standard supersymmetric Yukawa couplings for the electron and the muon. For the consistency of the theory, we rely on the absence of quadratic divergences in supersymmetric theories. Small violations of the $SU(2)_H$ symmetry will also have to be present, for example, in the scalar mixing terms, but they will be controlled by the small Yukawa couplings and they will not be quadratically divergent. Departing from the MSSM, we will assume for these mixings terms

$$\Delta_{11} \simeq \Delta_{22} \gg |\Delta_{11} - \Delta_{22}| \simeq \tilde{m} m_\mu \quad (10)$$

Let us now reconsider the diagrams of Fig.1. The $SU(2)_H$ symmetry forbids the diagrams 1 (c),(d), since the individual couplings entering the diagrams violate the horizontal

$SU(2)_H$. On the other hand, the allowed diagrams 1 (a),(b) in the exact $SU(2)_H$ limit contribute to the magnetic moment (after the photon insertion), which is an $SU(2)_H$ singlet term, but do not to the mass, which is a triplet. In practical terms, a cancellation takes place in the mass diagrams 1 (a),(b) between the electron (selectron) and the muon (smuon) diagram, whereas the related magnetic moment terms add coherently. This allow to evade the neutrino mass bound as the strongest restriction on the magnetic moment. With respect to the previous considerations in the MSSM, we offer a symmetry reason for a natural cancellation among different mass terms.

The cancellation is most efficient in the diagrams 1 (a), which are the interesting source of the sizeable magnetic moment. Their contribution to $m_{\nu_e \nu_\mu}$ is

$$\delta m_{\nu_e \nu_\mu} = \frac{\lambda_{123} \lambda_{232}}{16\pi^2} m_\tau \frac{\Delta_{22} - \Delta_{11}}{m_0^2} = \frac{\lambda_{123} \lambda_{232}}{16\pi^2} m_\tau m_\mu \frac{\tilde{m}}{m_0^2} \quad (11)$$

which is safe ($\delta m_{\nu_e \nu_\mu} < 10eV$), provided

$$\lambda_{123} \lambda_{232} < 9 \cdot 10^{-4} \left(\frac{m_0}{\tilde{m}} \right) \left(\frac{m_0}{100GeV} \right). \quad (12)$$

On the other hand, their contribution to $\mu_{\nu_e \nu_\mu}$ is sensitive to $\Delta_{22} \simeq \Delta_{11}$ and not to their difference. The constraint (12) is then not enough to bound $\mu_{\nu_e \nu_\mu}$. Rather it is the contribution to the electron mass shown in Fig. 2 that provides the main restriction. From this diagram, where a neutral gaugino is exchanged, we have

$$\delta m_e = \frac{e^2}{16\pi^2} \sum_i \left(V_{1i}^2 + \frac{1}{\sin^2 \theta_W \cos^2 \theta_W} V_{2i}^2 \right) F(m_i) \Delta_{11} \quad (13)$$

where θ_W is the weak mixing angle,

$$F(m) = \frac{m}{m^2 - m_0^2} \left[1 + \frac{2m^2}{m^2 - m_0^2} \log \left(\frac{m_0}{m} \right) \right] \quad (14)$$

and m_i is the i -th eigenvalue of the neutralino mass matrix, diagonalized by the matrix V_{ij} from the usual basis $\chi_N = (\tilde{\gamma}, \tilde{Z}, \tilde{h}_1, \tilde{h}_2)$. We are, of course, using the fact that only

the gaugino couplings are relevant in Fig. 2 and not the couplings to the Higgsinos $\tilde{h}_{1,2}$. From (13,14), requiring $\delta m_e < 0.5 MeV$, $m_{\chi_N} \gtrsim 20 GeV$ and $m_0 \gtrsim 50 GeV$, we infer

$$\Delta_{11} \lesssim 100 GeV^2 \quad (15)$$

This bound can either be saturated for light gaugino masses m_i 's or for heavy ones $m_i \lesssim 1 TeV$.

The same mass insertion as in Fig. 2 enters in the τ -neutrino mass diagram of Fig. 3 as well. However, requiring $\delta m_{\nu_\tau} < 100 eV$ for consistency with cosmology and using $\Delta_{22} \simeq \Delta_{11} \lesssim 100 GeV^2$ gives $\lambda_{232} \lesssim 0.05(m_0/100 GeV)$, which is only slightly more restrictive than the limit $\lambda_{232} \lesssim 0.09(m_0/100 GeV)$ from lepton universality [6] and does not prevent the saturation of the bound (12) even taking into account $\lambda_{123} \lesssim 0.04(m_0/100 GeV)$ from charged current quark-lepton universality [6].

Putting all together, from the diagrams of Fig. 1 (a), we find, using (3b, 12, 15)

$$\delta\mu_{\nu_e\nu_\mu} \lesssim O(10^{-12} \mu_B) \quad (16)$$

We think that this is the largest value one can obtain for the neutrino magnetic moment in a completely natural way in a supersymmetric theory with minimal particle content.

We have also considered, as in the MSSM cases, the possible contribution of the quark-squark diagrams with the $SU(2)_H$ -symmetry implemented. In this case a strong constraint comes from the 1-loop contribution to the d-quark mass, involving the creation and annihilation of a gluino-sdown pair, analogous to the electron mass diagram of Fig 2. Also in view of the bounds on the squarks masses which are stronger than the corresponding ones for the sleptons, the value (16) cannot be attained by the quark-squark contributions.

After completion of this work, we have received a preprint by Babu and Mohapatra,

Ref. [9], which considers the non-minimal supersymmetric model discussed in Section 3. This model takes over the model previously discussed by the same authors [2] with the essential modification of breaking $SU(2)_H$ by the electron and the muon Yukawa couplings rather than by the scalar mass terms. In our opinion, the claim that a neutrino magnetic moment in the range $(10^{-11} - 10^{-10})\mu_B$ can be obtained, underestimates the significance of the bound coming from the electron mass, as we have discussed.

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FIGURE CAPTIONS

Figure 1: Supersymmetric contributions to $m_{\nu_e \nu_\mu}$ and to $\mu_{\nu_e \nu_\mu}$, when a photon line is inserted on each internal line.

Figure 2: Supersymmetric contribution to the electron mass.

Figure 3: Supersymmetric contribution to the tau-neutrino mass.

Figure 1.

Figure 2.

Figure 3.