

Effects of the neutrino B term on slepton mixing and electric dipole moments

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The supersymmetric standard model with right-handed neutrino supermultiplets generically contains a soft supersymmetry breaking mass term: $\delta L = B_\nu \mathbf{M} \tilde{\nu}_R \tilde{\nu}_R / 2$. We call this operator the “neutrino B term.” We show that the neutrino B term can give the dominant effects from the neutrino sector to lepton-flavor-violating processes and to lepton electric dipole moments.

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I. INTRODUCTION

The minimal supersymmetric standard model (MSSM), like the standard model itself, predicts a zero mass for neutrinos, and this is not compatible with the recent neutrino observations. One of the most promising methods to attribute a tiny but nonzero mass to neutrinos is the seesaw mechanism [1], which requires three extremely heavy right-handed neutrinos. In the MSSM with right-handed neutrino supermultiplets, the leptonic part of the superpotential is

$$W = Y_l^{ij} \epsilon_{\alpha\beta} H_1^\alpha L_{Ri}^\beta + Y_\nu^{ij} \epsilon_{\alpha\beta} H_2^\alpha \nu_{Ri} L_{Rj}^\beta + \frac{1}{2} M_{ij} \nu_{Ri} \nu_{Rj}, \quad (1)$$

where L_j^β is the supermultiplet corresponding to the doublet (ν_{Lj}, l_{Lj}) . Without loss of generality, we can rephase and rotate the fields to make the matrices Y_l^{ij} and M_{ij} real and diagonal: $Y_l^{ij} = \text{diag}(Y_e, Y_\mu, Y_\tau)$ and $M^{ij} = \text{diag}(M_1, M_2, M_3)$. In this basis, Y_ν can have off-diagonal and complex elements. Soft supersymmetry breaking terms of the Lagrangian in the context of this model can include

$$-\mathcal{L}_{\nu_R}^{\text{soft}} = (m_0^2)_j (\tilde{\nu}_R^j)^\dagger \tilde{\nu}_R^j + \left[\frac{1}{2} B_\nu M^{ij} \tilde{\nu}_R^i \tilde{\nu}_R^j + \text{H.c.} \right]. \quad (2)$$

The second term in Eq. (2), the “neutrino B term,” is a lepton-number-violating term [2] which can cause profound effects including sneutrino-antisneutrino oscillation [3,4]. The parameter B_ν is allowed to be much larger than the electroweak scale because it is associated only with $\tilde{\nu}_R$, which is an electroweak singlet. It has been shown that a nonzero neutrino B term can create neutrino mass through one-loop diagrams [3]. The upper bound on the neutrino mass can then be translated into a bound on B_ν ,

$$B_\nu < 10^3 m_0. \quad (3)$$

If B_ν is large, some new effects are expected both in the e^+e^- accelerator experiments [3] and in cosmology [4]. In particular, values of B_ν close to the saturating bound (3) can

induce observable slepton-antislepton oscillation. In this paper, we show that large values of B_ν can also affect other observables.

It is well known that nonzero flavor-number-violating slepton mass terms in the soft Lagrangian ($m_{\alpha\beta}^2 \tilde{L}_\alpha^\dagger \tilde{L}_\beta$, $\alpha \neq \beta$) can give rise to rare decays such as $(\mu \rightarrow \gamma e)$, $(\tau \rightarrow \gamma e)$, and $(\tau \rightarrow \gamma \mu)$. One way to avoid flavor changing neutral current effects is to choose the off-diagonal mass terms to be small. In fact, theories such as minimal supergravity (MSUGRA) suggest that at the grand unified theory (GUT) scale the soft supersymmetry breaking terms are flavor blind; that is, at the GUT scale

$$\begin{aligned} -\mathcal{L}_{\text{soft}} = & m_0^2 (\tilde{L}_{L\alpha}^\dagger \tilde{L}_{L\alpha} + \tilde{L}_{R\alpha}^\dagger \tilde{L}_{R\alpha} + \tilde{\nu}_{R\alpha}^\dagger \tilde{\nu}_{R\alpha} + H_1^\dagger H_1 + H_2^\dagger H_2) \\ & + \frac{1}{2} m_{1/2} (\tilde{B}^\dagger \tilde{B} + \tilde{W}^{a\dagger} \tilde{W}^a) + (b H_1 H_2 + \text{H.c.}) \\ & + a_0 (Y_l^{ij} \epsilon_{\alpha\beta} H_1^\alpha \tilde{L}_{Ri}^\beta \tilde{L}_{Lj} + Y_\nu^{ij} \epsilon_{\alpha\beta} H_2^\alpha \tilde{\nu}_{Ri} \tilde{L}_{Lj}^\beta) \\ & + \left(\frac{1}{2} B_\nu M_{ij} \tilde{\nu}_R^i \tilde{\nu}_R^j + \text{H.c.} \right) \end{aligned} \quad (4)$$

with universal m_0^2 , $m_{1/2}$, and a_0 .

The off-diagonal elements of the neutrino Yukawa coupling radiatively produce nonvanishing off-diagonal mass terms for the left-handed slepton doublet:

$$\begin{aligned} m_{(1)\alpha\beta}^2 = & - \sum_k \frac{Y_\nu^{k\alpha} (Y_\nu^{k\beta})^*}{16\pi^2} \left\{ m_0^2 \left(3 \log \left[\frac{\Lambda_{\text{GUT}}}{M_k} \right]^2 - 1 \right) \right. \\ & \left. + a_0^2 \log \left[\frac{\Lambda_{\text{GUT}}}{M_k} \right]^2 \right\}. \end{aligned} \quad (5)$$

This effect was first noticed and studied in [5] and then worked out in a series of papers (e.g., see [6]). However, the contribution to $m_{\alpha\beta}^2$ from the neutrino B term was ignored. In Sec. II, we study this effect and show that if B_ν is large, its contribution will dominate over the effects in Eq. (5).

In the MSSM with flavor blind soft supersymmetry breaking terms, in addition to the phases in the Yukawa couplings, there are two other independent CP -violating phases, usually chosen to be the phases of the a_0 and μ parameters. These phases can create electric dipole moments (EDMs) for

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charged leptons and for the neutron [7]. In the presence of the neutrino B term, there is one more phase that can also give a contribution to the EDMs of charged leptons. In Sec. III, we show that, even if at the GUT scale no A term is present ($a_0=0$), through one-loop corrections, the neutrino B term creates A terms for leptons at the electroweak scale. This effect could be the dominant term in lepton EDMs.

In Sec. IV, we explore the upper bounds on the imaginary and real parts of B_ν resulting from the upper bounds on the branching ratios of the rare decays [$\text{BR}(l_\alpha \rightarrow l_\gamma + \gamma)$] and the EDMs of the charged leptons. The main limitation will be the uncertainty in the pattern of neutrino Yukawa couplings Y_ν .

II. EFFECTS OF THE NEUTRINO B TERM ON SLEPTON MIXING

It has been shown that the off-diagonal slepton masses ($m_{\alpha\beta}^2 \tilde{L}_\alpha^\dagger \tilde{L}_\beta$, $\alpha \neq \beta$) at the one-loop level can give rise to lepton-number-violating rare decays such as ($\mu \rightarrow e \gamma$), ($\tau \rightarrow \mu \gamma$), and ($\tau \rightarrow e \gamma$) [6,8]. In the mass insertion approximation, a simplified formula can be derived [9]:

$$\text{BR}(l_\alpha \rightarrow l_\beta + \gamma) \sim \frac{\alpha^3}{G_F^2} \frac{|m_{\alpha\beta}^2|^2}{m_{SUSY}^8} \tan^2 \beta. \quad (6)$$

The upper bounds on the branching ratios of the rare decays [10] can be interpreted as bounds on the off-diagonal elements of $|m_{\alpha\beta}^2|$:

$$\begin{aligned} |m_{e\mu}^2| &< \frac{2 \times 10^{-3}}{\tan \beta} \left(\frac{m_{SUSY}}{200 \text{ GeV}} \right)^2 m_{SUSY}^2, \\ |m_{\tau\mu}^2| &< \frac{0.7}{\tan \beta} \left(\frac{m_{SUSY}}{200 \text{ GeV}} \right)^2 m_{SUSY}^2, \end{aligned} \quad (7)$$

and

$$|m_{\tau e}^2| < \frac{1}{\tan \beta} \left(\frac{m_{SUSY}}{200 \text{ GeV}} \right)^2 m_{SUSY}^2. \quad (8)$$

The next generation of experiments [11] is expected to improve Eq. (7) to

$$\begin{aligned} |m_{e\mu}^2| &< \frac{6 \times 10^{-5}}{\tan \beta} \left(\frac{m_{SUSY}}{200 \text{ GeV}} \right)^2 m_{SUSY}^2, \\ |m_{\tau\mu}^2| &< \frac{0.07}{\tan \beta} \left(\frac{m_{SUSY}}{200 \text{ GeV}} \right)^2 m_{SUSY}^2. \end{aligned} \quad (9)$$

The off-diagonal mass terms for left-handed sleptons receive a contribution from the neutrino B term through the two diagrams shown in Fig. 1. The two lepton-number-violating vertices on the neutrino line are the neutrino B term and the standard $\tilde{\nu}_R$ mass term. The neutrino A term is also needed. The amplitude corresponding to diagram (a) is equal to

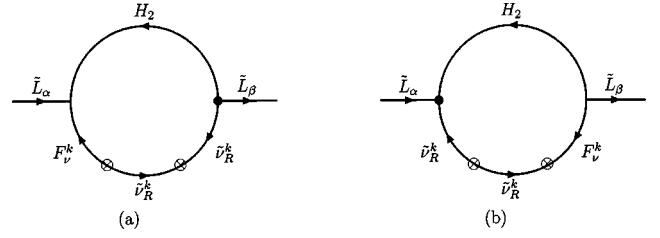


FIG. 1. Diagrams contributing to slepton masses. F_ν^k represents the auxiliary field associated with $\tilde{\nu}_R^k$. The A_ν vertices are marked with black circles.

$$\begin{aligned} -i\mathcal{M} &= \sum_k iY_\nu^{k\alpha} i(a_0 Y_\nu^{k\beta})^* \int \frac{d^4 k}{(2\pi)^4} \frac{i}{k^2} \frac{i}{k^2 - M_k^2} \\ &\quad \times (-iB_\nu M_k) \frac{i}{k^2 - M_k^2} iM_k i \\ &= \sum_k \frac{i}{(4\pi)^2} a_0^* Y_\nu^{k\alpha} (Y_\nu^{k\beta})^* B_\nu. \end{aligned} \quad (10)$$

Similarly, diagram (b) gives

$$-i\mathcal{M} = \sum_k \frac{i}{(4\pi)^2} a_0 Y_\nu^{k\alpha} (Y_\nu^{k\beta})^* B_\nu^*. \quad (11)$$

The mass correction is given by the sum of the two amplitudes:

$$m_{(2)\alpha\beta}^2 = -2 \sum_k \frac{Y_\nu^{k\alpha} (Y_\nu^{k\beta})^* \text{Re}[a_0 B_\nu^*]}{(4\pi)^2}, \quad (12)$$

which has to be added to $m_{(1)\alpha\beta}^2$ presented in Eq. (5). For $\text{Re}[a_0 B_\nu^*] \geq 10m_0^2$, $m_{(2)\alpha\beta}^2$ exceeds $m_{(1)\alpha\beta}^2$. Note that the contribution we have found does not depend on the heavy right-handed masses at all. This can be traced back to the form of the neutrino B term assumed in Eq. (4). Had we defined this term as $B_\nu^2 \tilde{\nu}_R \tilde{\nu}_R$, the result would have been proportional to $\text{Re}[a_0 (B_\nu^2)^*]/M_k$.

Up to factors of $\log(\Lambda_{GUT}/M_k)$, $m_{(1)\alpha\beta}^2$ and $m_{(2)\alpha\beta}^2$ [see Eqs. (5), (12)] have the same flavor structure. The structure can be different only if the masses of the right-handed neutrinos are hierarchical ($M_1 \ll M_2 \ll M_3$). Although the one-loop mass matrix presented in Eq. (5) is enhanced by a factor of $6 \log(\Lambda_{GUT}/M_k) \sim 10$, the neutrino B term contributions given in Eq. (12) dominate if $B_\nu \sim 10^3 m_0$ as allowed by Eq. (3).

The dependence of $m_{(2)\alpha\beta}^2$ on B_ν involves the combinations $\sum_k Y_\nu^{k\alpha} (Y_\nu^{k\beta})^*$. To derive conclusive bounds on B_ν , first we have to find some lower bounds on the $\sum_k Y_\nu^{k\alpha} (Y_\nu^{k\beta})^*$ combinations; however, this information is not available at the moment. If Y_ν are so large that $m_{(1)\alpha\beta}^2$ [see Eq. (5)] saturate the bounds (7), (8), B_ν has to be smaller than $10m_0$ [to be compared with Eq. (3)]. In Sec. IV, we will discuss this further.

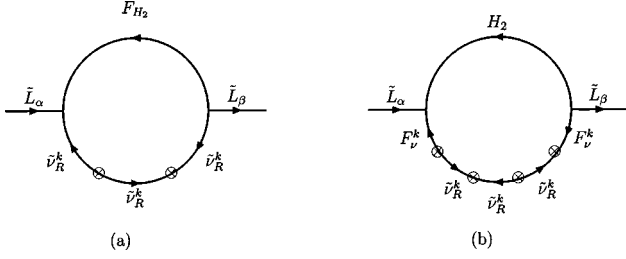


FIG. 2. Diagrams proportional to $|B|^2$ contributing to slepton masses. F_{H_2} represents the auxiliary field associated with H_2 .

In the discussion above, we have assumed $a_0 \sim m_0 \sim m_{SUSY}$ but it is possible that a_0 is much smaller than m_0 . In this case, $m_{(2)\alpha\beta}^2$ given in Eq. (12) will not be the dominant effect. At the one-loop level, there is no contribution to $m_{\alpha\beta}^2$ proportional to $|B_\nu|^2$: it can be shown that the two one-loop diagrams that are proportional to $|B_\nu|^2$ (depicted in Fig. 2) cancel each other at zero external momentum. However, at the two-loop level, there is a contribution proportional to $|B_\nu|^2$ which can dominate over $m_{(1)\alpha\beta}^2$ [Eq. (5)] provided that $|B_\nu|^2 Y_\nu Y_\nu^* / (4\pi)^2 > m_0^2$.

III. EFFECTS OF THE NEUTRINO B TERM ON A TERMS

In this section, we show that the neutrino B term creates an A term for leptons through one-loop diagrams. We then discuss how this will affect the EDMs.

As is discussed in the literature [7], the phases of μ and a_0 can create electric dipole moments for charged leptons as well as for the neutron. The current bounds on lepton EDMs are

$$d_e < 1.5 \times 10^{-27} \text{ e cm} \quad [12], \quad d_\mu < 7 \times 10^{-19} \text{ e cm} \quad [10], \quad (13)$$

and

$$d_\tau < 3 \times 10^{-16} \text{ e cm} \quad [10]. \quad (14)$$

Proposed future experiments are expected to set stronger bounds:

$$d_e < 10^{-32} \text{ e cm} \quad [13],$$

$$d_\mu < 10^{-24} \text{ } (5 \times 10^{-26}) \text{ e cm} \quad [14] \text{ } ([15]). \quad (15)$$

The bounds on the electric dipole moments of the charged leptons yield strong bounds on the imaginary parts of μ and a_0 [16].

The phase of the neutrino B term can provide yet another source of CP violation.¹ When we fixed the mass matrix M to be real, the phases of $\tilde{\nu}_R$ were fixed; therefore, the phase of B_ν in this convention cannot be removed. We expect the imaginary part of B_ν to give contribution to EDMs.

¹Within the extended MSSM, the Yukawa couplings (Y_ν) are another source of CP violation. This effect has been discussed in [17].

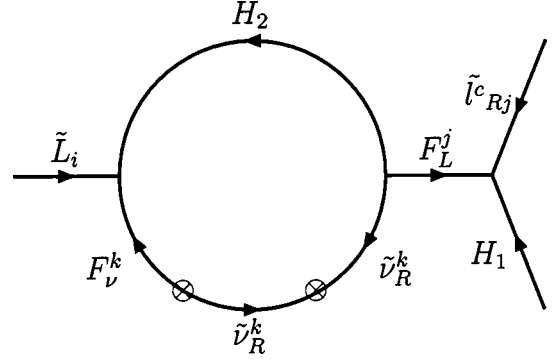


FIG. 3. Diagram contributing to A_l . F_ν^k and F_L^j represent the auxiliary fields associated with $\tilde{\nu}_R^k$ and \tilde{L}_j , respectively.

The parameter B_ν contributes to the A_l term through the diagram shown in Fig. 3. Adding this correction to the tree level A_l [see Eq. (4)], we find

$$\begin{aligned} -iA_l^{ji} &= -ia_0 Y_l^{ji} \delta_{ij} + (iY_l^{jj})(i)(iY_\nu^{kj})^*(iY_\nu^{ki}) \int \frac{i}{k^2} \frac{i}{k^2 - M_k^2} \\ &\quad \times (-iB_\nu M_k) \frac{-iM_k}{k^2 - M_k^2} \frac{d^4 k}{(2\pi)^4} \\ &= -ia_0 Y_l^{ji} \delta_{ij} - \frac{i}{(4\pi)^2} Y_l^{jj} (Y_\nu^{kj})^* Y_\nu^{ki} B_\nu. \end{aligned} \quad (16)$$

Similarly, the neutrino B term contributes to the A_ν term through the diagram shown in Fig. 4:

$$\begin{aligned} -iA_\nu^{ki} &= -ia_0 Y_\nu^{ki} + (iY_\nu^{qi})(i)(iY_\nu^{qj})^*(iY_\nu^{kj}) \int \frac{i}{k^2} \frac{i}{k^2 - M_q^2} \\ &\quad \times (-iB_\nu M_q) \frac{-iM_q}{k^2 - M_q^2} \frac{d^4 k}{(2\pi)^4} \\ &= -ia_0 Y_\nu^{ki} - \frac{i}{(4\pi)^2} Y_\nu^{qi} (Y_\nu^{qj})^* Y_\nu^{kj} B_\nu. \end{aligned} \quad (17)$$

According to [16], for $m_0 \sim 200$ GeV, the present bound ($d_e < 10^{-27}$ e cm) implies

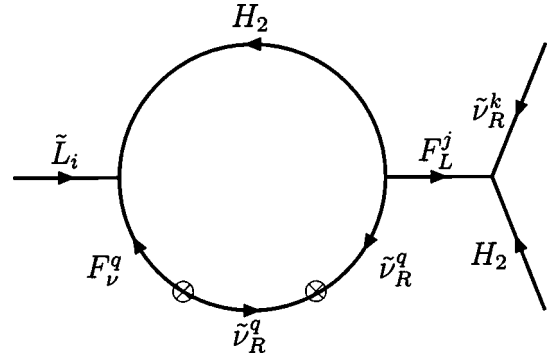


FIG. 4. Diagram contributing to A_ν . F_ν^q and F_L^j represent the auxiliary fields associated with $\tilde{\nu}_R^q$ and \tilde{L}_j , respectively.

$$\text{Im}(A_l^{ee})\langle H_1 \rangle / (m_e m_0) \lesssim 0.1. \quad (18)$$

Since the dependence of d_e on $\text{Im}(A_l^{ee})$ is linear, if in the future the bound $d_e < 10^{-32}$ e cm is obtained, the above bound will be improved to

$$\text{Im}(A_l^{ee})\langle H_1 \rangle / (m_e m_0) \lesssim 10^{-6}. \quad (19)$$

Assuming that B_ν is the only source of CP violation, the bound in Eq. (18) can be translated into

$$\text{Im}[B_\nu] \sum_k Y_l^{ee} (Y_\nu^{ke})^* Y_\nu^{ke} < 15 m_0 m_e / \langle H_1 \rangle, \quad (20)$$

which can be improved by five orders of magnitude in the future.

The present experimental data do not lead to any conclusive bounds on the values of the neutrino Yukawa couplings. In Sec. IV, we will discuss how future observations and developments can improve our knowledge on Y_ν . In principle, Y_ν can be as large as order 1. (In fact, Y_ν can be even larger than 1; however in this case we cannot treat it perturbatively.) For $Y_\nu^{ke} (Y_\nu^{ke})^* \sim 1$, the present bound (18) gives $\text{Im}[B_\nu] < 10 m_0$ [we have used Eq. (16) and $m_e = \langle H_1 \rangle Y_l^{ee}$]. Future EDM experiments can make the bound dramatically stronger.

Discovery of a lepton EDM could provide invaluable information on B_ν . The neutrino B term gives a contribution to A_l ; however, it has no impact on the A term of quarks. As a result, $\text{Im}(B_\nu)$ will not affect the EDM of the neutron. On the other hand, $\text{Im}(a_0)$ and $\text{Im}(\mu)$ give contributions to both the EDMs of charged leptons and neutrons. It is possible that the contributions of $\text{Im}(a_0)$ and $\text{Im}(\mu)$ cancel each other. However, it has been shown [18] that cancellation in the electron EDM occurs in the same regions as cancellation in the neutron EDM. Therefore, if d_e turns out to be nonzero while $d_n \ll d_e$, the effect cannot be attributed to the contribution of $\text{Im}(a_0)$ or $\text{Im}(\mu)$. Such a situation can be explained with a nonzero complex B_ν .

There is another point that is noteworthy: EDMs are proportional to $\sum_k |Y_\nu^{k\alpha}|^2$, and $\text{BR}(l_\alpha \rightarrow l_\beta \gamma)$ are given by $|\sum_k Y_\nu^{k\alpha} (Y_\nu^{k\beta})^*|^2$, which are both independent of M_l . To the author's best knowledge, there is no other observable that depends on these combinations. If $|B_\nu|$ is large, by studying these observables we can extract additional information on Yukawa couplings which will improve our current understanding of the seesaw mechanism and leptogenesis.

If the neutrino B term gives the dominant contribution to the electric dipole moments, we expect $d_\tau / (m_\tau \sum_k |Y_\nu^{k\tau}|^2) = d_\mu / (m_\mu \sum_k |Y_\nu^{k\mu}|^2) = d_e / (m_e \sum_k |Y_\nu^{ke}|^2)$; therefore, if d_e is close to its present upper bound, $d_e \sim 10^{-27}$ e cm, we expect $d_\mu \sim 10^{-25}$ e cm, which can be tested in proposed experiments [15].

IV. BOUNDS ON B_ν

In Secs. II and III, we have shown that large values of B_ν can lead to flavor-violating rare decays and EDMs of charged leptons. However, the dependence of these observables on

B_ν is through the unknown combination of Yukawa couplings $Y_\nu^{k\alpha} (Y_\nu^{k\beta})^*$. To derive upper bounds on B_ν , we have to find other observables that provide lower bounds on these combinations. In this section, we combine various pieces of information on the Yukawa couplings (some of them yet to be obtained) to derive a lower bound on the factors $Y_\nu^{k\alpha} (Y_\nu^{k\beta})^*$. We will then use the current upper bounds on the branching ratios of the rare decays and the values of EDMs to extract upper bounds on B_ν .

Neutrino masses depend on the Yukawa couplings through

$$m_{\alpha\beta}^{(\nu)} = \sum_k Y_\nu^{k\alpha} \frac{1}{M_k} Y_\nu^{k\beta} \langle H_2 \rangle^2. \quad (21)$$

All the parameters involved have to be evaluated at the electroweak scale. The effect of running from the GUT scale to the electroweak scale can change the details of the neutrino masses and mixing [19]; however, the order of magnitude of the masses will not be affected. Since, here, we want to estimate only the order of magnitude of the Yukawa couplings, we will neglect the running effects.

Currently, we have only bounds on the neutrino masses [20]:

$$\sqrt{\Delta m_{atm}^2} < \sum m_\nu < 1 \text{ eV},$$

where $\Delta m_{atm}^2 = 2.5 \times 10^{-3}$ eV² [21]. Future terrestrial and cosmological experiments will improve these bounds. Our knowledge of the masses of the right-handed neutrinos (M_k) is even less complete than the information on $m_{\alpha\beta}^{(\nu)}$. If leptogenesis is the mechanism behind the baryon asymmetry of the universe [22], it will be possible to derive a lower bound on M_k [23]. For a given value of $m_{\alpha\beta}^{(\nu)}$, there is at least one k such that

$$|Y_\nu^{k\alpha} (Y_\nu^{k\beta})^*| > \frac{1}{3} \left(\frac{m_{\alpha\beta}^{(\nu)}}{0.1 \text{ eV}} \right) \frac{2 \times 10^{-6}}{\sin^2 \beta} \left(\frac{M_k}{6 \times 10^8 \text{ GeV}} \right). \quad (22)$$

The parameter $m_{ee}^{(\nu)}$ can be extracted directly from neutrinoless double decay observations. If $m_{ee}^{(\nu)}$ is relatively large ($m_{ee}^{(\nu)} \sim 0.1$ eV), its effect should be observable in future experiments [24]. Using Eq. (16), $d_e < 10^{-27}$ e cm [see Eq. (18)] yields

$$\frac{\text{Im}(B_\nu)}{m_0} < 3 \times 10^7 \left(\frac{0.1 \text{ eV}}{m_{ee}^{(\nu)}} \right) \left(\frac{6 \times 10^8 \text{ GeV}}{M_k} \right).$$

In future, if the bound $d_e < 10^{-32}$ is obtained, this bound will be improved by five orders of magnitude [see Eq. (19)] which means it can be more restrictive than the bound in Eq. (3).

Extracting bounds on $Y_\nu^{k\alpha} (Y_\nu^{k\beta})^*$, $\alpha \neq \beta$, will be more challenging because, unlike $m_{ee}^{(\nu)}$, $m_{\alpha\beta}^{(\nu)}$, $\alpha \neq \beta$, cannot be directly measured. Even if forthcoming experiments find that the overall neutrino mass is of order of 0.1 eV (quasidegen-

erate mass scheme), it will be difficult to derive definite lower bounds on $[m_{\alpha\beta}^{(\nu)}, (\alpha \neq \beta)]$. In the case of a quasidegenerate mass scheme (the scheme for which the absolute values of the mass eigenvalues are much larger than $\sqrt{\Delta m_{\text{atm}}^2}$ with zero Majorana phases we expect $m_{e\mu}^{(\nu)}, m_{\mu\tau}^{(\nu)}, m_{e\tau}^{(\nu)} \ll m_{ee}^{(\nu)}, m_{\mu\mu}^{(\nu)}, m_{\tau\tau}^{(\nu)}$. Only in the framework of the quasidegenerate neutrino mass scheme with at least one nonzero Majorana phase, it is possible to have large off-diagonal neutrino masses $m_{e\mu}^{(\nu)}, m_{\mu\tau}^{(\nu)}, m_{e\tau}^{(\nu)} \gg \sqrt{\Delta m_{\text{atm}}^2}$ (a phase equal to π also works). On the other hand, determining the values of Majorana phases is very challenging, if possible at all [25]. Nevertheless, let us suppose that in the future some hypothetical experiment will be able to determine $m_{e\mu}^{(\nu)}$. Then, assuming that the factors $Y_{\nu}^{ke}(Y_{\nu}^{k\mu})^*$ do not cancel each other, the present bound on $|m_{e\mu}^2|$ given in Eq. (7) implies

$$\text{Re}[a_0 B_{\nu}^*]/m_0^2 < \frac{10^5}{\tan \beta} \left(\frac{m_0}{200 \text{ GeV}} \right)^2 \left(\frac{0.1 \text{ eV}}{m_{e\mu}^{(\nu)}} \right) \times \left(\frac{6 \times 10^8 \text{ GeV}}{M_k} \right). \quad (23)$$

The future possible bounds [inferred from Eq. (9)] can be more restrictive than Eq. (3):

$$\text{Re}[a_0 B_{\nu}^*]/m_0^2 < \frac{3 \times 10^3}{\tan \beta} \left(\frac{m_0}{200 \text{ GeV}} \right)^2 \left(\frac{0.1 \text{ eV}}{m_{e\mu}^{(\nu)}} \right) \times \left(\frac{6 \times 10^8 \text{ GeV}}{M_k} \right); \quad (24)$$

however, as we pointed out earlier, at the moment, measuring $m_{e\mu}^{(\nu)}$ seems to be impossible.

V. CONCLUDING REMARKS

We have studied the effects of the neutrino B term on the slepton mixing and EDMs of charged leptons in the frame-

work of seesaw model embedded in the MSSM with universal soft supersymmetry breaking terms.

If $B_{\nu} > 10m_0 \sim 10a_0$ but $a_0 > B_{\nu} Y_{\nu}(Y_{\nu})^*/(4\pi)^2$, the dominant flavor-violating slepton masses are given by Eq. (12) rather than Eq. (5). For values of B_{ν} satisfying $B_{\nu} Y_{\nu}(Y_{\nu})^*/(4\pi)^2 > a_0$ and $|B_{\nu}|^2 Y_{\nu}(Y_{\nu})^*/(4\pi)^2 > m_0^2$, the two-loop contribution proportional to $|B_{\nu}|^2$ can be dominant. The bounds on the Yukawa couplings and neutrino B term which have been discussed in the literature allow quite large contributions to the slepton masses, violating the upper bounds from rare flavor-violating decays. However, since there is no direct lower bound on the combinations of neutrino Yukawa couplings appearing in the formulations, it is not possible to derive any upper bound on the $\text{Re}[a_0 B_{\nu}^*]$.

The parameter B_{ν} can be considered as another source for CP violation and therefore EDMs. In fact, we have shown that the neutrino B term directly creates A terms for both neutrinos and charged leptons—but not quarks—even if $a_0 = 0$ at the GUT scale. The imaginary part of A_l gives a contribution to the EDMs of charged leptons. If the B_{ν} effect is dominant, we expect $d_{\tau}/(m_{\tau} \sum_k |Y_{\nu}^{k\tau}|^2) = d_{\mu}/(m_{\mu} \sum_k |Y_{\nu}^{k\mu}|^2) = d_e/(m_e \sum_k |Y_{\nu}^{ke}|^2)$; therefore, if d_e is close to its present upper bound, we expect that the proposed experiments [15] will be able to measure the value of d_{μ} . The discovery of nonzero d_e and d_{μ} while $d_n \ll d_e$ can be explained with large $\text{Im}(B_{\nu})$. In this case, $d_e \propto \sum_k Y_{\nu}^{ke}(Y_{\nu}^{ke})^* \text{Im}(B_{\nu})$. If $\text{Im}(B_{\nu})$ is determined through some other observation, information on d_e and $\text{Im}(B_{\nu})$ combined with $m_{ee}^{(\nu)} = \sum_k \langle H_2 \rangle^2 (Y_{\nu}^{ke})^2 / M_k$ (extracted from neutrinoless double beta decay searches) can provide us with information on the values of M_k , shedding light on the origins of neutrino masses and on leptogenesis.

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