Electric dipole moment of the electron in the left-right supersymmetric model

M. Frank*
Department of Physics, Concordia University, 1455 De Maisonneuve Boulevard W., Montreal, Quebec, Canada H3G 1M8
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An analysis of the supersymmetric contribution to the electron EDM in a general left-right model is given. We include one-loop contributions from the chargino, the neutralino and the doubly charged Higgsino diagrams. We discuss the dependence of the EDM on the phases of the model, as well as on masses in the left and right sectors. We show that in the unrestricted version of the model, the EDM imposes more stringent conditions on the supersymmetric spectrum for a certain range of the soft-breaking parameters, even if the right-hand scale is heavy. The electron EDM may be a clue to an extended gauge structure in supersymmetry. [S0556-2821(99)00405-1]

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I. INTRODUCTION

Large electric dipole moments for the electron and the neutron have always been of great interest because they provide information on new sources of *CP* violation and physics beyond the standard model [1]. The minimal supersymmetric standard model (MSSM), while solving the fine-tuning problem of the standard model (SM), seems to reintroduce it through the back door, through the electric dipole moments (EDMs) of the neutron and the electron. It is known that the MSSM predicts too large EDMs by about two orders of magnitude for scalar fermion masses close to the current experimental bounds and *CP* violating phases of $O(1)$. There are at present three solutions to this problem. One is to assume that supersymmetric phases are not of order unity, but rather of $O(10^{-2}-10^{-3})$ [2]. The second possibility is that the spectrum of the supersymmetric partners of quarks and leptons is heavy, i.e. of $O(1\text{ TeV})$ or more [3], and out of reach of even the CERN Large Hadron Collider (LHC). The third possibility, suggested recently, is that there are internal cancellations among the different components of the neutron EDM (the chargino and gluino contributions in particular) which can reduce the magnitude of the neutron EDM [4]. These solutions are in effect fine tuning, either for the scalar fermion masses, or for the phases, or for part of the parameter space. (The question of whether some region of parameter space will satisfy both electron and neutron EDM constraints is still open.)

Beyond the MSSM, there has been recently a renewed interest in EDMs as important signatures for supersymmetric unification [5]. In a generic supersymmetric extension of the SM, sources of *CP* violation come from either SM-type phases, such as the Cabibbo-Kobayashi-Maskawa (CKM) phase in the charged current interaction and the strong θ_{QCD} phase, or supersymmetric type, such as the phases appearing in the soft supersymmetry breaking Lagrangian and the vertex phases present in a unified supersymmetry theory [such as SU(5) or SO(10)]. Additional phases or a non-minimal particle content could provide additional sources of internal cancellations for non-minimal models such as supersymmetric grand unified theories (SUSY GUTs). Such is the case for example of theories with extra U(1) factors [6] where one has additional neutralinos. Supersymmetric theories beyond the MSSM also have additional gauginos or Higgsinos associated with the extra gauge groups and with the Higgs bosons required to break the symmetry to $SU(2)_L \times U(1)_Y$.

One of the most natural extensions of the standard model is the left-right symmetric model [7]. In addition to providing a framework for the spontaneous breaking of parity, it also provides a mechanism for a small left-handed neutrino mass (and a large right-handed neutrino mass) through the see-saw mechanism. In its supersymmetric reincarnation, it explains the absence of interactions leading to rapid proton decay without introducing some global *ad hoc* symmetry like the MSSM. It was also suggested that in certain circumstances it could cure supersymmetry of both its strong and weak *CP* problem [8]. In its most general framework, the model contains extra particles and extra phases in addition to those in the MSSM, especially in the leptonic sector. In this article we present an analysis of the electric dipole moment of the electron in a general supersymmetric left-right model. The advantage of studying the electron, rather than the neutron EDM, is that the former is free of QCD contributions, therefore offering a clear test of the electroweak sector of the model. In LR SUSY in particular, the electron EDM gets a new contribution from the doubly charged Higgsino sector, which could offer a window into a new gauge structure through a highly constrained EDM.

The paper is organized as follows: in Sec. II we describe the left-right supersymmetric model (LR SUSY). In Sec. III, we present the sources of *CP* violation in the model. We then proceed with the analytical calculation of the EDM of the electron and give the chargino, neutralino and doubly charged Higgsino contributions in Sec. IV. The numerical analysis is discussed in Sec. V, and we conclude in Sec. VI.

II. THE LEFT-RIGHT SUPERSYMMETRIC MODEL

The LR SUSY model, based on $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, has matter doublets for both left- and right-handed fermions and the corresponding left- and right-handed scalar partners (sleptons and squarks) [9]. In the

*Email address: mfrank@vax2.concordia.ca

gauge sector, corresponding to $SU(2)_L$ and $SU(2)_R$, there are triplet gauge bosons $(W^{+,-},W^0)_L,(W^{+,-},W^0)_R$ and a singlet gauge boson V corresponding to $U(1)_{B-L}$, together with their superpartners. The Higgs sector of this model consists of two Higgs bi-doublets, $\Phi_u(\frac{1}{2},\frac{1}{2},0)$ and $\Phi_d(\frac{1}{2},\frac{1}{2},0)$, which are required to give masses to both the up and down quarks. The phenomenology of the Higgs doublet is similar to the non-supersymmetric left-right model [7], except that the second pair of Higgs doublet fields, which provide new contributions to the flavor-changing neutral currents, must be heavy, in the 5–10 TeV range, effectively decoupling from the low-energy spectrum [10]. The spontaneous symmetry breaking of the group $SU(2)_R\times U(1)_{B-L}$ to the hypercharge symmetry group $U(1)_Y$ is accomplished by the vacuum expectation values of a pair of Higgs triplet fields $\Delta_L(1,0,2)$ and $\Delta_R(0,1,2)$, which transform as the adjoint representation of $SU(2)_R$. The choice of the triplets (versus four doublets) is preferred because with this choice a large Majorana mass can be generated (through the see-saw mechanism) for the right-handed neutrino and a small one for the left-handed neutrino [7]. In addition to the triplets $\Delta_{L,R}$, the model must contain two additional triplets $\delta_L(1,0,-2)$ and $\delta_R(0,1,-2)$, with quantum number $B-L=-2$ to insure cancellation of the anomalies that would otherwise occur in the fermionic sector. Given their strange quantum numbers, the δ_L and δ_R do not couple to any of the particles in the theory, so their contribution is negligible for any phenomenological studies. We list the field content of the model in Table I.

As in the standard model, in order to preserve $U(1)_{EM}$ gauge invariance, only the neutral Higgs fields acquire non-zero vacuum expectation values (VEV's). These values are

$$\langle\Delta_L\rangle=\begin{pmatrix}0 & 0\\ v_L & 0\end{pmatrix},\quad \langle\Delta_R\rangle=\begin{pmatrix}0 & 0\\ v_R & 0\end{pmatrix},$$

$$\langle\Phi\rangle=\begin{pmatrix}\kappa & 0\\ 0 & \kappa' e^{i\omega}\end{pmatrix}.$$

$\langle\Phi\rangle$ causes the mixing of W_L and W_R bosons with *CP*-violating phase ω. In order to simplify, we will take the VEV's of the Higgs fields as $\langle\Delta_L\rangle=0$ and

$$\langle\Delta_R\rangle=\begin{pmatrix}0 & 0\\ v_R & 0\end{pmatrix},\quad \langle\Phi_u\rangle=\begin{pmatrix}\kappa_u & 0\\ 0 & 0\end{pmatrix},$$

$$\langle\Phi_d\rangle=\begin{pmatrix}0 & 0\\ 0 & \kappa_d e^{i\omega}\end{pmatrix}.$$

Choosing $v_L=0$, $\kappa_d\gg\kappa_u$ satisfies the more loosely required hierarchy $v_R\gg max(\kappa,\kappa')\gg v_L$ and also the required cancellation of flavor-changing neutral currents. The Higgs fields acquire non-zero VEV's to break both parity and $SU(2)_R$. In the first stage of breaking, the right-handed gauge bosons, W_R and Z_R, acquire masses proportional to v_R and become much heavier than the usual (left-handed) neutral gauge bosons W_L and Z_L, which pick up masses proportional to κ_u and κ_d at the second stage of breaking.

TABLE I. The particle content of the left-right supersymmetric model.

Fields	Components	$SU(2)_L\times SU(2)_R\times U(1)_{B-L}$		
$Q_{L,R}$	$(u,d)_{L,R}$	$\frac{1}{2}(0)$	$0(\frac{1}{2})$	$\frac{1}{3}$
$L_{L,R}$	$(u,d)_{L,R}$	$\frac{1}{2}(0)$	$0(\frac{1}{2})$	-1
$\tilde{Q}_{L,R}$	$(\tilde{u},\tilde{d})_{L,R}$	$\frac{1}{2}(0)$	$0(\frac{1}{2})$	$\frac{1}{3}$
$\tilde{L}_{L,R}$	$(\tilde{\nu},\tilde{e})_{L,R}$	$\frac{1}{2}(0)$	$0(\frac{1}{2})$	-1
$W_{L,R}$	$(W^+,W^-,W^0)_{L,R}$	1(0)	0(1)	0
V	V	0	0	0
$\tilde{W}_{L,R}$	$(\tilde{W}^+,\tilde{W}^-,\tilde{W}^0)_{L,R}$	1(0)	0(1)	0
$\tilde{V}$	$\tilde{V}$	0	0	0
$\Phi_{u,d}$	$\begin{pmatrix}\Phi_1^0 & \Phi_1^+\\ \Phi_2^- & \Phi_2^0\end{pmatrix}_{u,d}$	$\frac{1}{2}$	$\frac{1}{2}$	0
$\Delta_{L,R}$	$\begin{pmatrix}\frac{1}{\sqrt{2}}\Delta^0 & \Delta^{++}\\ \Delta^0 & \frac{-1}{\sqrt{2}}\Delta^+\end{pmatrix}_{L,R}$	1(0)	0(1)	2
$\delta_{L,R}$	$\begin{pmatrix}\frac{1}{\sqrt{2}}\delta^0 & \delta^{++}\\ \delta^0 & \frac{-1}{\sqrt{2}}\delta^+\end{pmatrix}_{L,R}$	1(0)	0(1)	-2
$\tilde{\Phi}_{L,R}$	$\begin{pmatrix}\tilde{\Phi}_1^0 & \tilde{\Phi}_1^+\\ \tilde{\Phi}_2^- & \tilde{\Phi}_2^0\end{pmatrix}_{u,d}$	1(0)	0(1)	0
$\tilde{\Delta}_{L,R}$	$\begin{pmatrix}\frac{1}{\sqrt{2}}\tilde{\Delta}^0 & \tilde{\Delta}^{++}\\ \tilde{\Delta}^0 & \frac{-1}{\sqrt{2}}\tilde{\Delta}^+\end{pmatrix}_{L,R}$	1(0)	0(1)	2
$\tilde{\delta}_{L,R}$	$\begin{pmatrix}\frac{1}{\sqrt{2}}\tilde{\delta}^0 & \tilde{\delta}^{++}\\ \tilde{\delta}^0 & \frac{-1}{\sqrt{2}}\tilde{\delta}^+\end{pmatrix}_{L,R}$	1(0)	0(1)	-2

The supersymmetric sector of the model, while preserving left-right symmetry, has four singly-charged charginos (corresponding to $\tilde{\lambda}_L$, $\tilde{\lambda}_R$, $\tilde{\phi}_u$, and $\tilde{\phi}_d$), in addition to $\tilde{\Delta}_L^-$, $\tilde{\Delta}_R^-$, $\tilde{\delta}_L^-$ and $\tilde{\delta}_R^-$, which are presumed heavy. The model also has eleven neutralinos, corresponding to $\tilde{\lambda}_Z$, $\tilde{\lambda}_{Z'}$, $\tilde{\lambda}_V$, $\tilde{\phi}_{1u}^0$, $\tilde{\phi}_{2u}^0$, $\tilde{\phi}_{1d}^0$, $\tilde{\phi}_{2d}^0$, $\tilde{\Delta}_L^0$, $\tilde{\Delta}_R^0\tilde{\delta}_L^0$, and $\tilde{\delta}_R^0$. It has been shown that in the scalar sector, the left-triplet Δ_L couplings can be neglected in phenomenological analyses of muon and tau decays [12]. Although Δ_L is not necessary for symmetry breaking [13] and is introduced only for preserving left-right symmetry, both Δ_L^{--} and its right-handed counterpart Δ_R^{--} play very important roles in phenomenological studies of the LR SUSY model. It has been shown that these bosons, and possibly their fermionic counterparts, are light [11]. Also, these doubly charged Higgs bosons and their corresponding Higgsinos could lead to an enhancement in lepton-flavor violating decays, the anomalous magnetic moment of the muon [9] and possibly the electric dipole moment of the electron [15].

III. SOURCES OF *CP* VIOLATION IN LR SUSY

In the three-family $SU(2)_L \times U(1)_Y$ model of electroweak interactions, *CP* violation arises (apart from the θ-term in QCD) from the complex couplings of the charged weak current, i.e., the Kobayashi-Maskawa matrix V. An equivalent Kabayashi Maskawa matrix exists for the leptons in the LR symmetric model. An extra source of *CP* violation comes from phases of the complex $W_R - W_L$ transition mass term and from the complex Dirac masses of neutral leptons. The effects on the electron EDM expected to be induced from these couplings has been evaluated in non-supersymmetric LR models [14]. We shall concentrate in what follows on the phases in the supersymmetric sector. The supersymmetric sector of the model has a few interesting *CP* violating phases, arising from either the complex parameters in the superpotential, or the soft supersymmetry-breaking terms. The superpotential for the LR SUSY is

$$\begin{aligned} W = & \mathbf{h}_q^{(i)} Q^T \tau_2 \Phi_i \tau_2 Q^c + \mathbf{h}_l^{(i)} L^T \tau_2 \Phi_i \tau_2 L^c \\ & + i(\mathbf{h}_{LR} L^T \tau_2 \Delta L + \mathbf{h}_{LR} L^{cT} \tau_2 \Delta^c L^c) + M_{LR} \\ & \times [\mathrm{Tr}(\Delta \overline{\Delta}) + \mathrm{Tr}(\Delta^c \overline{\Delta^c})] + \mu_{ij} \, \mathrm{Tr}(\tau_2 \Phi_i^T \tau_2 \Phi_j) + W_{NR} \end{aligned} \tag{1}$$

where W_{NR} denotes (possible) non-renormalizable terms arising from higher scale physics or Planck scale effects [11]. The presence of these terms insures that, when the SUSY breaking scale is above M_{W_R}, the ground state is R-parity conserving [8]. The soft-breaking term is given by

$$\begin{aligned} \mathcal{L}_{soft} = & -[\mathbf{A}_q^i \mathbf{h}_q^{(i)} \tilde{Q}^T \tau_2 \Phi_i \tau_2 \tilde{Q}^c + \mathbf{A}_l^i \mathbf{h}_l^{(i)} \tilde{L}^T \tau_2 \Phi_i \tau_2 \tilde{L}^c \\ & + i \mathbf{A}_{LR}^i \mathbf{h}_{LR} (\tilde{L}^T \tau_2 \Delta \tilde{L} + \tilde{L}^{cT} \tau_2 \Delta^c \tilde{L}^c)] \\ & - [M_L \tilde{W}_L \tilde{W}_L + M_R \tilde{W}_R \tilde{W}_R + M_V \tilde{V} \tilde{V}] \\ & - M_\Delta^2 [\mathrm{Tr}(\Delta \overline{\Delta}) + \mathrm{Tr}(\Delta^c \overline{\Delta^c})] - B\mu_{ij} \Phi_i \Phi_j - \mu_{ij}^2 \Phi_i \Phi_j \end{aligned} \tag{2}$$

where $\mathbf{h}_u$, $\mathbf{h}_d$, $\mathbf{h}_\nu$ and $\mathbf{h}_e$ are the Yukawa couplings for the up and down quarks and neutrino and electron, respectively, and $\mathbf{h}_{LR}$ is the coupling for the Higgs triplet bosons. LR symmetry requires all **h**-matrices to be Hermitian in the generation space and the $\mathbf{h}_{LR}$ matrix to be symmetric. The Yukawa matrices have physical and geometrical significance and cannot be rotated away. The trilinear scalar couplings **A**-matrices (A_u, A_d, A_ν and A_e) are of a similar form to the Yukawa couplings, $B\mu_{ij}$ and μ_{ij}^2 are bilinear Higgs couplings. The soft supersymmetry breaking parameters run by the renormalization group flow. The Yukawa coupling matrices $\mathbf{h}_u$, $\mathbf{h}_d$, $\mathbf{h}_\nu$ and $\mathbf{h}_e$ lead, after diagonalization of the quark and lepton mass matrices, to the Kobayashi-Maskawa phase δ. The parameters: **A**, $B\mu_{ij}$ and μ_{ij}^2, as well as the gaugino Majorana masses M_L, M_R and M_V, are in general complex. However, by re-arrangement, not all the phases are physical, and some can be shifted into the interaction terms such that supersymmetric *CP*-violating phases can be expressed in terms of only two, $\alpha = \arg(A m_\lambda^*)$ and $\theta = \arg(\mu)$. We shall make the standard assumption of ignoring all the inter-generational mixings of leptons and sleptons; therefore we are interested only in the diagonal elements of Yukawa couplings and trilinear couplings **A**.

For simplicity we shall assume a universal form of supersymmetry breaking, namely a universal scalar mass m_0 and a trilinear scalar coupling with a universal parameter **A**.

The *CP*-violating effects arise from the slepton mass matrices which have the following form [10]:

$$\mathcal{L}_{M_{\tilde{e}}} = (\tilde{e}_L^\dagger \tilde{e}_R^\dagger) \begin{pmatrix} \mu_L^2 + m_e^2 + c_\nu h_\nu^2 & \mathcal{A}_e^* m_e \\ \mathcal{A}_e m_e & \mu_R^2 + m_e^2 + c_\nu h_\nu^2 \end{pmatrix} \begin{pmatrix} \tilde{e}_L \\ \tilde{e}_R \end{pmatrix}, \tag{3}$$

and

$$\mathcal{L}_{M_{\tilde{\nu}}} = (\tilde{\nu}_L^\dagger \tilde{\nu}_R^\dagger) \begin{pmatrix} \mu_L'^2 + m_\nu^2 + c_e h_e^2 & \mathcal{A}_\nu^* m_\nu \\ \mathcal{A}_\nu m_\nu & \mu_R'^2 + m_\nu^2 + c_e h_e^2 \end{pmatrix} \begin{pmatrix} \tilde{\nu}_L \\ \tilde{\nu}_R \end{pmatrix}, \tag{4}$$

where $\mathcal{A}_e \sim A_e - \mu \tan\beta$, $\mathcal{A}_e = |\mathcal{A}_e| \exp(i\alpha)$, $\mathcal{A}_\nu \sim A_\nu - \mu \cot\beta$, $\mathcal{A}_\nu = |\mathcal{A}_\nu| \exp(i\alpha)$, where A_e, A_ν are the trilinear scalar interaction, and $\tan\beta = \kappa_d/\kappa_u$ is the ratio of the vacuum expectation values of the Higgs bidoublet. The matrix for $\mathcal{L}_{M_{\tilde{e}}}$ differs from the corresponding one in the MSSM by the presence of the neutrino mass term in $[\mathcal{L}_{M_{\tilde{e}}}]_{22}$. The fields $\tilde{e}_L, \tilde{e}_R$ can be transformed into mass eigenstates $\tilde{e}_1, \tilde{e}_2$,

$$\tilde{e}_L = D_{e11} \tilde{e}_1 + D_{e12} \tilde{e}_2 = \cos\theta_e \tilde{e}_1 + \exp\left(-\frac{1}{2} i \phi_e\right) \sin\theta_e \tilde{e}_2,$$

$$\tilde{e}_R = D_{e21} \tilde{e}_1 + D_{e22} \tilde{e}_2 = \cos\theta_e \tilde{e}_2 - \exp\left(\frac{1}{2} i \phi_e\right) \sin\theta_e \tilde{e}_1, \tag{5}$$

where the mixing angle θ is given by

$$\tan 2\theta_e = 2 |\mathcal{A}_e| m_e / (\mu_L^2 - \mu_R^2) \tag{6}$$

and the physical masses, $M_{1,2}$, corresponding to the eigenvalues of the mass matrix in Eq. (3) are

$$\begin{aligned} M_{1,2}^2 = & \frac{1}{2} (\mu_L^2 + \mu_R^2) + 2m_e^2 + 2m_\nu^2 \\ & \pm [(\mu_L^2 - \mu_R^2)^2 + 4m_e^2 |\mathcal{A}_e|^2]^{1/2}. \end{aligned} \tag{7}$$

If L-R symmetry is conserved, then $\mathcal{A}_e = \mathcal{A}_e^*$, the masses of the left and right sleptons are equal and the contribution to the electric dipole moment is zero. After breaking, the trilinear coupling can develop an imaginary part and the mass difference between the scalar partners of the left handed and right handed leptons is much larger than the mass difference between the generations. We expect $(m_{\tilde{e}_1}^2 - m_{\tilde{e}_2}^2)/m_{\tilde{e}}^2 \sim 10^{-2}$

-10^{-1} [8,10]. Since $m_e \ll m_{\tilde{e}}$, and we can approximate reasonably $A_e - \mu \tan\beta \sim m_{\tilde{e}}$, the mixing angle is given by $\sin\theta \sim m_e(A_e - \mu\tan\beta)/m_{\tilde{e}}^2 \sim m_e/m_{\tilde{e}}$. The same discussion of mixings and splittings is valid for scalar neutrinos, i.e., Eqs. (5)–(7) are valid, with the replacement $\nu \leftrightarrow e$. Note that in the LR model, the neutrino mass is allowed to be nonzero, but can be made small through the see-saw mechanism, as long as the right-handed neutrino is very heavy (masses of order 1 TeV or so are consistent with the upper limit on the right-handed electron neutrino mass [7]). Despite the presence of the two scalar neutrinos, the mixing between the right-handed and the left-handed sneutrinos is small, due to the see-saw mechanism in the sfermion sector. The left-right elements of the sneutrino mass matrix are proportional to the Dirac neutrino mass, which can be significant. But the right-right element of the sneutrino mass matrix is very heavy, so the mixing of sneutrino will be suppressed by the inverse M_R^2. The only place where the right-handed scale could have a measurable effect is in the doubly charged Higgsino contribution. We shall proceed now to evaluate the contributions of these phases to the EDM of the electron.

IV. THE ELECTRIC DIPOLE MOMENT

The electric dipole moment of an elementary fermion is defined through its electromagnetic form factor $F_3(q^2)$ found from the (current) matrix element:

$$\langle f(p')|J_\mu(0)|f(p)\rangle = \overline{u}(p')\Gamma_\mu(q)u(p), \tag{8}$$

where $q = p' - p$ and

$$\Gamma_\mu(q) = F_1(q^2)\gamma_\mu + F_2(q^2) i\sigma_{\mu\nu}q^\nu/2m + F_A(q^2) \times(\gamma_\mu\gamma_5 q^2 - 2m\gamma_5 q_\mu) + F_3(q^2)\sigma_{\mu\nu}\gamma_5 q^\nu/2m, \tag{9}$$

with m the mass of the fermion. The EDM of the fermion field f is then given by

$$d_f = -F_3(0)/2m, \tag{10}$$

corresponding to the effective dipole interaction

$$\mathcal{L}_I = -\frac{i}{2} d_f \overline{f}\sigma_{\mu\nu}\gamma_5 f F^{\mu\nu}. \tag{11}$$

In the static limit this corresponds to an effective Lagrangian $\mathcal{L}_i = d_f \Psi_A^+ \vec{\sigma}\cdot\vec{E}\Psi_A$, where Ψ_A is the large component of the Dirac field. The effective Lagrangian is induced at one-loop level if the theory contains a *CP*-violating coupling at tree level. We can parametrize the interaction of a fermion Ψ_f with other fermions Ψ_i-s and scalars Φ_k-s with charges Q_i, Q_k in general as

$$-\mathcal{L}_{int} = \sum_{ik} \overline{\Psi}_f \left(A_{ik}\frac{1-\gamma_5}{2} + B_{ik}\frac{1+\gamma_5}{2}\right)\Psi_i\Phi_k + \text{H.c.} \tag{12}$$

If there is *CP* violation, then $\text{Im}(A_{ik}B_{ik}^*) \neq 0$, and the one-loop fermion EDM is given by

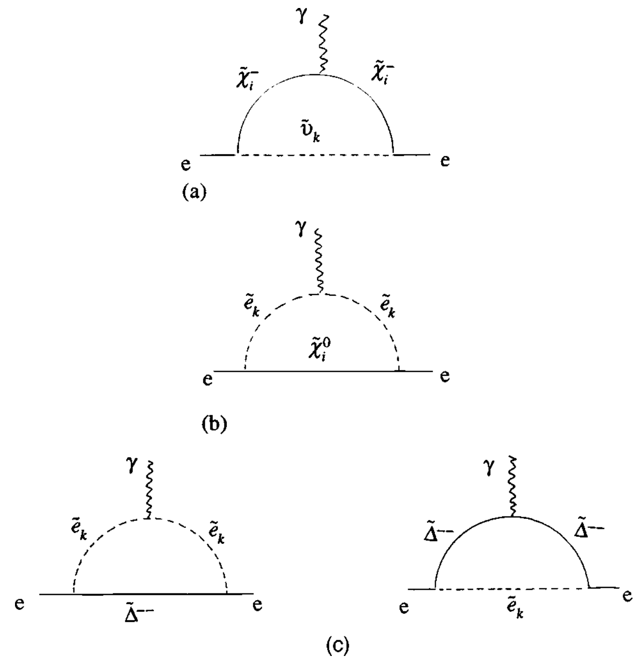

FIG. 1. (a) One-loop chargino contributions to the electric dipole moment of the electron. Here $\tilde{\chi}_i^\pm$ represents a chargino state and i runs from 1 to 4. (b) One-loop neutralino contributions to the electric dipole moment of the electron. Here $\tilde{\chi}_i^0$ represents a neutralino state and i runs from 1 to 6. (c) One-loop doubly-charged Higgsino contributions to the electric dipole moment of the electron. Here $\tilde{\Delta}_{L,R}^{--}$ represents a doubly-charged Higgsino state.

$$d_f^E = \sum_{ik} \frac{m_i}{(4\pi)^2 m_k^2}\text{Im}(A_{ik}B_{ik}^*)\left(Q_i J\left(\frac{m_i^2}{m_k^2}\right) + Q_k I\left(\frac{m_i^2}{m_k^2}\right)\right) \tag{13}$$

where the loop functions $I(r)$ and $J(r)$ are

$$I(r) = \frac{1}{2(1-r)^2}\left(1 + r + \frac{2r\ln r}{1-r}\right) \tag{14}$$

and

$$J(r) = \frac{1}{2(1-r)^2}\left(3 - r + \frac{2\ln r}{1-r}\right), \tag{15}$$

assuming charge conservation at the vertices $Q_k = Q_f - Q_i$. Since a non-vanishing d_f in the SM results in fermion chirality flip, it requires both *CP* violation and $SU(2)_L$ symmetry breaking. This can occur at the one-loop level if (1) either one vertex is a gauge-type, the other one a Yukawa type (and the fermion in the loop is a mixed gaugino-Higgsino state); or (2) the vertices come both from gauge or Yukawa type interactions, but the slepton states are mixed ($\tilde{l}_L$ and $\tilde{l}_R$).

The contributions to the one-loop fermion EDM in LR SUSY are shown in Fig. 1. In the subsequent sections, we shall analyze each of these contributions in turn.

A. The chargino contribution

The chargino contribution is shown in Figure 1(a). The electroweak gauginos and Higgsinos are all spin-1/2 weakly interacting charged particles which mix once the symmetry is broken. In the left-right supersymmetric model, the chargino matrix is a 4×4, non-symmetric, non-Hermitian matrix, M^c, from the Lagrangian:

$$\mathcal{L}_{ch}=-\frac{1}{2}(\Psi^+\Psi^-)\begin{pmatrix}0 & M^{cT}\\ M^c & 0\end{pmatrix}\begin{pmatrix}\Psi^+\\ \Psi^-\end{pmatrix}+\text{H.c.} \quad (16)$$

where

$$\Psi^+=(-i\lambda_L^+,-i\lambda_R^+,\tilde{\Phi}_u^+,\tilde{\Phi}_d^+); \quad (17)$$

$$\Psi^-=(-i\lambda_L^-,-i\lambda_R^-,\tilde{\Phi}_u^-,\tilde{\Phi}_d^-); \quad (18)$$

and

$$M^c=\begin{pmatrix} M_L & 0 & 0 & \sqrt{2}M_W\sin\beta\\ 0 & M_R & 0 & \sqrt{2}M_W\sin\beta\\ \sqrt{2}M_W\cos\beta & \sqrt{2}M_W\cos\beta & 0 & \mu\\ 0 & 0 & \mu & 0\end{pmatrix} \quad (19)$$

where $\tan\beta=\kappa_d/\kappa_u$, and M_L, M_R are the gaugino masses in the left- and right-handed sector. We have neglected here the contributions of the $\tilde{\Delta}_{L,R}^+$ Higgsinos since we expect them to be heavy and decouple from the low-lying spectrum.

The matrix M^c is not real because in general the gaugino masses are complex and so is μ. Re-defining the phases, one is still left with the phase of μ and the mixing phase between W_L and W_R from

$$W_1\simeq W_L=\cos\zeta W_L+e^{i\omega}\sin\zeta W_R$$

$$W_2\simeq W_R=-e^{i\omega}\sin\zeta W_L+\cos\zeta W_R \quad (20)$$

where experimental considerations restrict $0.005\leqslant\zeta\leqslant0.015$. As in the MSSM, we need two unitary matrices, U and V, to diagonalize M^c:

$$M_D=U^*M^cV^{-1}. \quad (21)$$

The eigenvalues of M^c can be either positive or negative, whereas we require M_D to have only non-negative entries. The matrices U and V are obtained analytically in [16] in the limit of large M_R, M_L, and μ such that

$$|M_R\mu|\gg M_W^2\cos^2\beta, \quad |M_L\mu|\gg M_W^2\cos^2\beta, \quad (22)$$

and also for $\sin^2\beta\leftrightarrow\cos^2\beta$. Here we shall only use numerical expressions for V_{ij} and U_{ij}.

With these expressions the chargino contribution to the dipole moment of the electron is

$$d_{e-ch}^E=\frac{\alpha_{EM}}{4\pi\sin^2\theta_W}\sum_{k=1}^{2}\sum_{i=1}^{4}\text{Im}(\gamma_{eik})\frac{m_{\chi_i^+}}{m_{\tilde{\nu}_k}^2}J\left(\frac{m_{\chi_i^+}^2}{m_{\tilde{\nu}_k}^2}\right) \quad (23)$$

where the electron contribution is

$$\gamma_{eik}=(U_{i2}^*D_{\nu2k}+\kappa_eU_{i3}^*D_{\nu1k})(V_{i1}^*D_{\nu1k}-\kappa_\nu V_{i4}^*D_{\nu2k}^*) \quad (24)$$

with

$$\kappa_\nu=\frac{m_\nu}{\sqrt{2}M_W\sin\beta}, \quad \kappa_e=\frac{m_e}{\sqrt{2}M_W\cos\beta}. \quad (25)$$

The sources of *CP* violation enter expression (23) from the imaginary parts of the matrices U_{ik} and V_{ik}, the off-diagonal sneutrino mixings $D_{\nu ik}$ being very small.

B. The neutralino contribution

The neutralino contribution is shown in Fig. 1(b). The neutralino Lagrangian can be written in matrix form as

$$\mathcal{L}_n=-\frac{1}{2}(\Psi^0)^TM^n(\Psi^0)+\text{H.c.} \quad (26)$$

using the basis

$$\Psi^0=(-i\lambda_B^0\cos\theta_W,-i\lambda_L^0,-i\lambda_R^0\sin\theta_W,\tilde{\Phi}_u^0,\tilde{\Phi}_d^0). \quad (27)$$

The neutralino mixing matrix is in general a complex symmetric matrix given by

$$M^N=\begin{pmatrix} M_V+M_R\tan^2\theta_W & 0 & 2(M_R-M_V) & C_1 & -C_2\\ 0 & M_L & 0 & -C_3 & C_2\\ 2(M_R-M_V) & 0 & M_V+\dfrac{M_R}{\tan^2\theta_W} & \dfrac{M_Z\sin\theta_W\cos\beta}{\tan^2\theta_W} & C_2\\ C_1 & -C_3 & \dfrac{M_Z\cos\theta_W\sin\beta}{\tan^2\theta_W} & -\mu & 0\\ -C_2 & C_2 & C_2 & 0 & -\mu\end{pmatrix} \quad (28)$$

where

$$C_1 = M_Z \sin\theta_W \cos\beta \tag{29}$$

$$C_2 = M_Z \cos\theta_W \sin\beta \tag{30}$$

$$C_3 = M_Z \cos\theta_W \cos\beta. \tag{31}$$

Note that this matrix differs from the one given in [16] because of the choice of basis states. Again we assume here that the Higgsino $\tilde{\Delta}^0_{L,R}$ is heavy and decouples. (Allowing for the mixing with the other neutralinos has no measurable effect on the value of the EDM.) Defining the mass eigenstates to be

$$\chi^0 = N_{il}\Psi_l, \quad i,l = 1,\ldots,5 \tag{32}$$

where N_{il} is a unitary matrix which diagonalizes the neutralino mass matrix:

$$M_D^N = N^* M^N B^{-1} \tag{33}$$

with M_D^N the diagonal neutralino mass matrix. With these expressions, the neutralino contribution to the dipole moment of the electron is

$$d^E_{e-n} = -\frac{\alpha_{EM}}{4\pi\sin^2\theta_W}\sum_{k=1}^{2}\sum_{i=1}^{5}\mathrm{Im}(\eta_{eik})\frac{m_{\chi_i^0}}{m^2_{\tilde{e}_k}} I\left(\frac{m^2_{\chi_i^0}}{m^2_{\tilde{e}_k}}\right) \tag{34}$$

where the electron contribution is

$$\begin{aligned}\eta_{eik} = \Bigg\{ & -\sqrt{2}\tan\theta_W N_{1i}\left[\tan^2\theta_W\left(-\frac{1}{2}-2\sin^2\theta_W\right)+(1-\tan^2\theta_W)^{1/2}\left(-\frac{3}{2}-2\sin^2\theta_W\right)\right]D^*_{e1k}+\frac{1}{\sqrt{2}}N_{2i}D^*_{e1k}-\sqrt{2}\tan\theta_W N_{3i}\\ & \times\left[\frac{1}{2}+2\sin^2\theta_W+(1-\tan^2\theta_W)^{1/2}\left(\frac{1}{2}+2\sin^2\theta_W\right)\right]D^*_{e1k}+\kappa_e N_{5i}D^*_{e2k}\Bigg\}\Bigg\{\sqrt{2}\tan\theta_W N_{1i}\Bigg[\left(\frac{1}{2}+2\sin^2\theta_W\right)(1\\ & -\tan^2\theta_W)^{1/2}+\tan^2\theta_W\left[-\frac{1}{2}(\cot^2\theta_W-1)-2\sin^2\theta_W\right]\Bigg]D_{e2k}+\sqrt{2}\tan\theta_W N_{3i}\Bigg[\left(\frac{1}{2}+\sin^2\theta_W\right)(1-\tan^2\theta_W)^{1/2}\\ & -\frac{1}{2}(\cot^2\theta_W-1)-2\sin^2\theta_W\Bigg]D_{e2k}-\kappa_e N_{4i}D_{e1k}\Bigg\}.\end{aligned} \tag{35}$$

The imaginary parts in expression (34) come from the selectron mixing D_{eik} and the mixing matrix elements for the neutralinos N_{ki}.

C. The doubly-charged Higgsino contribution

The supersymmetric left-right model has four doubly-charged Higgsinos [17], two of which have quantum numbers $B-L=-2$, therefore they can interact with leptons only, and two of which have $B-L=2$ and do not interact with the matter multiplets. In a fully left-right symmetric model, both Δ_L and Δ_R exist, and they contribute to lepton-flavor violating decays, such as $Z\to l\bar{l}$ [18], $\mu\to e\gamma$ as well as the anomalous magnetic moment of the muon [19]. In general the Higgsino contribution to the EDM is smaller than the corresponding gaugino one. We will however include this particular contribution here. One reason to do so is that this contribution is new to LR SUSY. The other reason why the doubly charged Higgsino contribution may be important is that its mass is not constrained by the experiment, and it is expected to be small, possibly smaller than that of charginos. The same is true for its Yukawa coupling. In this case, the doubly charged Higgsino contribution can become important. The doubly charged Higgsino contributions to the EDM of the electron at one loop are generated by the interactions in Fig. 1(c). Their contribution is given by

$$d^E_{e-dch} = -\frac{e\,\alpha_{EM}}{4\pi\sin^2\theta_W}\sum_{k=1}^{2}\kappa^2_{\Delta e}\frac{m_{\tilde{\Delta}}}{m^2_{\tilde{e}}}\left[I\left(\frac{m^2_{\tilde{\Delta}^{--}}}{m^2_{\tilde{e}_k}}\right)-2J\left(\frac{m^2_{\tilde{\Delta}^{--}}}{m^2_{\tilde{e}_k}}\right)\right]\mathrm{Im}(\omega_{e\Delta k}) \tag{36}$$

where

$$\omega_{e\Delta k} = (D_{e11}+D_{e12})(D^*_{e21}+D^*_{e22}) \tag{37}$$

and

$$\kappa_{\Delta e} = \frac{h_{LR}\sin\theta_W}{g m_{\tilde{\Delta}}}. \tag{38}$$

The only sources for *CP* violation in this case come from the selectron mixing.

V. NUMERICAL ANALYSIS AND RESULTS

We have given, in the preceding sections, the complete expressions for the contributions to the electron and neutron EDM in the framework of the unrestricted supersymmetric left-right model in terms of the following parameters: M_L, M_R, m_0, A, μ, $\tan\beta$, h_{LR}, $m_{\tilde{\Delta}}$, and the angles θ and α. We make the standard assumption of ignoring all the inter-generational mixing of leptons and sleptons. We take all universal trilinear couplings A to be equal. In general, as in the standard model, one can remove all phases except for two by a redefinition of the fields. We choose the two to be the phase of the μ term, $\sin\theta$ and the phase of the trilinear scalar coupling $\sin\alpha$. There is an extra physical phase in the non-supersymmetric sector coming from the W_L-W_R mixing, and an extra CKM phase in the lepton sector, but we could safely neglect these phases since the induced effects are known to be very small.

Experimentally, the EDMs of the electron and the neutron are some of the most restrictive parameters in the Particle Data Group Booklet, the present experimental upper limits being $d_e \leqslant 4.3\times 10^{-27} e$ cm and $d_n \leqslant 1.1\times 10^{-25} e$ cm [20].

We give first the simplified forms of the supersymmetric contributions to the EDM of the electron using the parameters and approximations outlined above. These are

$$d^E_{e-ch}=\frac{e\,\alpha_{EM}}{4\pi\sin^2\theta_W}\left\{\frac{m_e\sin\mu}{\sqrt{2}M_W\cos\beta}\,\frac{m_{\chi_1}}{m^2_\nu}J\left(\frac{m^2_{\chi_1}}{m^2_{\tilde{\nu}}}\right)U_{31}V_{11}+\frac{\mu}{M_L\tan\beta}\,\frac{m_{\chi_2}}{m^2_\nu}J\left(\frac{m^2_{\chi_2}}{m^2_\nu}\right)U_{43}V_{41}+\mathcal{O}(m_\nu)\right\} \tag{39}$$

$$\begin{aligned}d^E_{e-n}=&\frac{e\,\alpha_{EM}}{4\pi\sin^2\theta_W}\sum_{i=1}^{5}\frac{m_{\chi^0_i}}{m^2_{\tilde{e}}}\left\{\frac{m_e\sin\mu}{\sqrt{2}M_W\cos\beta}[N_{5i}(0.27N_{1i}-0.75N_{3i})-N_{4i}(1.5N_{1i}+0.71N_{2i}-1.4N_{3i})]\right.\\&\times\left[I\left(\frac{m^2_{\chi_i}}{m^2_{\tilde{e}}}\right)+\cos^2\theta_e\frac{(|A|m_0-|\mu|\tan\beta)m_e}{m_{\tilde{e}}}K\left(\frac{m^2_{\chi_i}}{m^2_{\tilde{e}}}\right)\right]+\left[\frac{m_e(m_0|A|\sin\alpha+|\mu|\sin\mu\tan\beta)}{m^2_{\tilde{e}}}\right.\\&\times\left(0.4N_{1i}N_{1i}+0.2N_{1i}N_{2i}-1.7N_{1i}N_{3i}-0.5N_{2i}N_{3i}-1.1N_{3i}N_{3i}+\left(\frac{m_e\sin\mu}{\sqrt{2}M_W\cos\beta}\right)^2N_{5i}N_{4i}\right)\\&+\frac{m_e\sin\mu}{\sqrt{2}M_W\cos\beta}N_{4i}(1.5N_{1i}+0.71N_{2i}-1.4N_{3i})\,\frac{m_e(|A|m_0-|\mu|\tan\beta)}{m^2_{\tilde{e}}}\\&\left.\left.+\left(\frac{m_e\sin\mu}{\sqrt{2}M_W\cos\beta}\right)^2\sin2\mu N_{4i}N_{5i}\right]K\left(\frac{m^2_{\chi_i}}{m^2_{\tilde{e}}}\right)\right\}\end{aligned} \tag{40}$$

$$d^E_{e-dch}=\frac{e\,\alpha_{EM}}{4\pi\sin^2\theta_W}\left\{\kappa^2_{\Delta e}\frac{m_e}{m^4_{\tilde{e}}}(m_0|A|\sin\alpha+|\mu|\sin\mu\tan\beta)m_{\tilde{\Delta}}G\left(\frac{m^2_{\tilde{\Delta}}}{m^2_{\tilde{e}}}\right)\right\} \tag{41}$$

where we have inserted, for simplicity, $\sin^2\theta_W=0.2315$ in the expression for the neutralino contribution. The functions K and G are obtained from expanding the expressions for the EDM loop functions around an average slepton mass $m^2_{\tilde{e}}=(m^2_{\tilde{e}_1}+m^2_{\tilde{e}_2})/2$ and are given by

$$K(r)=\frac{1}{2(1-r)^3}\left(1+5r+\frac{2r(2+r)\ln r}{1-r}\right) \tag{42}$$

and

$$G(r)=\frac{1}{2(1-r)^3}\left(9+r-4r^2+\frac{2(4-r^2)\ln r}{1-r}\right). \tag{43}$$

The coupling $\kappa_{\Delta e}$ depends on the value of the left-right coupling h_{LR} [as in Eq. (38)]. The corresponding Higgs triplets Δ_L,Δ_R are not involved in lepton mass generation. Therefore, their couplings to the leptons are not suppressed. We will assume, for simplicity, that the couplings of the left- and right-handed Higgs bosons are equal; we will also assume this coupling to be large, since these Higgs bosons are responsible for breaking left-right symmetry and giving mass to the right-handed boson W_R. One could reasonably set $h_{LR}\simeq\mathcal{O}(1)$, which is consistent with the bounds obtained from lepton-flavor violation in left-right models [21]. We will also take the mass of the $\tilde{\Delta}^{--}$ to be $m_{\tilde{\Delta}}=100$ GeV, which is roughly as expected [11].

The matrices U_{ij}, V_{ij} and N_{ij} are obtained numerically for given values of the parameters. The terms of $\mathcal{O}(m_\nu)$ can be nonzero, but are constrained to be small, much smaller than the rest of the terms in the chargino contributions, so taking the left-handed neutrino to have a (small) mass has no

TABLE II. The relative size of supersymmetric contributions to the EDM of the electron: $m_0 = 1000$ GeV; $|A| = 1$; $m_{\tilde{\Delta}} = 100$ GeV; $\theta = \alpha = \pi/10$ rad.

Case	M_L	μ	$d_{e-c}(10^{-27} e$ cm)	$d_{e-n}(10^{-27} e$ cm)	$d_{e-dch}(10^{-27} e$ cm)
(i)	200	200	22.1	5.1×10^{-1}	1.8×10^{-3}
(ii)	200	1000	7.04	1.4	4.5×10^{-3}
(iii)	1000	200	1.13	1.8×10^{-1}	3.4×10^{-3}
(iv)	1000	1000	4.9	1.6	4.5×10^{-3}

effect on the present calculation. In the sneutrino mass matrix, the effect is that of decoupling of the sneutrinos, their mixing being proportional to m_ν. The difference between the EDM in LR SUSY and the MSSM lies in the values for the mixing matrices U_{ij}, V_{ij} and N_{ij}, the contributions of the extra neutralinos, and the additional contribution of the doubly charged Higgsinos. Before we begin the complete analysis, we present below the numerical values of the contributions to the electron EDM from the chargino, neutralino and doubly charged Higgsino, for same-sign *CP*-violating angles α and θ (Table II), and opposite sign *CP*-violating angles α and θ (Table III). We take the right-handed scale to be $M_R = 10$ TeV in both scenarios.

One could see that in both scenarios the chargino contribution is dominant, as it is in the MSSM; this is due to their couplings. The neutralino contribution tends to be larger than the one in the MSSM, especially for large μ, where the chargino and the neutralino contributions are of the same order. The doubly charged Higgsino contribution is much smaller than the neutralino contribution, even for light doubly charged Higgsinos. (Even taking the mass of the doubly charged Higgsino to be the improbable value of $m_{\tilde{\Delta}} = 50$ GeV would increase the doubly charged Higgsino contribution by only an order of magnitude, still too small to influence the EDM in any significant way.) This is due to the fact that the chargino and neutralino contributions are dominated by the gaugino-gaugino, or the gaugino-Higgsino couplings, whereas the doubly charged Higgsino coupling is proportional to $\kappa^2_{\Delta e}$, a Higgsino-Higgsino coupling, so this contribution is much smaller even for large Yukawa couplings $h_{LR} \simeq \mathcal{O}(1)$. In the case in which the *CP*-violating angles have opposite signs, the neutralino contribution can change signs, and for small angles θ and moderate angles α, could be of the same order of magnitude as the chargino contribution. In this case we could obtain the same situation as in $N=1$ supergravity [4], where one would have a suppression of the EDM without large superpartner masses or small *CP*-violating angles.

As a general feature of the model, the electron EDM can exceed the present experimental bounds unless (a) the superpartner (selectron, sneutrino) masses are large, or (b) the *CP*-violating angles are small (barring unforeseen cancellations as discussed above). The exact restriction on the masses (or angles, depending on one's point of view) depends on the parameters M_L, μ, and $\tan\beta$. Figure 2 shows the variation of the electron EDM with the universal scalar mass m_0 for the low $\tan\beta$ scenario ($\tan\beta=3$) for the four (M_L,μ) scenarios presented in Table II. In what follows we take $M_R = 10$ TeV. As can be seen, the electron EDM falls off with increasing m_0, behavior understood since $K(r)/m^2_{\tilde{e}}$, $I(r)/m^2_{\tilde{e}}$ and $J(r)/m^2_{\tilde{\nu}}$ decrease with increasing m_0. If anything, the LR SUSY model is even more restrictive on the values of the scalar masses or angles. In particular, for left-gaugino masses of the order of the electro-weak scale, the constraint on the electron EDM would require that either the scalar masses must be $m_0 \geqslant 4.5$ TeV, or the angles must be of $\mathcal{O} \leqslant \pi/100$. The situation is somewhat alleviated for larger M_L,μ (in effect, for a heavier chargino-neutralino spectrum). The result is the same for same sign *CP*-violating angles [Fig. 2(a)], or opposite-sign angle [Fig. 2(b)], except that in the latter partial cancellations can occur.

Figure 3 shows the same situation in the high $\tan\beta$ scenario ($\tan\beta=50$). The restrictions there must be extremely severe; if not, for most reasonable values of the supersymmetric parameters the electron EDM is two orders of magnitude larger than the experimental result. The $\tan\beta$ dependence is very dramatic. Indeed, as seen in Fig. 4, for both same-sign and opposite-sign values of the *CP*-violating angles, the electron EDM scales like $\tan\beta$ for large values of $\tan\beta$. This can be explained through the inverse $\cos\beta$ dependence in chargino and neutralino contributions, which scale like $\tan\beta$ for large values of $\tan\beta$. This dependence is so strong that it offsets the chargino-neutralino spectrum dependence on $\tan\beta$. The only way out would be to expect a heavy scalar fermion spectrum for large $\tan\beta$, since the *CP*-

TABLE III. The relative size of supersymmetric contributions to the EDM of the electron: $m_0 = 1000$ GeV; $|A| = 1$; $m_{\tilde{\Delta}} = 100$ GeV; $\theta = -\alpha = \pi/10$ rad.

Case	M_L	μ	$d_{e-c}(10^{-27} e$ cm)	$d_{e-n}(10^{-27} e$ cm)	$d_{e-dch}(10^{-27} e$ cm)
(i)	200	200	22.1	1.3×10^{-1}	4.3×10^{-4}
(ii)	200	1000	7.04	-6.8×10^{-1}	-2.25×10^{-3}
(iii)	1000	200	1.13	1.6×10^{-1}	4.3×10^{-4}
(iv)	1000	1000	4.9	-7.9×10^{-1}	-2.25×10^{-3}

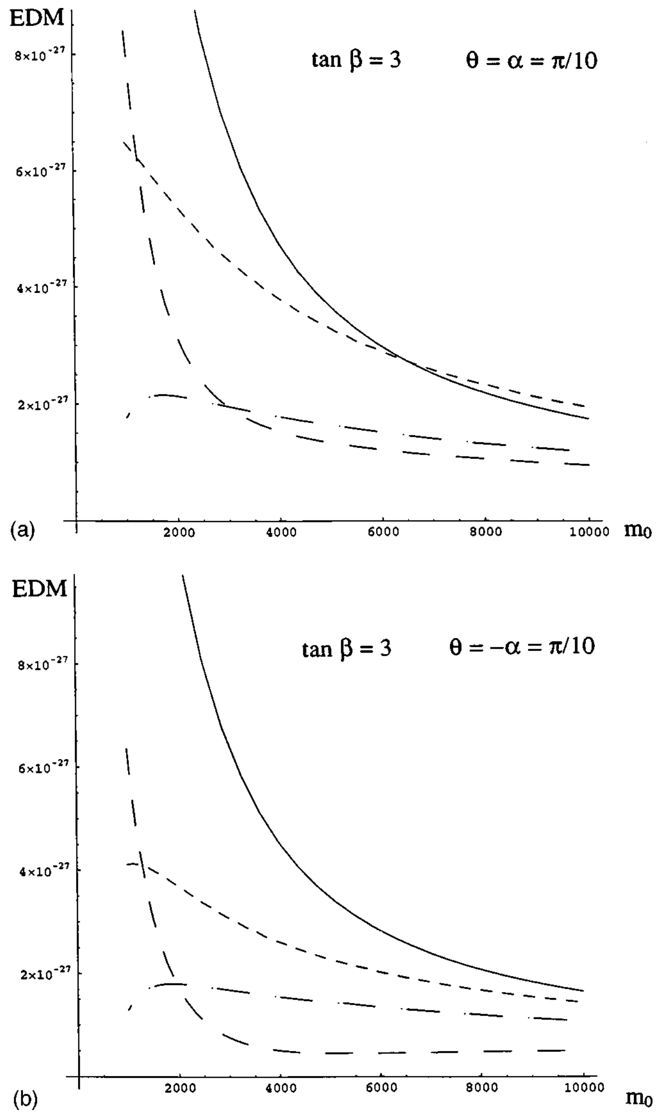

FIG. 2. Plot of the variation of the electron EDM with the universal scalar mass parameter m_0 for low $\tan\beta$, $\tan\beta=3$ for the four scenarios shown in Table II. The curves are marked: (solid curve) $M_L=200$ GeV, $\mu=200$ GeV; (large-dash curve) $M_L=200$ GeV, $\mu=1000$ GeV; (dot-dashed curve) $M_L=1000$ GeV, $\mu=200$ GeV; (small-dash curve) $M_L=1000$ GeV, $\mu=1000$ GeV. We take $|A|=1$, $M_R=10$ TeV. (a) shows the case in which the *CP*-violating angles have same sign $\theta=\alpha=\pi/10$ rad; (b) shows the case in which the *CP*-violating angles have opposite signs $\theta=-\alpha=\pi/10$ rad.

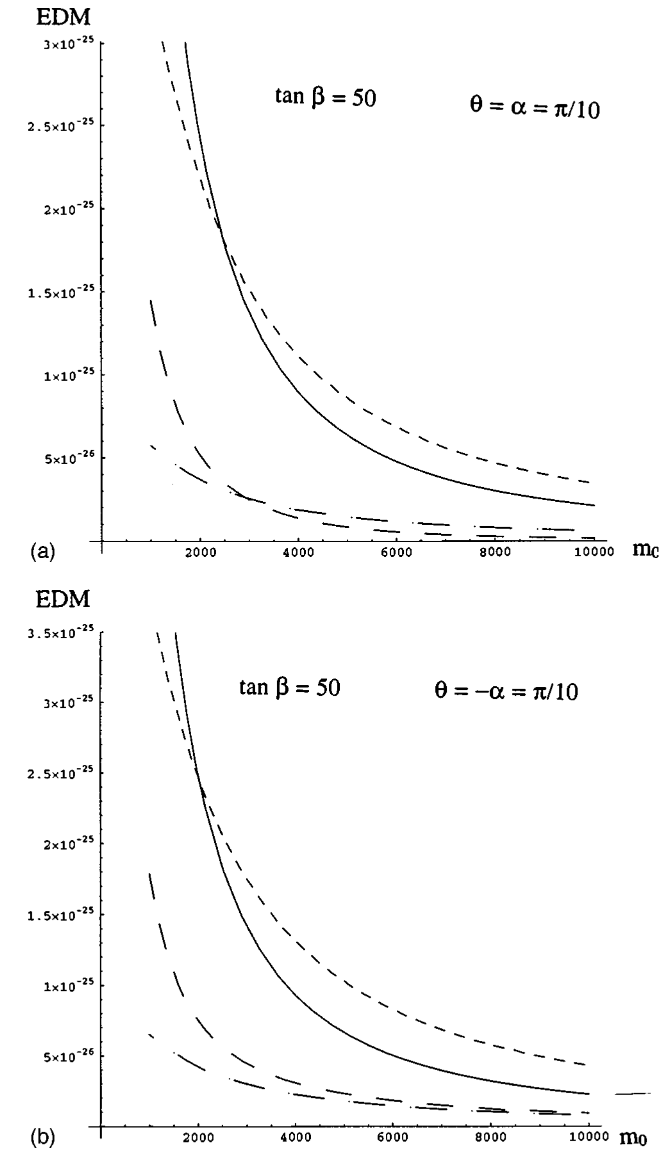

FIG. 3. Same plot as in Fig. 2, but for high $\tan\beta$ case, $\tan\beta=50$. The curves are marked: (solid curve) $M_L=200$ GeV, $\mu=200$ GeV; (large-dash curve) $M_L=200$ GeV, $\mu=1000$ GeV; (dot-dashed curve) $M_L=1000$ GeV, $\mu=200$ GeV; (small-dash curve) $M_L=1000$ GeV, $\mu=1000$ GeV.

violating angles α and θ are independent of $\tan\beta$, but this scenario would require even more fine tuning than the usual required to keep the EDM within experimental bounds.

Figure 5 investigates the dependence of the electron EDM of the *CP*-violating angles. In Fig. 5(a) we show the dependence on the angle $\sin\theta$ for a heavy scalar lepton spectrum $m_0=5$ TeV. Even for this value of m_0, the $\sin\theta$ angle is constrained to be in the $\{-0.4,0.4\}$ region. The dependence on α is less dramatic, due to the fact that the chargino contribution is practically independent of the angle α [Fig. 5(b)], but for moderate values of α and small values of θ some cancellations between the chargino and neutralino contributions can occur.

Figures 6 and 7 show the dependence of the electron EDM on the left-gaugino mass parameter M_L and the Higgsino parameter μ. These parameters affect the mass spectrum of the charginos and neutralinos. The dependence is shown for two values of the scalar mass, $m_0=1.5$ TeV and $m_0=5$ TeV. The increase in the masses of the charginos and neutralinos is offest by the decrease in the functions $I(r)$, $J(r)$ and $K(r)$, so the variation is not as pronounced as perhaps expected.

Figure 8 shows the dependence of the electron EDM on the trilinear coupling A. The EDM depends only on the electron A_e (and is therefore independent of any assumption about the universality of the trilinear couplings). Increasing the value of $|A|$ beyond 1, the neutralino contribution can become larger than the chargino, offseting the chargino

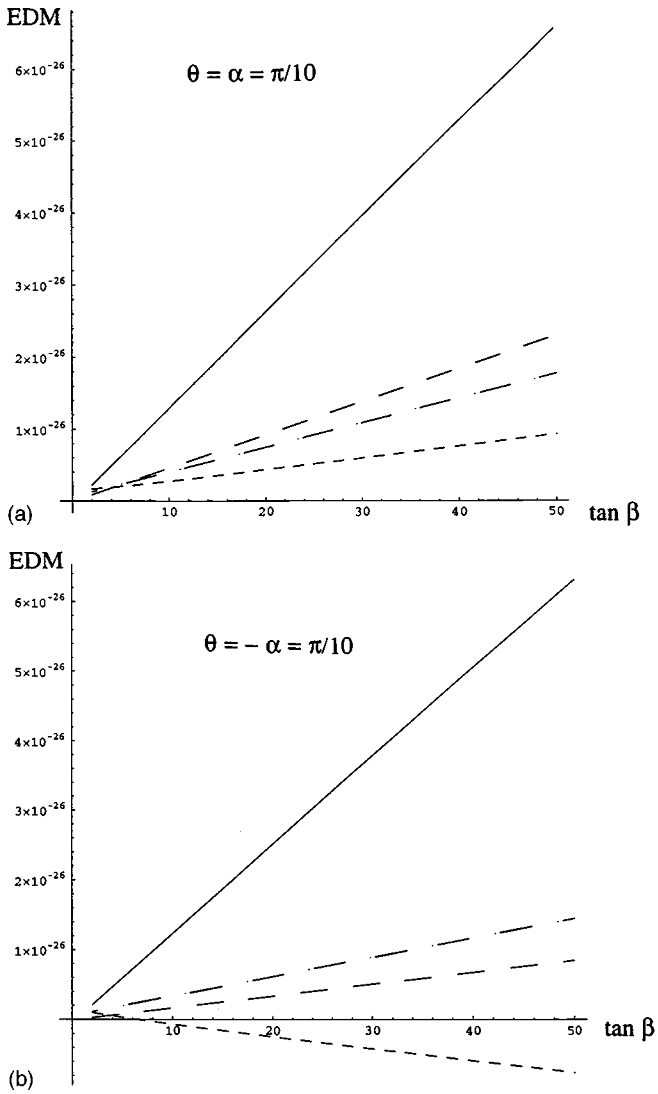

FIG. 4. Plot of the variation of the electron EDM with tan β. We consider the same four cases as in Table II for $M_R = 10$ TeV, $m_0 = 5$ TeV. (a) shows the case in which the *CP*-violating angles have same sign $\theta = \alpha = \pi/10$ rad; (b) shows the case in which the *CP*-violating angles have opposite signs $\theta = -\alpha = \pi/10$ rad.

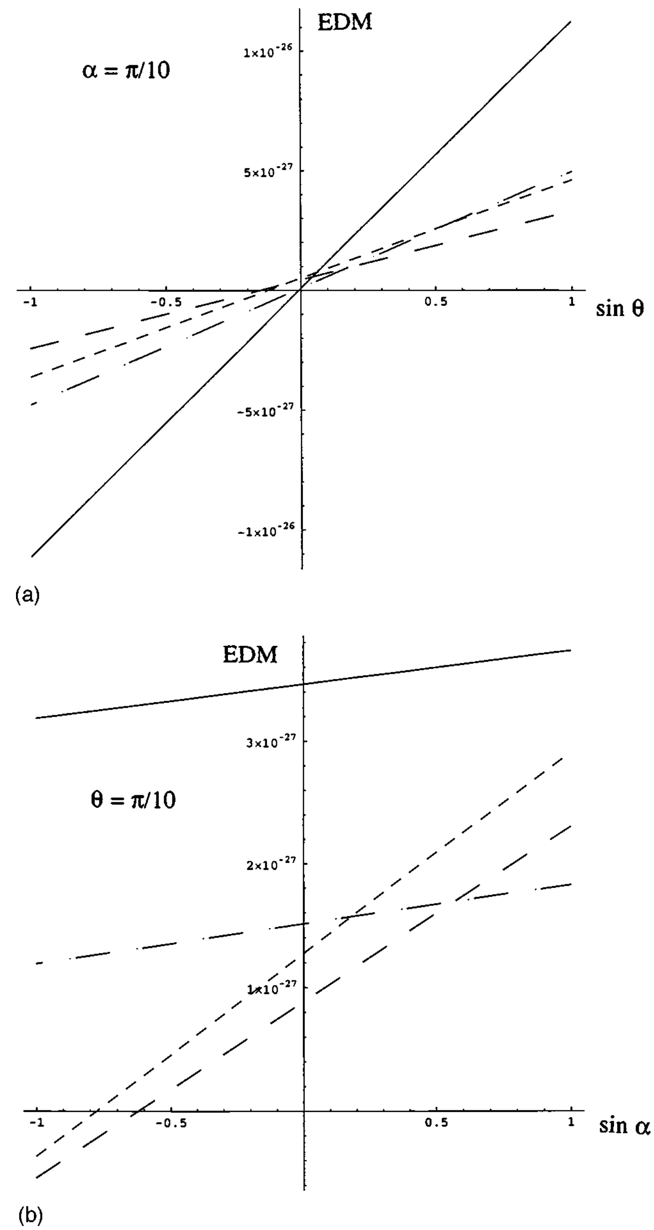

FIG. 5. Plot of the variation of the electron EDM with the *CP*-violating angles θ (a) and α (b), for the four scenarios in Table II. The other parameters are fixed at $M_R = 10$ TeV, $m_0 = 5$ TeV, $|A| = 1$. The curves are marked: (solid curve) $M_L = 200$ GeV, $\mu = 200$ GeV; (large-dash curve) $M_L = 200$ GeV, $\mu = 1000$ GeV; (dot-dashed curve) $M_L = 1000$ GeV, $\mu = 200$ GeV; (small-dash curve) $M_L = 1000$ GeV, $\mu = 1000$ GeV.

dominance seen for the spectrum otherwise. In particular, for a combination of larger values for A and opposite-sign *CP*-violating angles α and θ, the neutralino contribution can be of the same magnitude and opposite in sign to the chargino one, resulting in a region of parameter space where the electron EDM is very small.

Finally, a word on the variation of the electron EDM with the mass of the right-handed scale M_R. At first sight, the results are practically independent of the values of M_R. The explanation comes from the dominance of the chargino contribution in most cases. The chargino contribution is not much affected by variations in the values of the right-hand scale. The physical chargino state corresponding to the right-handed W-ino $\tilde{\lambda}_R$ does not contribute much to the EDM. For large values of M_R, the EDM contains a remnant of the LR SUSY in the double Higgsino contribution, whereas the neutralino spectrum resembles more the MSSM [16]. However, it is expected that the right-hand scale will have an important effect on the *CP*-violating phases in the model, as explained below.

As in the case of the MSSM, one might question the naturalness of the parameter choice, as well as the naturalness of making either the scalar spectrum heavy, or the *CP*-violating angles very small; and in both cases the restrictions of the LR SUSY are more stringent than those of the MSSM. In the MSSM, the parameters associated with soft supersymmetry breaking are the least understood parameters. In a commonly-used version of the model the parameters of the

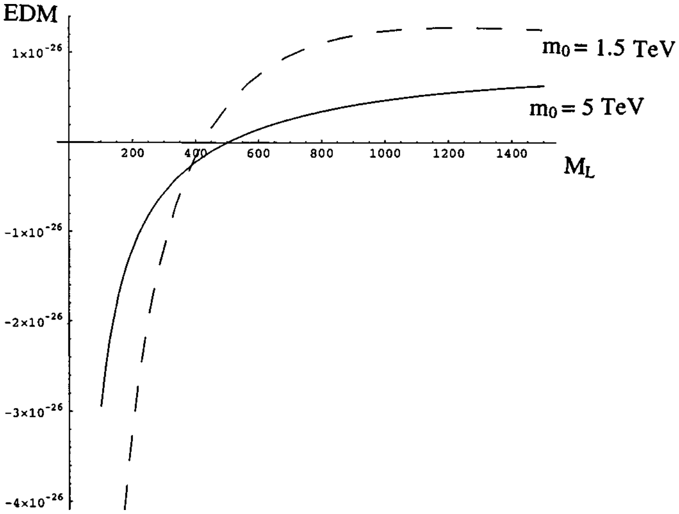

FIG. 6. Plot of the variation of the electron EDM with the left-gaugino mass parameter M_L. The other parameters are fixed as $M_R = 10$ TeV, $\mu = 1$ TeV, $|A| = 1$, $\theta = \alpha = \pi/10$ rad. The solid curve is for $m_0 = 5$ TeV; the dashed for $m_0 = 1.5$ TeV.

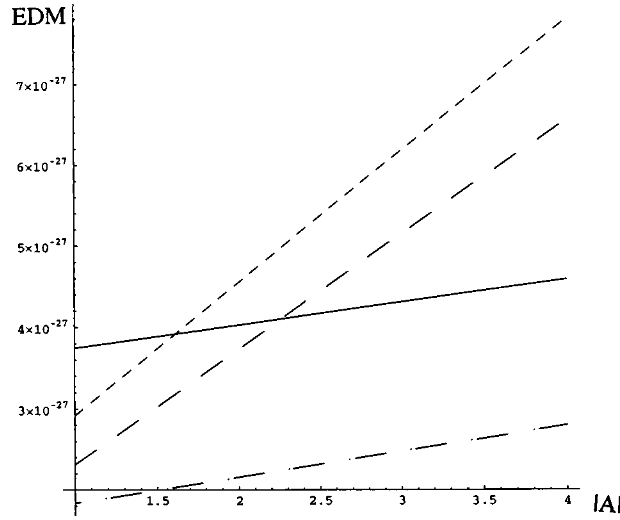

FIG. 8. Plot of the variation of the electron EDM with the trilinear coupling $|A|$. The other parameters are set to $M_R = 10$ TeV, $M_L = 1$ TeV, $\mu = 1$ TeV, $\theta = \alpha = \pi/10$ rad. The curves are marked: (solid curve) $M_L = 200$ GeV, $\mu = 200$ GeV; (large-dash curve) $M_L = 200$ GeV, $\mu = 1000$ GeV; (dot-dashed curve) $M_L = 1000$ GeV, $\mu = 200$ GeV; (small-dash curve) $M_L = 1000$ GeV, $\mu = 1000$ GeV.

model are real, except for the Cabibbo-Kobayashi-Maskawa-(CKM) type phase at the unification or the Planck scale. The phases are then radiatively induced from the renormalization group equations involving the CKM matrix. What could one expect in the LR model? In the model of Ref. [8] the authors show that, above the M_R scale, there are no one-loop contributions to the neutron EDM. If one imposes left-right symmetry, the trilinear coupling A and the gaugino masses are real above M_R. However, below M_R, *CP*-violating phases are generated through symmetry breaking. *If* one expects the angles generated this way to be naturally small, of $\mathcal{O} \approx 2 \times 10^{-5}$, one might argue that, in this case, it is reasonable to assume that the LR SUSY model is a means to select the small scenario of *CP*-violating angles naturally, and the model would predict an electron EDM safely within the experimental bounds. To predict exactly what the angles will be would depend on the right-handed scale M_R. On one hand, the requirement of small *CP*-violating angles seem to favor a right-hand scale of the order of the electroweak, or supersymmetry scale; on the other hand, constraints coming from the absence of flavor changing neutral current (FCNC) effects and R-parity conservation tend to push this scale to the unification scale [11]. For this reason, it is important to have a general calculation of the electron EDM in the left-right supersymmetric model.

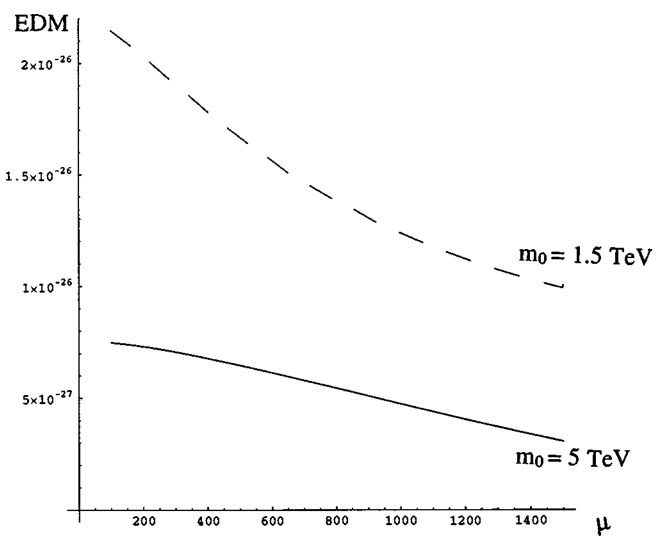

FIG. 7. Plot of the variation of the electron EDM with the Higgsino mass parameter μ. The other parameters are fixed as $M_R = 10$ TeV, $M_L = 1$ TeV, $|A| = 1$, $\theta = \alpha = \pi/10$ rad. The solid curve is for $m_0 = 5$ TeV; the dashed for $m_0 = 1.5$ TeV.

VI. CONCLUSION

We analyzed the electron EDM in the LR SUSY model including all one-loop contribution from charginos, neutralinos and doubly charged Higgsinos. We found that the chargino contribution dominates in almost all cases. This will provide an independent restriction on that contribution, apart from possible cancellations with the gluino contribution that exists when one considers the neutron EDM. We found that in all cases the new contribution from the doubly charged Higgsino is much smaller than the other two. We found that cancellations between the neutralino and chargino contributions can occur, in the case in which the *CP*-violating angles are opposite, or $|A|$ is large, but the cancellation is most of the time only partial and significant only for certain combinations of the parameters. Without any restriction, we find the neutralino contribution to be larger than the one in the MSSM (perhaps making the cancellation more likely). In the absence of any special considerations, the LR SUSY model imposes even stricter limits on the masses of the scalar partners of the leptons, or on the smallness of the *CP*-violating angles. A recent proposal to justify the smallness of the *CP*-violating angles [8] is an attractive solution to the electron EDM problem.

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