Electric dipole moments from flavored *CP* violation in supersymmetry

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The so-called supersymmetric flavor and *CP* problems are deeply related to the origin of flavor and hence to the origin of the standard model Yukawa couplings themselves. We show that realistic $SU(3)$ flavor symmetries with spontaneous *CP* violation reproducing correctly the standard model Yukawa matrices can simultaneously solve both problems without *ad hoc* modifications of the supersymmetric model. We analyze the leptonic electric dipole moments and lepton flavor violation processes in these models. We show that the electron electric dipole moment and the decay $\mu \to e\gamma$ are naturally within reach of the proposed experiments if the sfermion masses are measurable at the LHC.

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I. INTRODUCTION

The so-called supersymmetry (SUSY) flavor and *CP* problems are usually taken as the main naturalness problems SUSY has to face up to [1,2]. The formulation of the supersymmetric flavor problem is well known: since we have no information regarding the structure of SUSY soft-breaking terms, we could, in principle, expect that entries in a soft-mass matrix are all of the same order. In particular, this can happen in the basis where the Yukawa couplings are diagonal. In such a situation, flavor-changing neutral current (FCNC) and flavor-dependent *CP* violation observables would receive too large contributions from loops involving SUSY particles to satisfy the stringent phenomenological bounds on these processes [1,2]. We can formulate the SUSY *CP* problem in a similar way. If *CP* is not a symmetry of the model, we naturally expect all complex parameters in the model to have $O(1)$ phases. In this case, the phases in flavor-independent terms typically generate too large contributions to the so-far unobserved electric dipole moment (EDM) of the electron and neutron [3,4].

The basis of both problems clearly lies in our total ignorance about the origin of the observed flavor and *CP* violation in our theory. However, notice that these problems are not restricted to supersymmetry. Even the standard model (SM) shares the flavor problem with SUSY in exactly the same terms. If we had not measured the quark and lepton masses and mixings, we would naturally expect all the elements in the Yukawa matrices to be $O(1)$. Yet, if one gives such a structure to the Yukawa matrices, the predicted fermion masses and mixings would never agree with the observed ones. Therefore, we have to conclude that there is a much stronger flavor problem in the SM than in the minimal supersymmetric standard model (MSSM). The real flavor problem is simply our inability to understand the complicated structures in the quark and lepton Yukawa couplings, and likewise for the soft-breaking flavor structures in the MSSM. On the other hand, there seems to be no direct analog of the SUSY *CP* problem in the SM. In fact, the phases in the SM Lagrangian are already $O(1)$ without being in conflict with the experimental measurements. However, this apparent "fact" is also misleading. Notice that, due to the particle content of the SM, the only complex parameters in the Lagrangian are the Yukawa couplings themselves and we have measured them to be small. Once more, if we had not known the fermion masses and mixings beforehand and wrote arbitrary complex Yukawa parameters, we would also have a severe SM *CP* problem. Since the SUSY *CP* problem is basically due to the flavor-independent phases in the MSSM, both facts can suggest the idea that the flavor and the *CP* problems are indeed related, and solving the flavor problem while restricting the *CP* phases to the flavor sector would also solve the *CP* problem.

A particularly attractive solution to these problems (both in the SM and in SUSY) is found in models based on flavor symmetries. In these models, the flavor structure of the Yukawa matrices is only generated after the breaking of a flavor symmetry [5–14], and the flavor structure of the SUSY soft-breaking terms would also originate from the same mechanism [15–19]. Thus, finding a solution to the SM flavor problem will generally solve, at the same time, the so-called SUSY flavor problem to a sufficient degree, although probably still allowing naturally suppressed contributions that might bring more information about the flavor sector. Regarding the SUSY *CP* problem, if we want to restrict all *CP* phases in SUSY to the flavor sector, this can be achieved by postulating an exact *CP* symmetry spontaneously broken in the flavor sector. This would remove all flavor-independent phases, but still produce interesting observable sources of *CP* violation.

In this work we shall analyze how a flavor symmetry can solve the flavor and *CP* problems in a SUSY scenario. We

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shall also distinguish typical signatures of such a symmetry in the lepton sector.[1] In Sec. II we shall analyze a general flavor symmetry model which can solve both problems, based on [18]. In Sec. III we shall identify the most important flavor-dependent contributions to leptonic EDMs, in order to quantify the expected amount of *CP* violation in such models. In Sec. IV we shall study in a similar manner the relevant lepton flavor violation (LFV) processes. Finally, in Sec. V we will show the regions of the supergravity parameter space that are sensitive to future EDM and LFV experiments.

II. FLAVOR SYMMETRIES AND SPONTANEOUS *CP* BREAKING

Following the original ideas of Froggatt and Nielsen [6], flavor symmetries explain the peculiar structure of the SM Yukawa couplings as the result of a spontaneously broken symmetry associated with flavor. The three generations of SM fields are charged under this symmetry such that the SM Yukawa couplings are not allowed in the limit of exact symmetry. One or several scalar vacuum expectation values (vevs) breaking this symmetry must be inserted in a nonrenormalizable operator, suppressed by a heavy mediator mass, to compensate the charges. If the scalar vev is smaller than the mediator scale, this provides a small expansion parameter that can be used to explain the hierarchy of the observed Yukawa couplings.

In the context of a supersymmetric theory, an unbroken flavor symmetry would apply equally to the fermion and scalar sectors. This implies that in the limit of exact symmetry the soft-breaking scalar masses and the trilinear couplings must be invariant under the flavor symmetry. This has different implications in the case of the scalar masses and the trilinear couplings.

The scalar masses are couplings $\phi\phi^\dagger$; thus flavor-diagonal couplings are clearly invariant under any symmetry, i.e. diagonal soft-masses are always allowed by flavor symmetry. Therefore diagonal scalar masses will be of the order of the SUSY soft-breaking scale. However, in general, this does not guarantee that they are family universal. Being universal or not will depend on the considered family symmetry. In the case of an Abelian [6,7,10,12,13,21–23] family symmetry, the symmetry does not relate different generations and therefore diagonal masses can be different. In this case, FCNC and *CP* violation phenomenology set very strong constraints on the differences between these flavor-diagonal masses, and Abelian family symmetries have serious difficulties satisfying these constraints [24]. On the other hand, a non-Abelian family symmetry groups two or three generations in a single multiplet with a common mass, thus solving the FCNC problem. This was one of the main motivations for the construction of the first $SU(2)$ flavor models [11,25], where the first- and second-generation sfermions, facing the strongest constraints, share a common mass. In the case of an $SU(3)$ flavor symmetry [18,26–28], all three generations have the same mass in the unbroken family symmetry limit. For this reason, in the following we will consider non-Abelian family symmetries and, more precisely, $SU(3)$ flavor symmetries.

On the contrary, trilinear couplings are completely equivalent to the Yukawa couplings from the point of view of the symmetry because they involve exactly the same fields (scalar or fermionic components). Thus, they are forbidden by the symmetry [with the possible exception of the (3, 3) component in $SU(2)$ models] and generated only after symmetry breaking as a function of small vevs.

In addition to the renormalizable mass operators in the Lagrangian, we can construct nonrenormalizable operators that are neutral under flavor symmetry inserting an appropriate number of flavon fields. The flavon fields, charged under the symmetry, are responsible for the spontaneous symmetry breaking once they acquire a vev. Then, higher dimensional operators involving two SM fermions and a Higgs field, with several flavon vevs suppressed by a large mediator mass, generate the observed Yukawa couplings. In the same way, these flavon fields will couple to the scalar fields in all possible ways allowed by the symmetry and, after spontaneous symmetry breaking, they will generate a nontrivial flavor structure in the soft-breaking parameters. Given that they are generated by insertions of the same flavon vevs, we can expect the structures in the soft-breaking matrices and the Yukawa couplings to be related. Our starting point in our analysis of the soft-breaking terms must then involve an analysis of the texture in the Yukawas, in order to reproduce first the correct masses and mixings.

To fix the Yukawa couplings, we accept that the smallness of Cabibbo-Kobayashi-Maskawa (CKM) mixing angles is due to the smallness of the off-diagonal elements in the Yukawa matrices with respect to the corresponding diagonal elements, and we make the additional simplifying assumption of choosing the matrices to be symmetric. With these two theoretical assumptions, and using the ratio of masses at a high scale to define the expansion parameters in the up and down sectors as $\bar{\varepsilon} = \sqrt{m_s/m_b}$ and $\varepsilon = \sqrt{m_c/m_t}$, we can fix the Yukawa textures in the quark sector to be

$$Y_d \propto \begin{pmatrix} 0 & b\bar{\varepsilon}^3 & c\bar{\varepsilon}^3 \\ b\bar{\varepsilon}^3 & \bar{\varepsilon}^2 & a\bar{\varepsilon}^2 \\ c\bar{\varepsilon}^3 & a\bar{\varepsilon}^2 & 1 \end{pmatrix}, \quad Y_u \propto \begin{pmatrix} 0 & b'\varepsilon^3 & c'\varepsilon^3 \\ b'\varepsilon^3 & \varepsilon^2 & a'\varepsilon^2 \\ c'\varepsilon^3 & a'\varepsilon^2 & 1 \end{pmatrix}, \tag{1}$$

[1]Observables in the quark sector, such as the neutron EDM, shall be studied in a future work [20].

with $\bar{\varepsilon} \simeq 0.15$, $\varepsilon \simeq 0.05$, $b = 1.5$, $a = 1.3$, $c = 0.4$, and with a', b', c' poorly fixed from experimental data [18,22,29]. Unfortunately, the Yukawa couplings in the leptonic sector cannot be determined from the available phenomenological data. The left-handed neutrino masses and mixings cannot unambiguously fix the neutrino Yukawa couplings in a seesaw mechanism. Therefore only the charged-lepton masses provide useful information on leptonic Yukawas. For simplicity, we choose to work with a grand unified model at high scales, a possibility which is favored by the unification of the bottom and tau Yukawa couplings. In this case, charged-lepton and down-quark (and the neutrino and up-quark) flavor matrices are the same except for the different vev of a Georgi-Jarlskog Higgs field [30] to unify the second- and first-generation masses.

With this Yukawa structure as our starting point, we will generate the flavor structure of the soft-breaking terms in an explicit example based on $SU(3)$ flavor symmetry. Under this symmetry, the three generations of SM fields, including both $SU(2)_L$ doublets and singlets, are triplets **3** and the Higgs fields are singlets. Therefore Yukawa couplings and trilinear terms, $\mathbf{3}\times\mathbf{3}\times\mathbf{1}$, are not allowed by $SU(3)$. In the theory, we have several flavon fields, that we call θ_3, θ_{23} (antitriplets $\bar{\mathbf{3}}$), $\bar{\theta}_3$, and $\bar{\theta}_{23}$ (triplets **3**). The symmetry is broken in two steps: first θ_3 and $\bar{\theta}_3$ get a vev, $\propto (0, 0, 1)$, breaking $SU(3)$ into $SU(2)$. Subsequently, a smaller vev of θ_{23} and $\bar{\theta}_{23}$, $\propto (0, 1, 1)$, breaks the remaining symmetry [26].

To reproduce the Yukawa textures, the large third-generation Yukawa couplings require a θ_3 (and $\bar{\theta}_3$) vev of the order of the mediator scale, M_f (slightly smaller in the up sector and $\tan\beta$ dependent in the down sector as shown in the Appendix), while θ_{23}/M_f (and $\bar{\theta}_{23}/M_f$) have vevs of order ε in the up sector and $\bar{\varepsilon}$ in the down sector, with different mediator scales in both sectors. In this way, third-generation Yukawa couplings are generated by $\theta_3^i\theta_3^j$, while couplings in the 2–3 block of the Yukawa matrix are always given by $\theta_{23}^i\theta_{23}^j(\theta_3\bar{\theta}_3)^2$ (possibly with the Georgi-Jarlskog field Σ to unify quark and leptonic Yukawa couplings, not included here for simplicity; see Refs. [18,26,27]). The couplings in the first row and the first column of the Yukawa matrix are given by $\epsilon^{ikl}\bar{\theta}_{23,k}\bar{\theta}_{3,l}\theta_{23}^j(\theta_{23}\bar{\theta}_3)^n$, where to reproduce the texture in Eq. (1) we must force $n = 1$. Unfortunately, the $SU(3)_{\rm fl}$ symmetry is not enough to reproduce the textures in Eq. (1), and we must impose some additional global symmetries (typically Z_N symmetries) to guarantee the correct power structure and to forbid unwanted terms, like a mixed $\theta_3^i\theta_{23}^j$ term, in the effective superpotential. The basic structure of the Yukawa superpotential (for quarks and leptons) is then given by

$$W_{\rm Y} = H\psi_i\psi_j^c[\theta_3^i\theta_3^j + \theta_{23}^i\theta_{23}^j(\theta_3\bar{\theta}_3) + \epsilon^{ikl}\bar{\theta}_{23,k}\bar{\theta}_{3,l}\theta_{23}^j(\theta_{23}\bar{\theta}_3) + \ldots], \tag{2}$$

where, to simplify the notation, we have normalized the flavon fields to the corresponding mediator mass; i.e., all the flavon fields in this equation should be understood as θ_i/M_f. This structure is quite general for the different $SU(3)$ models we can build, and for additional details we refer to [18,26,27].

In the same way, after $SU(3)$ breaking, the scalar soft masses deviate from exact universality. As explained above, $\phi_i^\dagger\phi_i$ is completely neutral under gauge and global symmetries and gives rise to a common contribution for the family triplet. However, after $SU(3)$ breaking, terms with additional flavon fields give rise to important corrections [18,31–33]. Any invariant combination of flavon fields can also contribute to the sfermion masses. In this case, it is easy to see that the following terms will always contribute to the sfermion mass matrices[3]:

$$(M_{\tilde{f}}^2)^{ij} = m_0^2\bigg(\delta^{ij} + \frac{1}{M_f^2}[\theta_{3,i}^\dagger\theta_{3,j} + \bar{\theta}_3^i\bar{\theta}_3^{j\dagger} + \theta_{23,i}^\dagger\theta_{23,j} + \bar{\theta}_{23}^i\bar{\theta}_{23}^{j\dagger}] + \frac{1}{M_f^4}(\epsilon^{ikl}\bar{\theta}_{3,k}\bar{\theta}_{23,l})^\dagger(\epsilon^{jmn}\bar{\theta}_{3,m}\bar{\theta}_{23,n}) + \ldots\bigg), \tag{3}$$

where f represents the $SU(2)$ quark and lepton doublets or the up (neutrino) and down (charged-lepton) singlets. Notice that we have three different mediator masses, $M_f = M_L, M_u, M_d$, because the flavor symmetry must commute with the SM symmetry and therefore the vectorlike mediator fields must have the SM quantum numbers of the usual particles.[4] With these terms, we can see some similarities between the flavor structures of the Yukawa and soft-mass matrices. In particular, the off-diagonal (2, 3) elements in the Yukawa matrix are given by $\bar{\theta}_{23}^2\bar{\theta}_{23}^3/M_f^2 \simeq \epsilon_f^2$ [the term $(\theta_3\bar{\theta}_3)$ factorizes from all the contributions in the Yukawa matrix], while in the soft masses it is $\bar{\theta}_{23}^{i\dagger}\bar{\theta}_{23}^j/M_f^2$, also of order ϵ_f^2. Still, it is possible to build other invariant combinations with different flavon fields that cannot be present in the superpotential. This is due to the fact that the superpotential must be holomorphic, i.e. cannot include daggered fields, while the soft masses, coming from the Kähler

[2]Notice that we add the scalar product $(\theta_3\bar{\theta}_3)$ to this operator with respect to Refs. [18,26,27] to be able to generalize to different $\tan\beta = (\langle\theta_3^u\rangle/\langle\theta_3^d\rangle)^2$ values.

[3]Discrete non-Abelian subgroups of $SU(3)$ [34–41] would have a similar leading order structure in the soft-mass matrices.

[4]Nevertheless, in all our numerical calculations below we take only two different mediator masses, M_d and $M_u = M_L$, for simplicity

potential, are only a real combination of fields. In some cases, the global symmetries can allow a contribution like $\theta_{23,i}\bar{\theta}_3^j$ + H.c. to the soft masses, even though the combination $\theta_{23,i}\theta_{3,j}$ is forbidden by global symmetries in the superpotential. This structure would strongly affect the phenomenology of third-generation physics [20]. In any case, although possible, this situation is very rare and usually the structure of the soft terms follows the Yukawa structure. It is important to emphasize at this point that these deviations from universality in the soft-mass matrices proportional to flavor symmetry breaking always come through corrections in the Kähler potential. Therefore, these effects will be important only in gravity-mediation SUSY models where the low-energy soft-mass matrices are mainly generated through the Kähler potential [42]. In other mediation mechanisms, such as gauge mediation [43] or anomaly mediation [44–46], where these Kähler contributions to the soft masses are negligible, flavor effects in the soft-mass matrices will be basically absent.

In the case of the trilinear couplings we have to emphasize that from the point of view of flavor symmetry these couplings are completely equivalent to the corresponding Yukawa coupling. This means that they necessarily involve the same combination of flavon vevs, although order one coefficients are generically different because they require at least an additional coupling to a field mediating SUSY breaking (in general, coupled in different ways in the various contributions). Therefore, from our point of view, we expect that the trilinear couplings have the same structure as the Yukawa matrices in the flavor basis. However, in general, they are *not* proportional to the Yukawas, because of different $O(1)$ coefficients in the different elements. Thus, we can expect that going to the super-CKM (SCKM) basis does not diagonalize the trilinear matrices. In fact, the trilinear matrices maintain the same structure as in the flavor basis, and only the $O(1)$ coefficients are modified.

We should now take into account, especially in a gravity-mediated scenario, the canonical normalization of the kinetic terms (Kähler potential). However, as shown in Refs. [32,47] these canonical transformations do not modify the general structure of the different flavor matrices and change only the unknown $O(1)$ coefficients. As these coefficients cannot be fixed by flavor symmetry at this level, the previous discussion on the different flavor matrices remains valid in the canonical basis.

We have now fixed the flavor structure that we can expect in the soft-breaking matrices. Now, we are ready to consider the problem of CP violation in these flavor matrices. In fact, in the SM, CP violation is deeply related to flavor. The only possible complex parameters in the SM are the Yukawa couplings themselves, and all observed CP violation is consistent with a single observable phase in the CKM matrix. Therefore, we need complex flavon vevs to reproduce the observed CP in the CKM mixing matrix. On the other hand, the supersymmetric CP problem concerns the fact that, in the presence of flavor-independent phases in the μ term and trilinear couplings, supersymmetry gives rise to contributions to the EDMs at one-loop order with no suppression associated with flavor [48–55]. These one-loop contributions, for SUSY masses below several TeV, can easily exceed the present experimental bounds shown in Table I by 1 or 2 orders of magnitude. This fact forces most models to demand unnatural requirements, such as vanishing phases, very large mediator masses, or engineered cancellations between different contributions to the process [64].

TABLE I. Current constraints to EDMs (left) and reach of future experiments (right).

EDM	Current bound (e cm)	Future bound (e cm)
$\vert d_e\vert$	$\leq 1.4\times 10^{-27}$ [56]	$\sim 10^{-32}$ [57]
$\vert d_\mu\vert$	$\leq 7.1\times 10^{-19}$ [58]	$\sim 10^{-23}$ [59]
$\vert d_\tau\vert$	$\leq 2.5\times 10^{-17}$ [60]	$\sim 10^{-20}$ [61]
$\vert d_n\vert$	$\leq 2.9\times 10^{-26}$ [62]	$\sim 10^{-28}$ [63]

In this aspect, flavor symmetries with spontaneous breaking of CP provide an interesting solution. CP is an exact symmetry of the theory before the breaking of flavor symmetry. Thus, above the scale of flavor breaking, all terms in the Kähler potential, which give rise to the soft masses and the μ term (by the Giudice-Masiero mechanism [65]), are real. Even after the breaking of flavor and CP symmetries, μ receives flavor-suppressed complex corrections only at the two-loop level [66]. Finally, trilinear terms are only generated after symmetry breaking with the same phase structure as the Yukawa couplings, and it can be proven that diagonal elements in A_{ij} are real at leading order in the SCKM basis [18]. In this way, the supersymmetric CP problem is naturally solved. Nevertheless, flavor-dependent phases do exist and can contribute to the fermion EDMs as we will see below.

In the following sections, even though we use exact expressions to calculate all our observables, we shall use the mass insertion (MI) approximation [67,68] to quantify and explain all important contributions (as in [69]). In terms of the slepton mass matrix terms, these MIs are defined as

$$(\delta^e_{\rm LR})_{ij} = \frac{v_d}{\sqrt{2}M^2_{\tilde{e}}}[A^*_{e,ij} - \delta_{ij}Y_i\mu\tan\beta] = (\delta^e_{\rm RL})^*_{ji},$$
$$(\delta^e_{\rm LL})_{ij} = \frac{(m^2_{\tilde{L}})_{ij}}{M^2_{\tilde{e}}}, \qquad (\delta^e_{\rm RR})_{ij} = \frac{(m^2_{\tilde{e}})_{ji}}{M^2_{\tilde{e}}}, \tag{4}$$

where $M^2_{\tilde{e}}$ is the average slepton mass. From the soft slepton mass matrices of Eq. (A1), we can read the approximate MIs in the $\mu - e$, $\tau - e$, and $\tau - \mu$ sectors, respectively:

$$(\delta^e_{\mathrm{LL}})_{12} \approx \frac{\epsilon^2 \bar{\epsilon}}{3}, \qquad (\delta^e_{\mathrm{RR}})_{12} \approx \frac{\bar{\epsilon}^3}{3}, \qquad (\delta^e_{\mathrm{LR}})_{12} \approx A_0 \frac{m_\tau}{M^2_{\tilde{e}}} \bar{\epsilon}^3, \tag{5a}$$

$$(\delta^e_{\mathrm{LL}})_{13} \approx \bar{\epsilon}^3 y^\nu_{33}, \qquad (\delta^e_{\mathrm{RR}})_{13} \approx \frac{\bar{\epsilon}^3}{3}, \qquad (\delta^e_{\mathrm{LR}})_{13} \approx A_0 \frac{m_\tau}{M^2_{\tilde{e}}} \bar{\epsilon}^3, \tag{5b}$$

$$(\delta^e_{\mathrm{LL}})_{23} \approx \bar{\epsilon}^2 y^\nu_{33}, \qquad (\delta^e_{\mathrm{RR}})_{23} \approx \bar{\epsilon}^2, \qquad (\delta^e_{\mathrm{LR}})_{23} \approx A_0 \frac{m_\tau}{M^2_{\tilde{e}}} 3\bar{\epsilon}^2, \tag{5c}$$

where we have not included the renormalization group (RG) running from the unification scale down to low energy. Such running effects can be sizable in the LL and LR sectors, due to the presence of heavy right-handed neutrinos with large Yukawa couplings [70–72]. Moreover, in the present case there are new contributions to the running given by the nonuniversality of the soft-mass and trilinear matrices. These effects can be important even in the RR sector. As an example, the element $(m^2_{\tilde{e}})_{32}$ of the right-handed slepton mass matrices gets the following RG correction in the SCKM basis:

$$(m^2_{\tilde{e}})_{32}(M_{\mathrm{SUSY}}) \simeq \left[(m^2_{\tilde{e}})_{32}(M_U)\left(1 - \frac{1}{16\pi^2}(2y^2_\tau)\right) - 4\frac{A^e_{33}A^{e\dagger}_{32}}{16\pi^2}\right] \log\left(\frac{M_U}{M_{\mathrm{SUSY}}}\right) \tag{6}$$

where M_U is the unification scale, M_{SUSY} the average SUSY mass, and the matrix elements are evaluated at M_U using the SCKM expressions given in the Appendix. In the SCKM basis we have $A^e_{13}/A_0 \sim (Y_\nu)_{13} \sim \mathcal{O}(\bar{\epsilon}^3)$, $A^e_{23}/A_0 \sim (Y_\nu)_{23} \sim \mathcal{O}(\bar{\epsilon}^2)$, $A^3_{33}/A_0 \sim m_\tau/(v\cos\beta)$ and $(Y_\nu)_{33} \sim \mathcal{O}(1)$, and $y_e \sim \mathcal{O}(\bar{\epsilon}^4)$, $y_\mu \sim \mathcal{O}(\bar{\epsilon}^2)$, $y_\tau \sim \mathcal{O}(1)$. Therefore, we expect sizable corrections from the running only in the $\tau - \mu$ and $\tau - e$ sectors, i.e. where the third generation is involved.

III. EDMS AND FLAVOR PHASES

Fermion EDMs, d_ψ, are induced through effective dimension-six operators (with an implicit Higgs insertion providing the chirality charge) of the form

$$\mathcal{L} = -\frac{d_\psi}{2}[\bar{\psi}\sigma^{\mu\nu}\gamma^5\psi]F_{\mu\nu}, \tag{7}$$

being related to the imaginary part of a chirality-changing, flavor-diagonal loop process. In the SM these processes are greatly suppressed: the prediction for the electron EDM is lower than $O(10^{-40})e$ cm [73]. This makes EDMs very convenient observables in which to look for new physics related to *CP* violation.

As we have seen above, if *CP* is spontaneously broken in the flavor sector, the usual flavor-independent phases coming from μ and A_f are approximately zero and we expect EDMs to be under control.[5] Nevertheless, flavor-dependent phases in the soft-mass matrices and trilinears can also give large contributions [74,75]. In order to be sure that the SUSY *CP* problem is solved, it will be necessary to quantify the expected order of magnitude of the EDMs produced by $O(1)$ phases on these terms. If the current constraints are respected, we can then contrast these predictions with the expected sensitivity of future EDM experiments, shown also in Table I.

In order to identify the dominant terms in d_e, one needs to know the size and phases of the different mass insertions. In fact, observable phases will correspond to rephasing-invariant combinations of mass insertions and Yukawa elements [76]. Even in the general $SU(3)$ framework presented in the previous section, these combinations depend strongly on the particular model one takes into account. Thus, as a first step, it is of interest to study the magnitude of each potential contribution to d_e assuming a generic phase of order 1 for the whole rephasing-invariant combination, using the expected size of the off-diagonal elements in flavor symmetry. We will then proceed with a second analysis, considering the explicit model by Ross, Velasco-Sevilla, and Vives (RVV) [18] presented in the Appendix.

One-loop MSSM contributions to charged-lepton EDMs d_l ($l = e,\ \mu,\ \tau$) involve diagrams with neutralinos and charginos [77,78]. However, the chargino contribution only involves a flavor-diagonal left-handed sneutrino propagator, and, due to Hermiticity, the sensitivity to the phases within the sneutrino mass matrix is lost. With a vanishing phase for μ, we can neglect the chargino contribution to d_l, and concentrate on neutralinos completely. Neutralino contributions to d_e are

$$d_e^{\chi^0} = \left(\frac{e}{16\pi^2}\right)\Im m(A^L_{eij}A^R_{eij})\frac{1}{m_{\chi^0_j}}F\left(\frac{m^2_{\chi^0_j}}{m^2_{\tilde{l}_i}}\right), \tag{8}$$

$$A^L_{eij} = Y_e\mathcal{N}^*_{3j}\mathcal{R}_{\tilde{e}_R i} - \frac{g'}{\sqrt{2}}\mathcal{N}^*_{1j}\mathcal{R}_{\tilde{e}_L i} - \frac{g}{\sqrt{2}}\mathcal{N}^*_{2j}\mathcal{R}_{\tilde{e}_L i}, \tag{9}$$

[5]Notice that, since we are assuming gaugino mass unification, we can always take the unified gaugino mass as real, corresponding to the usual convention in the constrained MSSM.

$$A^R_{eij} = Y_e \mathcal{N}^*_{3j} \mathcal{R}^*_{\tilde{e}_L i} + \sqrt{2} g' \mathcal{N}^*_{1j} \mathcal{R}^*_{\tilde{e}_R i} \quad (10)$$

with $\mathcal{R}_{\tilde{e}i}$ and $\mathcal{N}_{kj}$ being elements of the charged slepton and neutralino mixing matrices [77], respectively, and the loop function $F(x) = \frac{x}{2(x-1)^3}(x^2 - 1 - 2x \log x)$.

In terms of mass insertions, the most relevant contributions from Eq. (8) can be written as

$$\frac{d_e}{e} = \frac{\alpha M_1}{8\pi \cos^2\theta_W M^2_{\tilde{e}}} \Im m[(\delta^e_{\rm LL})_{1i}(\delta^e_{\rm LR})_{i1} f_1 + (\delta^e_{\rm LR})_{1i}(\delta^e_{\rm RR})_{i1} f_2 + (\delta^e_{\rm LL})_{1i}(\delta^e_{\rm LR})_{ij}(\delta^e_{\rm RR})_{j1} f_3] \quad (11)$$

where $M^2_{\tilde{e}}$ is the average slepton mass and the loop functions f_i can be derived from [69].

Let us explain briefly why these are the relevant contributions, and identify the dominant ones. All phases in this $SU(3)$ flavor model are contained within the sfermion mass matrices and thus we shall need at least one mass insertion on the slepton line. Since all flavor-conserving insertions $(\delta^e_{\rm LR})_{ii}$ are real to leading order, we will need to combine at least two flavor-changing mass insertions, $(\delta^e_{AB})_{ij}$.

Regardless of the number of mass insertions, we shall always have two situations: one in which the neutralino line couples to the fermion through a gaugino and a Higgsino, and one that couples through two b-inos (although interactions with two Higgsinos also contribute, they are suppressed by at least an additional Yukawa coupling, so we shall not discuss them). In the gaugino-Higgsino case, the slepton line will need to maintain its handedness and again, due to the Hermiticity of the full slepton mass matrix, loses all dependency on the flavon phases. Thus, d_e is due entirely to diagrams with pure b-inos as the vertices, where a LR transition is required.

With two mass insertions, the only contribution with physical phases comes from a combination of insertions like $(\delta^e_{\rm LR})_{1i}(\delta^e_{\rm RR})_{i1}$ or $(\delta^e_{\rm LL})_{1i}(\delta^e_{\rm LR})_{i1}$. With three mass insertions, we consider only contributions with a single LR insertion. It is well known that each $\delta^e_{\rm LR}$ insertion is suppressed by a cumulative factor $m_\tau/M_{\tilde{e}}$. Therefore the dominant contribution comes from $(\delta^e_{\rm LL})_{1i}(\delta^e_{\rm LR})_{ij}(\delta^e_{\rm RR})_{j1}$. Obviously, the largest contribution is the one that involves a central $(\delta^e_{\rm LR})_{33}$, due to the $m_\tau \tan\beta$ enhancement. Thus, the most important contribution to d_e with three mass insertions is the pure b-ino $(\delta^e_{\rm LL})_{13}(\delta^e_{\rm LR})_{33}(\delta^e_{\rm RR})_{31}$.

Regarding whether the two- or three-insertion contribution dominates, it shall depend on the magnitude of $\tan\beta$ and the size of the off-diagonal terms in the trilinears. Evidently, for a large enough $\tan\beta$, the three-insertion contribution shall dominate no matter what the size of $(\delta^e_{\rm LR})_{ij}$ is, but for low values the situation is not so clear, especially if A_0 is large. Using Eq. (A1) in the Appendix, we can quantify the magnitude of each insertion as

$$(\delta^e_{\rm LR})_{1i}(\delta^e_{\rm RR})_{i1} \approx A_0 \bar{\epsilon}^6 \frac{m_\tau}{M^2_{\tilde{e}}}, \quad (12)$$

$$(\delta^e_{\rm LL})_{13}(\delta^e_{\rm LR})_{33}(\delta^e_{\rm RR})_{31} \approx \bar{\epsilon}^6 y^\nu_{33} \frac{m_\tau}{M_{\tilde{e}}} \tan\beta, \quad (13)$$

and therefore for $A_0 \simeq M_{\tilde{e}}$ the triple mass insertion is dominant except for small values of $y^\nu_{33} \tan\beta$.

It is also important to take into account that Eq. (12) has the same structure for the 12 and 13 elements. This means that, with a maximum phase on each element, we can double the two-insertion contribution. In any case, as these terms are all proportional to A_0, in order to do an appropriate study one needs to take the case where $A_0 = 0$ as a standard, and then understand further deviations when $A_0 \neq 0$.

Let us turn now to d_μ. For $A_0 = 0$, the structure of the dominant terms for d_μ shall be quite similar to the one for d_e. We shall have the main contribution coming from the triple mass insertion $(\delta^e_{\rm LL})_{23}(\delta^e_{\rm LR})_{33}(\delta^e_{\rm RR})_{32}$, which is enhanced by $m_\tau \tan\beta$. However, due to the flavor structure of our model, we should expect a suppression of order $\bar{\epsilon}^4$, instead of $\bar{\epsilon}^6$. Thus, d_μ is about 2 orders of magnitude larger than d_e. This is similar to the usual mass scaling relation, which predicts d_μ to be larger by m_μ/m_e, also 2 orders of magnitude. When $A_0 \neq 0$, the double insertion can be relevant for low $\tan\beta$, similar to the case of d_e analyzed in Eqs. (12) and (13). However, contrary to d_e, where both $(\delta^e_{\rm LR})_{12}(\delta^e_{\rm RR})_{21}$ and $(\delta^e_{\rm LR})_{13}(\delta^e_{\rm RR})_{31}$ are of the same order ($\bar{\epsilon}^6$), in this case the dominant contribution only comes from $(\delta^e_{\rm LR})_{23}(\delta^e_{\rm RR})_{32}$, which is again of order $\bar{\epsilon}^4$.

The situation for d_τ is critically different when $A_0 = 0$. In this case the $m_\tau \tan\beta$ enhancement is lost, and the main triple MI contribution is due to $(\delta^e_{\rm LL})_{32}(\delta^e_{\rm LR})_{22}(\delta^e_{\rm RR})_{23}$. This would be smaller than d_μ by a factor m_μ/m_τ, almost 2 orders of magnitude, and thus d_τ clearly violates the naive scaling relation. In contrast, when $A_0 \neq 0$, the main contributions from the double MIs with a flavor-changing $\delta_{\rm LR}$ insertion are identical in magnitude to those for d_μ, and so we expect d_τ to be of a size comparable to d_μ. In this case, we should take into account the possible presence of subdominant phases in the SCKM basis in flavor-diagonal trilinear couplings [31]. In any case, this breaks again the mass scaling relation, even though not so drastically as in the previous situation. The observation of such bizarre behavior would be a very clear signal favoring these types of flavor models.

IV. LEPTON FLAVOR-VIOLATING DECAYS

As discussed in the previous sections, supersymmetric flavor models are characterized by nonuniversal scalar masses at the scale where the SUSY breaking terms appear. Moreover, the trilinear A_f matrices are, in general, not aligned with the corresponding Yukawa matrices. This leads to the appearance of potentially large mixing among flavors. In particular, in the lepton sector, the same mass insertions which induce EDMs are sources of lepton flavor violation, again via neutralino or chargino loop diagrams

[79]. As a consequence, we expect a correlation between EDM and LFV processes and the allowed parameter space to be strongly constrained by the experimental limits on LFV decays such as $\mathrm{BR}(l_i \to l_j\gamma)$.

The branching ratio of $l_i \to l_j\gamma$ can be written as

$$\frac{\mathrm{BR}(l_i \to l_j\gamma)}{\mathrm{BR}(l_i \to l_j\nu_i\bar{\nu}_j)} = \frac{48\pi^3\alpha}{G_F^2}(|A_L^{ij}|^2 + |A_R^{ij}|^2), \tag{14}$$

with the SUSY contribution to each amplitude given by the sum of two terms $A_{L,R} = A^n_{L,R} + A^c_{L,R}$, where $A^n_{L,R}$ and $A^c_{L,R}$ denote the contributions from the neutralino and chargino loops, respectively. In the mass insertion approximation, and taking only the dominant terms, we can write the amplitudes as follows:

$$A_L^{ij} = \frac{\alpha_2}{4\pi}\frac{(\delta^e_{\mathrm{LL}})_{ij}}{m^2_{\tilde{l}}}\bigg[\frac{\mu M_2\tan\beta}{(M_2^2-\mu^2)}F_{2\mathrm{LL}}(a_2, b) + \tan^2\theta_W\frac{\mu M_1\tan\beta}{(M_1^2-\mu^2)}F_{1\mathrm{LL}}(a_1, b)\bigg] + \frac{\alpha_1}{4\pi}\frac{(\delta^e_{\mathrm{RL}})_{ij}}{m^2_{\tilde{l}}}\bigg(\frac{M_1}{m_{l_i}}\bigg)F_{1\mathrm{LR}}(a_1), \tag{15}$$

$$A_R^{ij} = \frac{\alpha_1}{4\pi}\bigg(\frac{(\delta^e_{\mathrm{RR}})_{ij}}{m^2_{\tilde{l}}}\frac{\mu M_1\tan\beta}{(M_1^2-\mu^2)}F_{1\mathrm{RR}}(a_1, b) + \frac{(\delta^e_{\mathrm{LR}})_{ij}}{m^2_{\tilde{l}}}\bigg(\frac{M_1}{m_{l_i}}\bigg)F_{1\mathrm{LR}}(a_1)\bigg), \tag{16}$$

where θ_W is the weak mixing angle, $a_{1,2} = M^2_{1,2}/\tilde{m}^2$, $b = \mu^2/m^2_{\tilde{l}}$, and the loop functions $F_{1\mathrm{LL}}$, $F_{2\mathrm{LL}}$, $F_{1\mathrm{RR}}$, and $F_{1\mathrm{LR}}$ can be obtained from the expressions in Refs. [80,81].

Here we can see that $(\delta^e_{\mathrm{LL}})_{ij}$ and $(\delta^e_{\mathrm{RR}})_{ij}$ contributions are $\tan\beta$ enhanced. In contrast to minimal flavor violation models with RH neutrinos, in our model the largest contribution to $\mu \to e\gamma$ comes from the RR sector, being $(\delta^e_{\mathrm{RR}})_{12} \simeq \bar{\varepsilon}^3$ while $(\delta^e_{\mathrm{LL}})_{12} \simeq \varepsilon^2\bar{\varepsilon}$ (see Sec. II). The $\tau \to \mu\gamma$ and $\tau \to e\gamma$ decays shall have similar LL and RR contributions.

On the other hand, the only term proportional to $(\delta^e_{\mathrm{LR}})_{ij}$ arises from pure $\tilde{B}$ exchange and it is completely independent of $\tan\beta$. However, in these flavor models LR mass insertions can still be important. In the case of the $\mu \to e\gamma$ decay and taking into account the necessary chirality change in the amplitude, we have to compare $(\delta^e_{\mathrm{RR}})_{12}m_\mu\tan\beta$ with $(\delta^e_{\mathrm{LR}})_{12}M_1$, as we can see from Eq. (16). Using the expression for the mass insertions of Eq. (5c), we see that in these models

$$\frac{(\delta^e_{\mathrm{RR}})_{12}m_{l_i}\tan\beta}{(\delta^e_{\mathrm{LR}})_{12}M_1} \simeq \frac{(\bar{\varepsilon}^3/3)m_\mu\tan\beta}{\bar{\varepsilon}^3A_0(m_\tau/M^2_{\tilde{e}})M_1} = \frac{m_\mu\tan\beta}{3m_\tau}\frac{M^2_{\tilde{e}}}{A_0M_1}. \tag{17}$$

Therefore, we can see that if $M^2_{\tilde{e}}/(A_0M_1) \sim O(1)$, the LR mass insertion will dominate the $\mu \to e\gamma$ decay up to $\tan\beta \sim 30$. In fact, these contributions can easily bring $\mathrm{BR}(\mu \to e\gamma)$ to the level of the present experimental reach, and therefore, we expect that the $A_0 \neq 0$ scenarios will be very strongly constrained by the present and future limits on $\mathrm{BR}(\mu \to e\gamma)$. This is the main consequence of the misalignment between A_e and Y_e. Let us notice that here the LR contribution, even if not enhanced by $\tan\beta$, becomes dominant due to an enhancement by a factor of order m_τ/m_μ with respect to the other contributions to the amplitude. This is clearly peculiar of $\mu \to e\gamma$, and it is not verified in the case of $\tau \to \mu\gamma$. For $\tau \to \mu\gamma$, even with $A_0 \neq 0$ the LR contribution is subdominant with respect to the other ones, mainly proportional to $(\delta^e_{\mathrm{LL}})_{32}$, which are $\tan\beta$ enhanced. Therefore, we expect the $\mathrm{BR}(\tau \to \mu\gamma)$ for $A_0 \neq 0$ to be approximately equal to the case $A_0 = 0$.

It is also interesting to compare the different LFV channels. In the case $A_0 = 0$, the dominant LFV source should be δ^e_{LL}, which contributes to both chargino and neutralino diagrams. Thus a rough estimation for the relative sizes of the branching ratios can be

$$\frac{\mathrm{BR}(\tau \to e\gamma)}{\mathrm{BR}(\mu \to e\gamma)} \approx \bigg(\frac{m_\tau}{m_\mu}\bigg)^5\frac{\Gamma_\mu}{\Gamma_\tau}\frac{(\delta^e_{\mathrm{LL}})^2_{13}}{(\delta^e_{\mathrm{LL}})^2_{12}} \approx \mathcal{O}(1), \tag{18}$$

$$\frac{\mathrm{BR}(\tau \to \mu\gamma)}{\mathrm{BR}(\mu \to e\gamma)} \approx \bigg(\frac{m_\tau}{m_\mu}\bigg)^5\frac{\Gamma_\mu}{\Gamma_\tau}\frac{(\delta^e_{\mathrm{LL}})^2_{23}}{(\delta^e_{\mathrm{LL}})^2_{12}} \approx \mathcal{O}(10^3) \tag{19}$$

where Γ_μ (Γ_τ) is the μ (τ) full width. Given the present limit $\mathrm{BR}(\tau \to e\gamma) < 1.1 \times 10^{-7}$ [82], we can see from Eq. (18) that $\tau \to e\gamma$ is not able to constrain the parameter space of the model better than $\mu \to e\gamma$ whose experimental bound is $\mathrm{BR}(\mu \to e\gamma) < 1.2 \times 10^{-11}$ [83] (which will be improved by 2 orders of magnitude by MEG [84]). On the other hand, we expect from Eq. (19) that the present constraints given by $\mu \to e\gamma$ and $\tau \to \mu\gamma$ are comparable, once the combined *BABAR* + Belle limit $\mathrm{BR}(\tau \to \mu\gamma) < 1.6 \times 10^{-8}$ is considered [85,86].

It is important also to clarify the dependence of our results on the chosen value of Y_ν. Notice that, in our $SU(3)$ model, the values of the neutrino Yukawa couplings are fixed by the symmetry following [18]. However, different values of Y_ν could be possible in other examples with respect to observed values of neutrino masses and mixings. In any case, as can be seen in Eqs. (A4), only the (1, 3) and (2, 3) elements of $m^2_{\tilde{L}}$ depend on the value of y^ν_{33} at one loop. Therefore only the predictions on $\tau \to l_i\gamma$ can be affected by a change on y^ν_{33}, and even in this case the contributions from $m^2_{\tilde{e}}$, independent of y^ν_{33}, will be of a similar size.

V. NUMERICAL RESULTS

In the following, we shall use the expressions for d_e and LFV processes to put bounds on the supergravity parameter space through a scan in m_0 and $M_{1/2}$ for fixed values of $\tan\beta$ and $a_0 = A_0/m_0$. Since all of our predictions will depend on arbitrary $O(1)$ coefficients, it is not possible to provide a precise numerical result. Thus, the following

TABLE II. Applied constraints coming from EDMs and LFV (left) and the neutral meson sector (right).

Observable	Bound	Observable	Bound
d_e	$<1.4\times10^{-27}e$ cm	Δm_K	$<3.48\times10^{-15}$ GeV
$\mathrm{BR}(\mu\to e\gamma)$	$<1.2\times10^{-11}$	Δm_D	$<4.61\times10^{-14}$ GeV
$\mathrm{BR}(\tau\to\mu\gamma)$	$<4.5\times10^{-8}$	Δm_B	$<3.34\times10^{-13}$ GeV
$\mathrm{BR}(\tau\to e\gamma)$	$<1.1\times10^{-7}$	Δm_{B_s}	$<1.17\times10^{-11}$ GeV
$\mathrm{BR}(b\to s\gamma)^{\mathrm{SUSY}}$	$<0.88\times10^{-4}$	ϵ_K	$<2.239\times10^{-3}$

discussion intends to point out the expected order of magnitude for these observables, and factors of 2 or even 4 can appear.

Our numerical analysis presented below is done defining the Yukawa, trilinear, and soft-mass matrices at $M_{\mathrm{flav}}=2\times10^{16}$ GeV, as explained in Sec. II and explicitly shown in the Appendix. Then we evolve the different flavor matrices to the electroweak scale using one-loop renormalization group equations (RGEs) [87]. $O(1)$ coefficients in the Yukawa matrices are determined by requiring a good fit on the fermion masses and quark mixings at M_Z [29]. Other $O(1)$ terms in the soft matrices are taken as random, varying between 0.5 and 2. The different values of $\tan\beta$ are fixed by the ratio of the vevs (a_3^u/a_3^d), which in our model are global factors in the Yukawa (and trilinear) matrices.

After running the resulting structures down to the M_Z scale, we diagonalize the Yukawas in order to obtain the left and right mixing matrices and rotate the soft matrices into the SCKM basis. Notice that this SCKM rotation generates the off-diagonal elements in the first row and the first column in all soft-mass matrices, and in the case of left-handed sleptons it also generates the dominant contribution to $(M_{\tilde{L}}^2)_{23}$. At this scale, we apply the CERN LEP bounds on the lightest sparticle and Higgs masses, and we require a neutral lightest supersymmetric particle (LSP). In the plots, regions that fail to satisfy any of these requirements are shown in dark brown (black). Then, we also apply the constraints set by the current bounds from d_e, LFV processes, and FCNC measurements on the hadronic sector: Δm_K, ϵ_K, Δm_D, Δm_B, Δm_{B_s}, and $b\to s\gamma$ [88], as shown in Table II. Nonetheless, only $\mu\to e\gamma$, $\tau\to\mu\gamma$, and ϵ_K shall exclude regions in the parameter space above the LEP and LSP bounds; we show these regions, stretching above the dark brown regions at low m_0, in green (gray).

Initially, we present the results in a generic $SU(3)$ model with phases $O(1)$ in the flavor soft terms, as shown in the Appendix, Eq. (A1). Then, we take the RVV model as an explicit example with a well-defined phase structure, as shown in Eqs. (A4).[6]

[6]In this model, even though each term in the soft matrices receives corrections from the RGEs, the leading terms from these corrections have an identical phase structure, so we can expect the phases not to change much by the running.

A. Generic $SU(3)$ model

In this model, we assume generic $O(1)$ phases in the soft-mass matrices at the flavor-breaking scale as specified in Eq. (A1) of the Appendix. In this way, we try to include generic models with different numbers of flavons or different contributions to the soft-mass matrices.

As concluded on Sec. III, the most important contributions to d_e come first from the $(\delta^e_{\mathrm{LL}})_{13}(\delta^e_{\mathrm{LR}})_{33}(\delta^e_{\mathrm{RR}})_{31}$ insertion, and then from $(\delta^e_{\mathrm{LR}})_{1i}(\delta^e_{\mathrm{RR}})_{i1}$. How significant the latter is depends on the values of A_0 and $\tan\beta$. In the numerical analysis we assume one $O(1)$ phase on each rephasing-invariant combination, and thus, here it is enough to put it on $(\delta^e_{\mathrm{RR}})_{31}$. This allows us to estimate the largest area in the m_0-$M_{1/2}$ plane into which EDM experiments could probe.

In Figs. 1 and 2 we show contours for expected values of $|d_e|$ in the m_0-$M_{1/2}$ plane. The red, orange, and yellow (dark gray, medium gray, and light gray) regions show contours for $|d_e|$ equal to 1×10^{-28}, 5×10^{-29}, and $1\times10^{-29}e$ cm, respectively.

Figure 1 takes $A_0=0$ and $\tan\beta=10,30$. In this case all off-diagonal δ_{LR} terms are generated by the running, and are thus small. d_e is then basically due to $(\delta^e_{\mathrm{LL}})_{13}(\delta^e_{\mathrm{LR}})_{33}\times(\delta^e_{\mathrm{RR}})_{31}$ insertions. It is important to emphasize that the present bound on d_e does not provide a constraint on the SUSY parameter space even for $\tan\beta=30$. However, by reaching a sensitivity of $10^{-29}e$ cm, we can explore values of $M_{1/2}$ and m_0 of the order of 1500 GeV. A value of $5\times10^{-29}e$ cm for the electron EDM would explore the parameter space up to values of $M_{1/2}$ and m_0 of order 700 GeV. In other words, a reasonable value of the electron EDM in these flavor models in the presence of large phases would be of the order of $1.1\times10^{-28}e$ cm for $\tan\beta=10$ and $3.0\times10^{-28}e$ cm for $\tan\beta=30$ with $(m_0,M_{1/2})=(500,300)$ GeV in both cases, corresponding to an accessible sfermion spectrum at the LHC with squark masses around 900 GeV. Thus, if large flavor phases are present and SUSY is to solve the SM hierarchy problem, we can hope to find some signature of d_e in the upcoming experiments, even for low values of $\tan\beta$ [20].

In Fig. 2 we set $A_0=m_0$ $(a_0=1)$. As expected, the $(\delta^e_{\mathrm{LR}})_{13}(\delta^e_{\mathrm{RR}})_{31}$ insertion comes into play especially for large m_0, due to the influence of the trilinears (remember we take $A_0=m_0$, i.e. it is not fixed to a single value), as they lower slepton masses in the RGEs. The expected

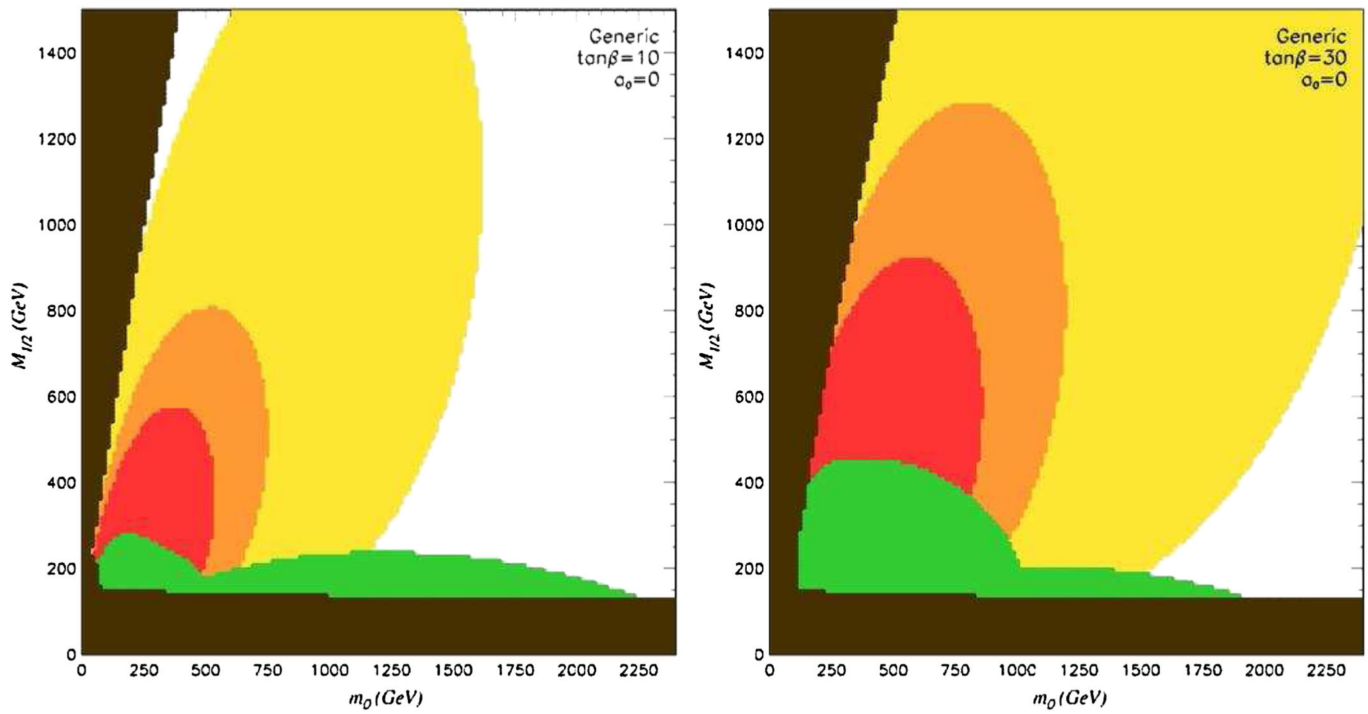

FIG. 1 (color online). Contours of $|d_e| = 1 \times 10^{-28} e$ cm (red/dark gray), $|d_e| = 5 \times 10^{-29} e$ cm (orange/medium gray), and $|d_e| = 1 \times 10^{-29} e$ cm (yellow/light gray) in the m_0-$M_{1/2}$ plane for $\tan\beta = 10$, 30 and $A_0 = 0$. Current FCNC constraints and direct LEP bounds are also shown in green (gray) and dark brown (black), respectively.

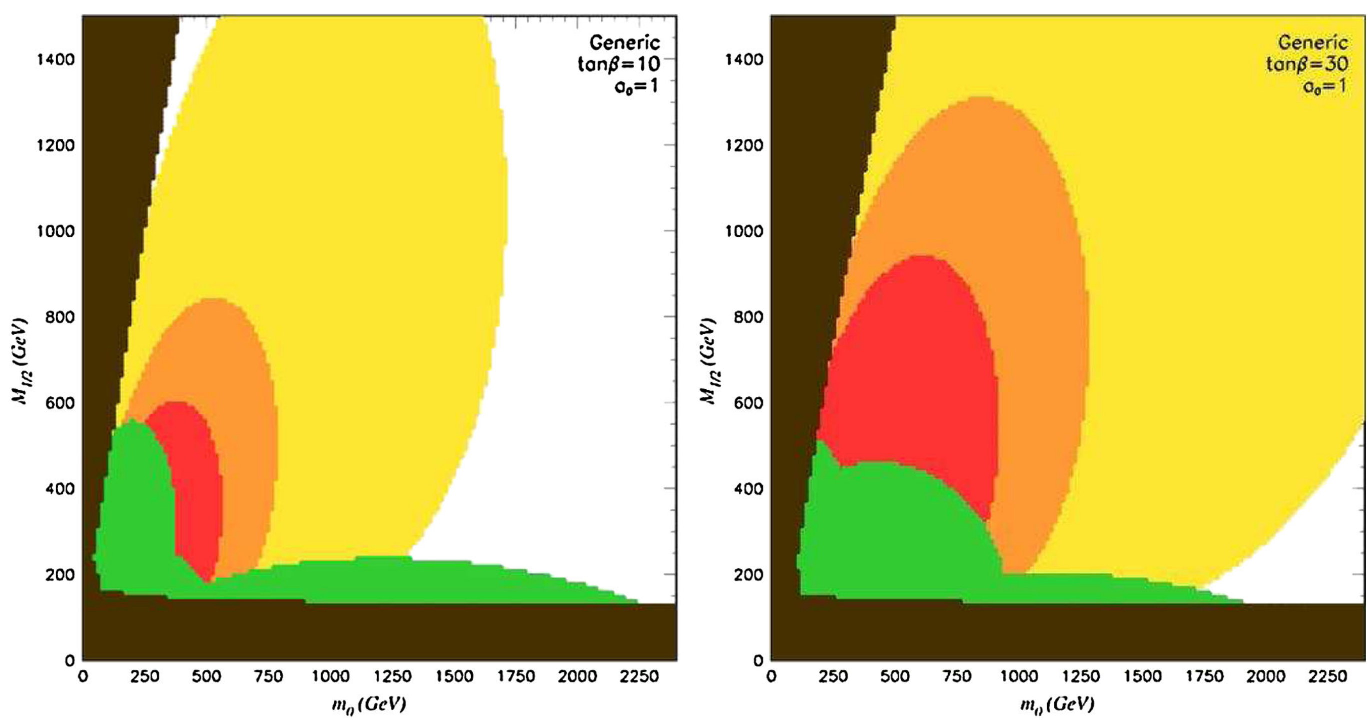

FIG. 2 (color online). Contours of $|d_e|$ and current constraints in the m_0-$M_{1/2}$ plane for $\tan\beta = 10$, 30 and $a_0 = 1$. See caption of Fig. 1 for the meaning of different regions.

values of d_e are now similar to the values found in the $A_0 = 0$ case, although they can be slightly increased in the regions of low $\tan\beta$ and large m_0 (A_0 for $a_0 = 1$). For comparison with the $A_0 = 0$ case, with $(m_0, M_{1/2}) = (500, 300)$ GeV we obtain $d_e \simeq 1.2 \times 10^{-28} e$ cm for $\tan\beta = 10$ and $3.4 \times 10^{-28} e$ cm for $\tan\beta = 30$. Therefore, we see that, in general, they are slightly increased, although we must keep in mind that having a positive or negative interference will depend on the phases.

Figure 3 gives details on the current constraints given by LFV experiments and ϵ_K, which were shown previously in green (gray). We show the bounds of $\mu \to e\gamma$ (green/medium gray), $\tau \to \mu\gamma$ (yellow/light gray), and ϵ_K (red/dark gray). It is interesting to notice that, even though strong, the bounds still allow a very large area of the parameter space which is compatible with the observation of SUSY at the LHC. The reach of the MEG experiment is also shown in Fig. 3, dotted in green (medium gray). We assume it is capable of reaching a sensitivity to the $\mu \to e\gamma$ branching ratio of 10^{-13} [84]. We also show the reach of $\tau \to \mu\gamma$ experiments at the Super Flavour Factory, hatched in yellow (light gray). We assume the experiment will be able to measure the branching ratio down to 2×10^{-9} [89].

The impact of MEG in these $SU(3)$ flavor models on the evaluated parameter space is impressive, covering values of $M_{1/2} \lesssim 1500$ GeV and $m_0 \lesssim 2500$ GeV for $\tan\beta = 30$. Thus, if any evidence of SUSY is to be found at the LHC, $\mu \to e\gamma$ decay should be seen at MEG. The same can be said for $\tau \to \mu\gamma$ at the Super Flavour Factory, even though such constraints are not as strong for low $\tan\beta$.

The main effect of $A_0 \neq 0$ can be clearly seen in Fig. 4 in the decay $\mu \to e\gamma$. The appearance of a $(\delta^e_{\rm LR})_{12}$ term implies a considerable new neutralino contribution. This new contribution can then interfere with the previous neutralino-chargino diagrams. Positive or negative interference depends on the relative sign (phase) between $(\delta^e_{\rm LR})_{12}$ and $(\delta^e_{\rm LL})_{12}$, as can be seen in Eqs. (15) and (16). From this figure we can see that this new contribution is indeed large and even dominant for $a_0 = 1$, especially for low values of $\tan\beta$.

This new contribution is almost $\tan\beta$ independent, so we are allowed to put strong bounds directly on the value of A_0. Notice that the constraints coming from the *current* bounds of $\mu \to e\gamma$ at small m_0 and $M_{1/2}$ are already very strong. The MEG prediction, shown in Fig. 4, will now cover the parameter space up to values of $M_{1/2} \lesssim 1500$ GeV and $m_0 \lesssim 1000$ GeV for $\tan\beta = 10$ and $M_{1/2} \lesssim 1500$ GeV and $m_0 \lesssim 2500$ GeV for $\tan\beta = 30$.

The $\tau \to \mu\gamma$ branching ratio is not affected as much by the flavor violating trilinear terms. The reason for this is that the other dominant insertions are proportional to $m_\tau \tan\beta$, as explained in Sec. IV. In contrast, d_μ and d_τ, even though larger than d_e, cannot be probed by the upcoming experiments. As explained in Sec. III, in these flavor models, the muon EDM is typically 2 orders of

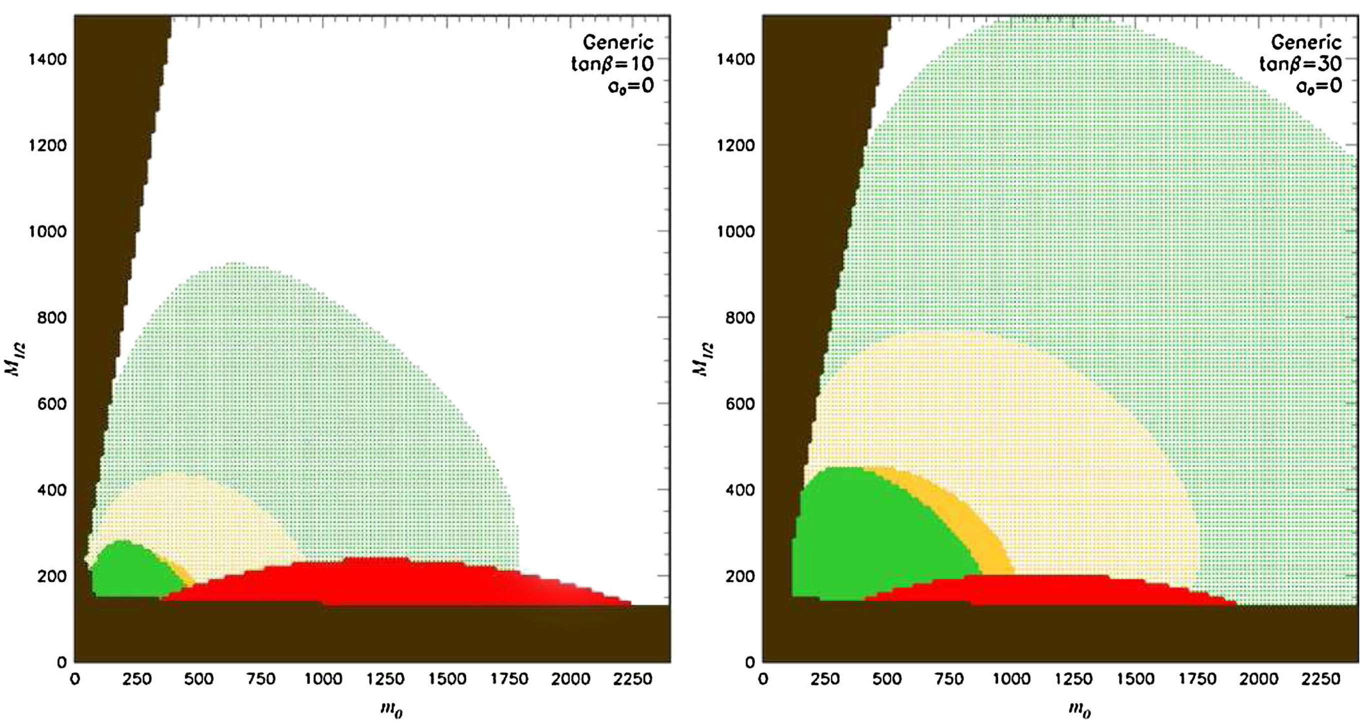

FIG. 3 (color online). Current constraints due to $\mu \to e\gamma$ (green/medium gray), $\tau \to \mu\gamma$ (yellow/light gray), and ϵ_K (red/dark gray) in the m_0-$M_{1/2}$ plane for $\tan\beta = 10$, 30 and $A_0 = 0$. The green (medium gray) dotted region corresponds to the reach of $\mu \to e\gamma$ in the MEG experiment, while the yellow (light gray) hatched region is the reach of $\tau \to \mu\gamma$ at the Super Flavour Factory.

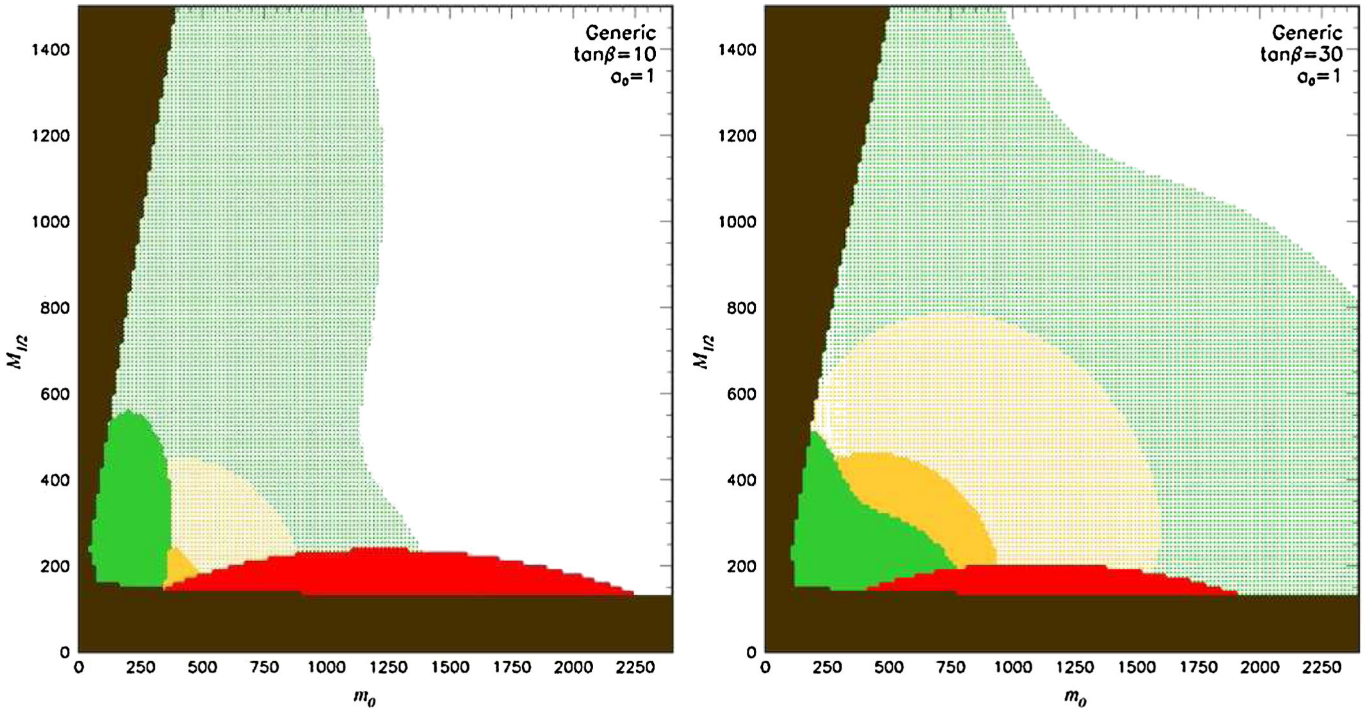

FIG. 4 (color online). Current constraints due to $\mu \to e\gamma$ (green/medium gray), $\tau \to \mu\gamma$ (yellow/light gray), and ϵ_K (red/dark gray) in the m_0-$M_{1/2}$ plane for $\tan\beta = 10$, 30 and $a_0 = 1$. The green (medium gray) dotted region corresponds to the reach of $\mu \to e\gamma$ in the MEG experiment, while the yellow (light gray) hatched region is the reach of $\tau \to \mu\gamma$ at the Super Flavour Factory.

magnitude larger than the electron EDM. In this case this would imply that d_μ can be at most of the order of $10^{-26}e$ cm, still orders of magnitude below the reach of the proposed experiments.

B. Generalization of the RVV model with spontaneous *CP*

In this section, we analyze a variation of the model defined in [18] and presented in the Appendix. As can be seen in Eqs. (A4), we have only one physical phase at leading order in the soft-mass matrices in the lepton sector: $(\beta_3 - \chi)$. Notice that this is due to the fact that, in this model, the soft-breaking terms have the "minimal" structure given by Eq. (3). This can change if different operators like $\theta_{23,i}\bar{\theta}_3^j$ + H.c. contribute to the soft-mass matrices.

With $A_0 = 0$, we can see that the leading phases in the $(\delta^e_{\rm LL})_{13}$ and $(\delta^e_{\rm RR})_{31}$ are equal, and therefore they cancel in $(\delta^e_{\rm LL})_{13}(\delta^e_{\rm LR})_{33}(\delta^e_{\rm RR})_{31}$.[7] The most significant contribution to d_e shall depend on the phase of the subleading term L_1, $2(\beta_3 - \chi)$. Because of this fact, and taking into account $\varepsilon/\bar{\varepsilon} = 1/3$, we can expect d_e to be smaller than in the generic case roughly by a factor of 10. This is confirmed by the numerical results shown in Fig. 5.

On the other hand, for $A_0 \neq 0$, the phase $2(\beta_3 - \chi)$ appears at leading order in the combination $(\delta^e_{\rm LR})_{13} \times (\delta^e_{\rm RR})_{31}$. As we have seen in the previous section, with phases $O(1)$ in all MIs, this contribution is comparable to the $(\delta^e_{\rm LL})_{13}(\delta^e_{\rm LR})_{33}(\delta^e_{\rm RR})_{31}$ contribution in the large m_0 (and thus large A_0) region. Therefore, in this model, where $(\delta^e_{\rm LL})_{13}(\delta^e_{\rm LR})_{33}(\delta^e_{\rm RR})_{31}$ is reduced by roughly an order of magnitude with respect to the generic case, we can expect the double mass insertion to dominate in most of the parameter space. In fact, for small $\tan\beta$, the moduli of this double MI is of the same order of magnitude as the triple MI and so we can expect to reach d_e values similar to those obtained in the generic case. Even for larger $\tan\beta \simeq 30$ the double MI is comparable to the triple MI in the large m_0 region where again the results for d_e are similar to those obtained in the generic case. Therefore, with $A_0 \neq 0$ we can expect similar values for d_e as in the generic model. This can be seen in Fig. 6.

The discussions of the constraints coming from MEG and Super Flavour Factories are completely analogous to those of the generic model, since these LFV processes do not depend heavily on the presence of sizable phases. Therefore Figs. 3 and 4 remain valid also in this model.

In summary, in this generalization of the RVV model with fixed phases, we would explore values of $M_{1/2}$ and m_0 of the order of 800 and 600 GeV with a value of $1 \times 10^{-29}e$ cm for the electron EDM with $\tan\beta = 10$ and $A_0 = 0$. This means we would need an increase of 10 in

[7]Notice that these phases are observable and will contribute to other CP violation observables [76].

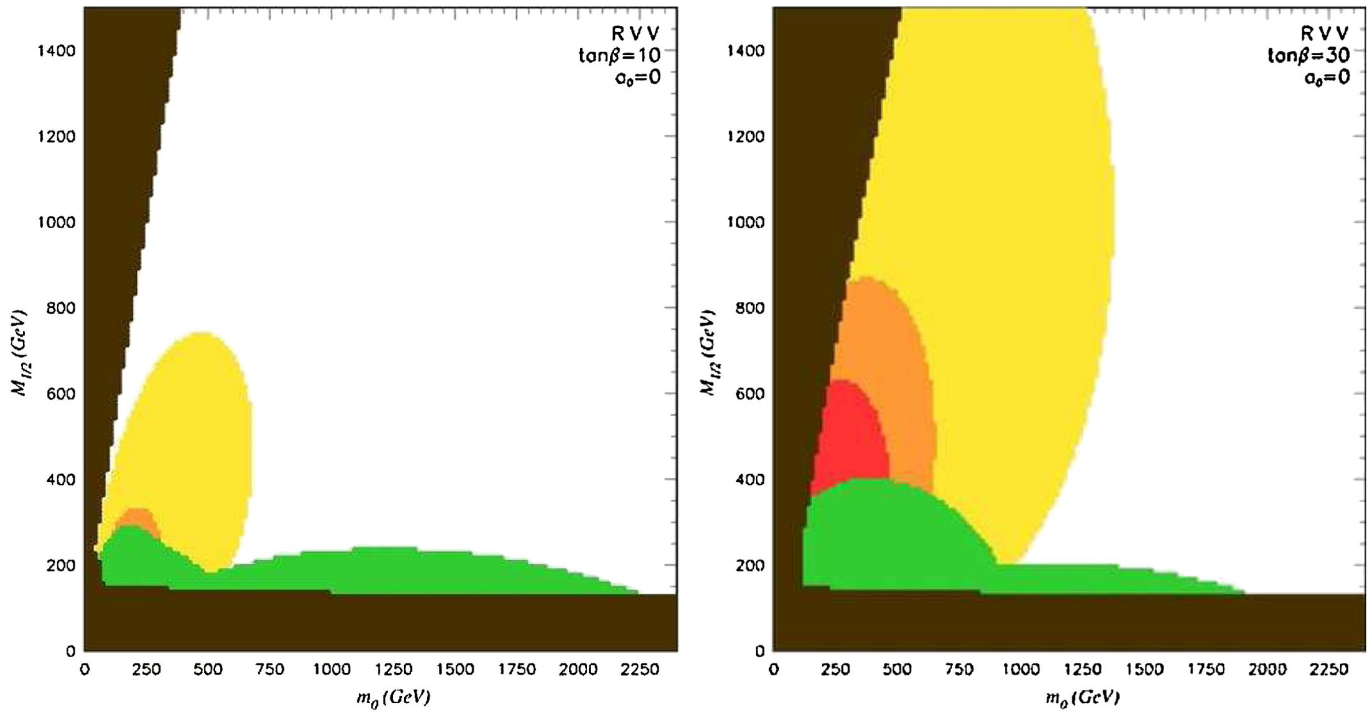

FIG. 5 (color online). Contours of $|d_e|$ as in Fig. 1, for $A_0 = 0$, but with phases following the RVV model.

the sensitivity to d_e to explore the same region of parameter space as in the generic model. However, if $a_0 = 1$ we would explore a region very similar to the region explored in the generic $SU(3)$ model, with similar values of $M_{1/2}$ and values of m_0 roughly smaller by a factor of 2. Reasonable values of the electron EDM in this explicit example for $(m_0, M_{1/2}) = (500, 300)$ GeV and $A_0 = 0$ would be of the order of $1.8 \times 10^{-29} e$ cm for $\tan\beta = 10$ and $8.3 \times 10^{-29} e$ cm for $\tan\beta = 30$. Thus, we see d_e is reduced by roughly a factor of 5 with respect to the generic

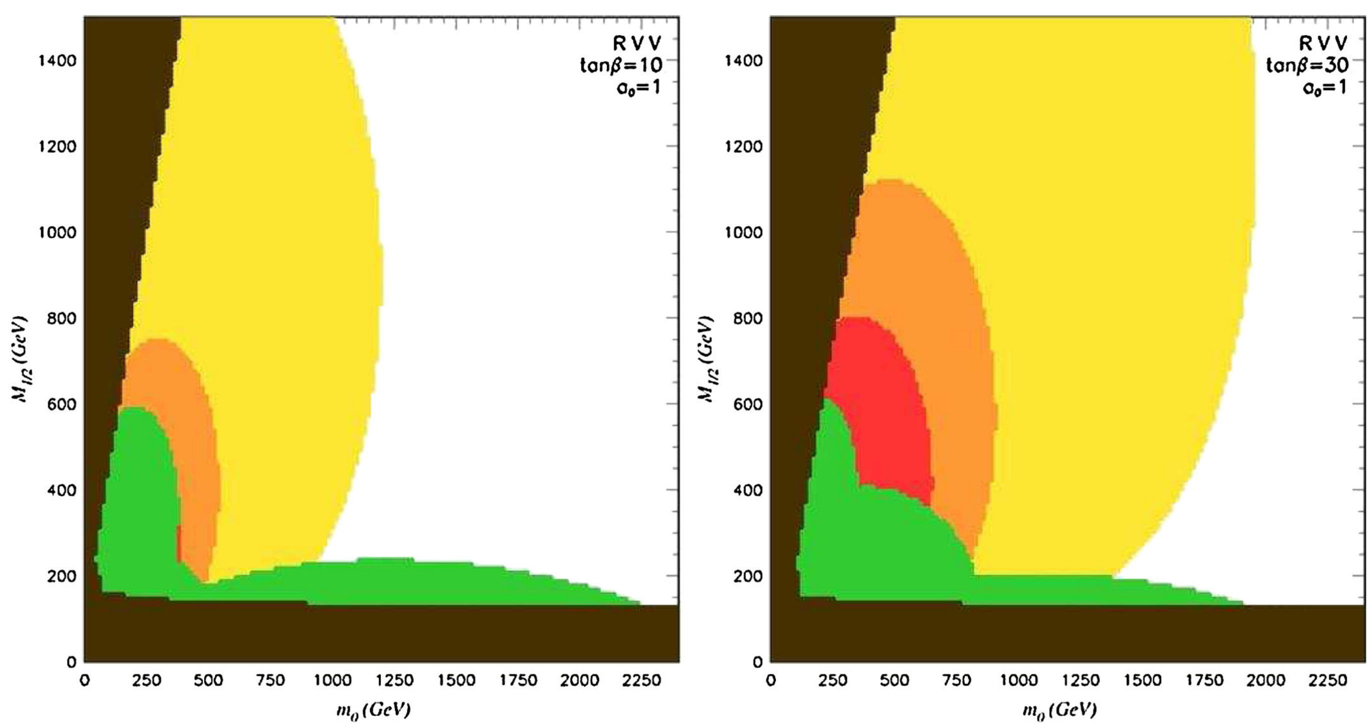

FIG. 6 (color online). Contours of $|d_e|$ as in Fig. 1, for $A_0 = m_0$, but with phases following the RVV model.

model. For the same scalar and gaugino masses and $A_0 = m_0$, $d_e \simeq 5.6 \times 10^{-29} e$ cm for $\tan\beta = 10$ and $d_e \simeq 1.7 \times 10^{-28} e$ cm for $\tan\beta = 30$, i.e., only a factor of 2 smaller than in the generic model.

VI. CONCLUSIONS

We have shown that the flavor and *CP* problems in the supersymmetric extensions of the SM are deeply related to the origin of flavor (and *CP*) in the Yukawa matrices. It is natural to think that the same mechanism generating the flavor structures and giving rise to *CP* violation in the Yukawa couplings is responsible for the structure and phases in the SUSY soft-breaking terms. A flavor symmetry with spontaneous *CP* violation in the flavor sector can simultaneously solve both problems.

In this paper, we have analyzed the phenomenology of non-Abelian $SU(3)$ flavor symmetry. In this model, flavor-independent phases are naturally zero and only flavor-dependent phases are present in the soft-breaking terms. We have studied the contributions to the leptonic electric dipole moments from these flavor phases, and we have shown that the future bounds on the electron EDM will be able to explore a large part of the SUSY parameter space in these models. Simultaneously, we have analyzed the reach of the future MEG and Super Flavour Factories through lepton flavor violation processes. We have shown that we can expect signals of new physics both in EDM and LFV experiments if the SUSY masses are accessible at the LHC.

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APPENDIX: SOFT MATRICES IN THE SCKM BASIS

In this appendix we present the structure of the soft-mass matrices in the SCKM basis in generic flavor symmetry models with a symmetric texture in the Yukawa matrices. We then present an explicit example based on the model of Ref. [18].

In order to make a mass insertion analysis of a given process, one must have all soft matrices in the SCKM basis where Yukawa matrices are diagonal at the electroweak scale. However, we present here the structure of the soft matrices in the SCKM basis at the flavor scale, and we will add the running effects later. For a generic model, assuming general phases (although taking into account that Yukawa and trilinear matrices are symmetric) and neglecting $O(1)$ constants, we have

$$Y_e = \begin{pmatrix} \frac{\bar\varepsilon^4}{3} & 0 & 0 \\ 0 & 3\bar\varepsilon^2 & 0 \\ 0 & 0 & 1 \end{pmatrix} y_{33}^e, \tag{A1a}$$

$$\frac{A_e}{A_0} = \begin{pmatrix} \frac{\bar\varepsilon^4}{3} & \bar\varepsilon^3 e^{i\alpha_1} & \bar\varepsilon^3 e^{i\beta_1} \\ \bar\varepsilon^3 e^{i\alpha_1} & 3\bar\varepsilon^2 & 3\bar\varepsilon^2 e^{i\gamma_1} \\ \bar\varepsilon^3 e^{i\beta_1} & 3\bar\varepsilon^2 e^{i\gamma_1} & 1 \end{pmatrix} y_{33}^e, \tag{A1b}$$

$$\frac{(m^2_{\tilde e_R})^T}{m_0^2} = \begin{pmatrix} 1+\bar\varepsilon^2 y_{33}^e & \frac13 \bar\varepsilon^3 e^{i\alpha_2} & \frac13 \bar\varepsilon^3 e^{i\beta_2} \\ \frac13 \bar\varepsilon^3 e^{-i\alpha_2} & 1+\bar\varepsilon^2 & \bar\varepsilon^2 e^{i\gamma_2} \\ \frac13 \bar\varepsilon^3 e^{-i\beta_2} & \bar\varepsilon^2 e^{-i\gamma_2} & 1+y_{33}^e \end{pmatrix}, \tag{A1c}$$

$$\frac{m^2_{\tilde L}}{m_0^2} = \begin{pmatrix} 1+\varepsilon^2 y_{33}^\nu & \frac13 \varepsilon^2\bar\varepsilon e^{i\alpha_3} & \bar\varepsilon^3 y_{33}^\nu e^{i\beta_3} \\ \frac13 \varepsilon^2\bar\varepsilon e^{-i\alpha_3} & 1+\varepsilon^2 & \bar\varepsilon^2 y_{33}^\nu e^{i\gamma_3} \\ \bar\varepsilon^3 y_{33}^\nu e^{-i\beta_3} & \bar\varepsilon^2 y_{33}^\nu e^{-i\gamma_3} & 1+y_{33}^\nu \end{pmatrix} \tag{A1d}$$

where $y_{33}^e = (\langle\theta_3^d\rangle/M_d)^2 = m_\tau/(v\cos\beta)$ and $y_{33}^\nu = (\langle\theta_3^u\rangle/M_d)^2 = m_t/(v\sin\beta)$.

In the following we will present an explicit example of an $SU(3)$ flavor symmetry model with spontaneous *CP* violation and the phase structure of the different matrices completely determined. This model is a generalization of the RVV model of Ref. [18] with arbitrary values of $\tan\beta$. The charges of the different flavon fields under the $SU(3)$ and global symmetries are shown in Table III.

TABLE III. Charges required to build satisfactory Yukawa matrices in the RVV model.

Field	ψ	ψ^c	H	Σ	θ_3	θ_{23}	$\bar\theta_3$	$\bar\theta_{23}$
$SU(3)$	3	3	1	1	$\bar 3$	$\bar 3$	3	3
R	1	1	0	3	-2	-2	-1	2
$U(1)$	1	1	-2	-1	0	1	-1	0
Z_3	1	1	2	0	1	1	-1	2

After spontaneous breaking of the flavor symmetry (and *CP* symmetry), the vevs of the different fields are

$$\begin{aligned} \langle\theta_3\rangle &= \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix} \otimes \begin{pmatrix} a_3^u & 0 \\ 0 & a_3^d e^{i\chi} \end{pmatrix}, \\ \langle\bar\theta_3\rangle &= \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix} \otimes \begin{pmatrix} a_3^u e^{i\alpha_u} & 0 \\ 0 & a_3^d e^{i\alpha_d} \end{pmatrix}, \\ \langle\theta_{23}\rangle &= \begin{pmatrix} 0 \\ b_{23} \\ b_{23} e^{i\beta_3} \end{pmatrix}, \\ \langle\bar\theta_{23}\rangle &= \begin{pmatrix} 0 \\ b_{23} e^{i\beta_2'} \\ b_{23} e^{i(\beta_2'-\beta_3)} \end{pmatrix}, \end{aligned} \tag{A2}$$

where we require the following relations:

$$\left(\frac{a_3^u}{M_u}\right)^2 = y_{33}^\nu, \qquad \left(\frac{a_3^d}{M_d}\right)^2 = y_{33}^e, \qquad \frac{b_{23}}{M_u} = \varepsilon, \qquad \frac{b_{23}}{M_d} = \bar{\varepsilon}. \tag{A3}$$

In the SCKM basis, and still neglecting $O(1)$ constants, the SUSY breaking matrices are

$$\frac{A_e}{A_0} = \begin{pmatrix} \frac{\bar{\varepsilon}^4}{3} & \bar{\varepsilon}^3 & \bar{\varepsilon}^3 e^{i(\beta_3-\chi)} \\ \bar{\varepsilon}^3 & 3\bar{\varepsilon}^2 & 3\bar{\varepsilon}^2 e^{i(\beta_3-\chi)} \\ \bar{\varepsilon}^3 e^{i(\beta_3-\chi)} & 3\bar{\varepsilon}^2 e^{i(\beta_3-\chi)} & 1 \end{pmatrix} y_{33}^e, \tag{A4a}$$

$$\frac{(m^2_{\tilde{e}_R})^T}{m_0^2} = \begin{pmatrix} 1+\bar{\varepsilon}^2 y_{33}^e & \frac{1}{3}\bar{\varepsilon}^3 & \frac{1}{3}\bar{\varepsilon}^3 e^{-i(\beta_3-\chi)} \\ \frac{1}{3}\bar{\varepsilon}^3 & 1+\bar{\varepsilon}^2 & E^*\bar{\varepsilon}^2 e^{-i(\beta_3-\chi)} \\ \frac{1}{3}\bar{\varepsilon}^3 e^{i(\beta_3-\chi)} & E\bar{\varepsilon}^2 e^{i(\beta_3-\chi)} & 1+y_{33}^e \end{pmatrix}, \tag{A4b}$$

$$\frac{m^2_{\tilde{L}}}{m_0^2} = \begin{pmatrix} 1+\varepsilon^2 y_{33}^\nu & \frac{1}{3}\varepsilon^2\bar{\varepsilon} & L_1^*\bar{\varepsilon}^3 y_{33}^\nu e^{-i(\beta_3-\chi)} \\ \frac{1}{3}\varepsilon^2\bar{\varepsilon} & 1+\varepsilon^2 & 3L_2^*\bar{\varepsilon}^2 y_{33}^\nu e^{-i(\beta_3-\chi)} \\ L_1\bar{\varepsilon}^3 y_{33}^\nu e^{i(\beta_3-\chi)} & 3L_2\bar{\varepsilon}^2 y_{33}^\nu e^{i(\beta_3-\chi)} & 1+y_{33}^\nu \end{pmatrix}, \tag{A4c}$$

where the terms E, L_1, and L_2 include important subdominant contributions with physical phases [L_1 and L_2 differ by $O(1)$ constants]:

$$E = 1 - 3y_{33}^e e^{-2i(\beta_3-\chi)}, \tag{A5}$$

$$L_i = 1 - \frac{1}{3y_{33}^\nu}\frac{\varepsilon^2}{\bar{\varepsilon}^2} e^{-2i(\beta_3-\chi)}. \tag{A6}$$

In order to reproduce down-quark and electron masses for all values of $\tan\beta$, we take $y_{33}^e = \langle\theta_3^d\rangle^2$ different from $y_{33}^\nu = \langle\theta_3^u\rangle^2$.

Notice that, within this structure, the only physical phase is $\beta_3 - \chi$. This is an additional phase with respect to the CKM phase ω in [18].

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