Electric dipole moments from flavoured CP violation in SUSY

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

I. INTRODUCTION

The so-called Supersymmetry (SUSY) flavour and CP problems are usually taken as the main naturalness problems SUSY has to face up. The formulation of the supersymmetric flavour 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, FCNC and flavour-dependent CP violation observables would receive too large contributions from loops involving SUSY particles to satisfy the stringent phenomenological bounds on these processes. 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 flavour-independent terms typically generate too large contributions to the so-far unobserved electric dipole moment (EDM) of the electron and neutron.

The basis of both problems lies clearly on our total ignorance about the origin of the observed flavour and CP-violation in our theory. However, notice that these problems are not restricted
to supersymmetry. Even the Standard Model (SM) shares the flavour 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 flavour problem in the SM than in the MSSM. The real flavour problem is simply our inability to understand the complicated structures in the quark and lepton Yukawa couplings, and likewise for the soft-breaking flavour 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 flavour-independent phases in the MSSM, both facts can suggest the idea that the flavour and the CP problems are indeed related and solving the flavour problem while restricting the CP phases to the flavour sector would also solve the CP problem.

A particularly attractive solution to these problems (both in the SM and in SUSY) is found on models based on flavour symmetries. In these models, the flavour structure of the Yukawa matrices is only generated after the breaking of a flavour symmetry \[ 5, 6, 7, 8, 9, 10, 11, 12, 13, 14 \] and the flavour structure of the SUSY soft-breaking terms would also originate through the same mechanism \[ 15, 16, 17, 18, 19 \]. Thus, finding a solution to the SM flavour problem will generally solve at the same time the so-called SUSY flavour problem to a sufficient degree, although probably still allowing naturally suppressed contributions that might bring more information about the flavour sector. Regarding the SUSY CP problem, if we want to restrict all CP phases in SUSY to the flavour sector, this can be achieved by postulating an exact CP symmetry spontaneously broken in the flavour sector. This would remove all flavour-independent phases, but still produce interesting observable sources of CP violation.

In this work we shall analyze how a flavour symmetry can solve the flavour and CP problems in a SUSY scenario. We shall also distinguish typical signatures of such a symmetry in the lepton sector\(^1\). In Section \[\text{II}\] we shall analyze a general flavour symmetry model which can solve both

\(^1\) Observables in the quark sector, such as the neutron EDM, shall be studied in a future work \[20\].
problems, based on [18]. In Section III we shall identify the most important flavour-dependent contributions to leptonic EDMs, in order to quantify the expected amount of CP violation in such models. In Section IV we shall study in a similar manner the relevant lepton flavour violation (LFV) processes. Finally, in Section V we will show the regions of the SUGRA parameter space that are sensitive to future EDM and LFV experiments.

II. FLAVOUR SYMMETRIES AND SPONTANEOUS CP BREAKING

Following the original ideas of Froggatt-Nielsen [6], flavour symmetries explain the peculiar structure of the SM Yukawa couplings as the result of a spontaneously broken symmetry associated with flavour. 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 non-renormalizable 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 flavour 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 flavour 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 flavour-diagonal couplings are clearly invariant under any symmetry, i.e. diagonal soft-masses are always allowed by the flavour 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, 21, 22, 23] family symmetry, the symmetry does not relate different generations and therefore diagonal masses can be different. In this case, Flavour Changing Neutral Current (FCNC) and CP violation phenomenology set very strong constraints on the differences between these flavour diagonal masses and Abelian family symmetries have serious difficulties to satisfy 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 $SU(2)$ flavour models [11, 25], where the first two generation sfermions, facing the strongest constraints, share a common mass. In the case of an $SU(3)$ flavour symmetry [18, 26, 27, 28], all
three generations have the same mass in the unbroken family symmetry limit. For this reasons, in
the following we will consider non-Abelian family symmetries, and more precisely $SU(3)$ flavour
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 non-renormalizable
operators neutral under the flavour 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 non-trivial flavour structure in the soft-breaking parameters. Therefore, by being
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 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 sector 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:

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

with $\bar{\varepsilon} \approx 0.15$, $\varepsilon \approx 0.05$, $b = 1.5$, $a = 1.3$, $c = 0.4$ and $a', b', c'$ are poorly fixed from experimental
data [18, 22, 37]. Unfortunately, the Yukawa couplings in the leptonic sector can not 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 in 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) flavour matrices are the same except for the different vev of a Georgi-Jarlskog Higgs field \([38]\) to unify the second and first generation masses.

With this Yukawa structure as our starting point, we will generate the flavour structure of the soft-breaking terms in an explicit example based in an \(SU(3)\) flavour 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: \(3 \times 3 \times 1\) are not allowed by the \(SU(3)\). In the theory, we have several flavon fields, that we call \(\theta_3, \theta_{23}\) (anti-triplets \(\overline{3}\)), \(\overline{\theta}_3\) and \(\overline{\theta}_{23}\) (triplets \(3\)). The symmetry is broken in two steps, first \(\theta_3\) and \(\overline{\theta}_3\) get a vev, \(\propto (0, 0, 1)\), breaking \(SU(3)\) into \(SU(2)\). Subsequently a smaller vev of \(\theta_{23}\) and \(\overline{\theta}_{23}\), \(\propto (0, 1, 1)\), breaks the remaining symmetry \([26]\).

To reproduce the Yukawa textures, the large third generation Yukawa couplings require a \(\theta_3\) (and \(\overline{\theta}_3\)) 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 \(\theta_{23}/M_f\) (and \(\overline{\theta}_{23}/M_f\)) have vevs of order \(\epsilon\) in the up sector and \(\bar{\epsilon}\) 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 \(\theta_{23}^i \theta_{23}^j\left(\theta_3 \overline{\theta}_3\right)^2\) (with possibly the Georgi-Jarlskog field \(\Sigma\) to unify quark and leptonic Yukawa couplings, not included here for simplicity, see Refs. \([18, 26, 27]\)). The couplings in the first row and column of the Yukawa matrix are given by \(\epsilon^{ikl} \theta_{23}^i \theta_{23}^j \left(\theta_3 \overline{\theta}_3\right)^2\) where to reproduce the texture in Eq. (1) we must force \(n = 1\). Unfortunately, the \(SU(3)_f\) 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

\[
W_Y = H \psi_1 \psi_2^C \left[\theta_3^i \theta_3^j + \theta_{23}^i \theta_{23}^j \left(\theta_3 \overline{\theta}_3\right)^2 + \epsilon^{ikl} \theta_{23}^i \theta_{23}^j \left(\theta_3 \overline{\theta}_3\right)^2 + \ldots\right],
\]

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 \(\theta_i/M_f\). This structure is quite general for the different \(SU(3)\) models we can build, and for additional details we refer to

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\(2\) Notice that we add the scalar product \((\theta_3 \overline{\theta}_3)\) to this operator with respect to Refs. \([18, 26, 27]\) to be able to generalize to different \(\tan \beta = ((\theta_{23}^i)/\theta_{23}^j)^2\) values.
In the same way, after \( SU(3) \) breaking the scalar soft masses deviate from exact universality. As explained above, \( \phi_i^+ \phi_i \) 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 \cite{18,39,40,41}. 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 \(^3\):

\[
(M_f^2)^{ij} = m_0^2 \left( \delta^{ij} + \frac{1}{M_f^2} \left[ \bar{\theta}_{23}^3 \bar{\theta}_{23}^3 + \bar{\theta}_{23}^3 \bar{\theta}_{23}^3 + \bar{\theta}_{23}^3 \bar{\theta}_{23}^3 \right] + \frac{1}{M_f^2} \left( e^{ikl} \bar{\theta}_{23}^3 \bar{\theta}_{23}^3 + \bar{\theta}_{23}^3 \bar{\theta}_{23}^3 \right) \right),
\]

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 flavour symmetry must commute with the SM symmetry and therefore the vector-like mediator fields must have the SM quantum numbers of the usual particles\(^4\). With this terms, we can see some similarities between the flavour structure of the Yukawa and soft-mass matrices. In particular, the off-diagonal (2,3) elements in the Yukawa matrix is given by \( \bar{\theta}_{23}^3 \bar{\theta}_{23}^3 / M_f^2 \simeq \epsilon_f^2 \) (the term \( \bar{\theta}_{23}^3 \bar{\theta}_{23}^3 \) factorizes from all the contributions in the Yukawa matrix), while in the soft masses it is \( \bar{\theta}_{23}^3 \bar{\theta}_{23}^3 / M_f^2 \), also of order \( \epsilon_f^2 \). Still, it is possible to build other invariant combinations with different flavon fields that can not be present in the superpotential. This is due to the fact that the superpotential must be holomorphic, i.e. can not include daggered fields, while the soft masses, coming from the Kähler potential, are only a real combination of fields. In some cases, the global symmetries can allow a combination like \( \theta_{23,i} \bar{\theta}_{23,j}^\dagger + h.c. \) to the soft masses, even though the combination \( \theta_{23,i} \bar{\theta}_{23,j} \) is forbidden by global symmetries in the superpotential. This structures would affect strongly the phenomenology of third generation physics \cite{20}. In any case, although possible, this situation is very rare and usually the structure of the soft terms follow the Yukawa structure. It is important to emphasize at this point that these deviations from universality in the soft-mass matrices proportional to flavour symmetry breaking come always 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.

\(^3\) Discrete non-Abelian subgroups of \( SU(3) \) \cite{29,30,31,32,33,34,35,36} 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 mass \( M_d \) and \( M_u = M_L \) for simplicity.
| EDM  | Current Bound (e cm) | Future Bound (e cm) |
|------|----------------------|---------------------|
| \(|d_e|\) | \(\leq 1.4 \times 10^{-27} [48]\) | \(\sim 10^{-32} [52]\) |
| \(|d_\mu|\) | \(\leq 7.1 \times 10^{-19} [49]\) | \(\sim 10^{-23} [53]\) |
| \(|d_\tau|\) | \(\leq 2.5 \times 10^{-17} [50]\) | \(\sim 10^{-20} [54]\) |
| \(|d_n|\) | \(\leq 2.9 \times 10^{-26} [51]\) | \(\sim 10^{-28} [55]\) |

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

In other mediation mechanism, as gauge-mediation \([43]\) or anomaly mediation \([44, 45, 46]\), where these Kähler contributions to the soft masses are negligible, flavour 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 the flavour 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 flavour 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 SCKM basis does not diagonalize the trilinear matrices. In fact, the trilinear matrices maintain the same structure as in the flavour basis and only the \(O(1)\) coefficients are modified.

We should now take into account, specially in a gravity-mediated scenario, the canonical normalization of the kinetic terms (Kähler potential). However as shown in Refs. \([40, 47]\) these canonical transformations do not modify the general structure of the different flavour matrices and change only the unknown \(O(1)\) coefficients. As, at this level, these coefficients can not be fixed by the flavour symmetry, the previous discussion on the different flavour matrices remains valid in the canonical basis.

We have now fixed the flavour structure that we can expect in the soft-breaking matrices. Now, we are ready to consider the problem of CP violation in these flavour matrices. In fact, in the SM, CP violation is deeply related with flavour. 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 flavour independent phases in the \( \mu \) term and trilinear couplings, Supersymmetry gives rise to contributions to the EDMs at 1 loop order with no suppression associated to flavour. These one-loop contributions, for SUSY masses below several TeV, can easily exceed the present experimental bounds shown on Table by one or two 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.

In this aspect, flavour symmetries with spontaneous breaking of CP provide an interesting solution. CP is an exact symmetry of the theory before the breaking of the flavor symmetry. Thus, above the scale of flavour breaking, all terms in the Kähler potential, which gives rise to the soft masses and the \( \mu \) term (by the Giudice-Masiero mechanism), are real. Even after the breaking of flavour and CP symmetries, \( \mu \) receives flavour-suppressed complex corrections only at the two-loop level. 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. In this way, the Supersymmetric CP problem is naturally solved. Nevertheless, flavour-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 to quantify and explain all important contributions (as in [70]). In terms of the slepton mass matrix terms, these MIs are defined as:

\[
\begin{align*}
(\delta_{LR}^e)_{ij} &= \frac{v_d}{\sqrt{2}M_\tilde{e}^2} [A_{e,ij}^* - \delta_{ij} Y_i \mu \tan \beta] = (\delta_{RL}^e)^*_{ji}, \\
(\delta_{LL}^e)_{ij} &= \frac{(m_\tilde{l}^2)_{ij}}{M_\tilde{e}^2}, \quad (\delta_{RR}^e)_{ij} = \frac{(m_\tilde{\tau}^2)_{ij}}{M_\tilde{e}^2}
\end{align*}
\]

(4)

where \( M_\tilde{e}^2 \) 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 \( \tau - \mu \) sector respectively:

\[
\begin{align*}
(\delta_{LL}^e)_{12} &\approx \frac{\epsilon_e^2}{3}; \quad (\delta_{RR}^e)_{12} \approx \frac{\epsilon_\tau^2}{3}; \quad (\delta_{LR}^e)_{12} \approx A_0 \frac{m_\tilde{\tau}^2 \epsilon_\tau^2}{M_\tilde{e}^2} \\
(\delta_{LL}^e)_{13} &\approx \epsilon_e^2 \gamma_{33}^\nu; \quad (\delta_{RR}^e)_{13} \approx \epsilon_\tau^2 \gamma_{33}^\nu; \quad (\delta_{LR}^e)_{13} \approx A_0 \frac{m_\tilde{\tau}^2 \epsilon_\tau^2 \gamma_{33}^\nu}{M_\tilde{e}^2} \\
(\delta_{LL}^e)_{23} &\approx \epsilon_e^2 \gamma_{33}^\nu; \quad (\delta_{RR}^e)_{23} \approx \epsilon_\tau^2; \quad (\delta_{LR}^e)_{23} \approx A_0 \frac{m_\tilde{\tau}^2 \epsilon_\tau^2 \gamma_{33}^\nu}{M_\tilde{e}^2}
\end{align*}
\]

(5a,b,c)

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

$$(m_{\tilde{e}}^2)_{32}(M_{\text{SUSY}}) \simeq \left( (m_{\tilde{e}}^2)_{32}(M_U) \left(1 - \frac{1}{16\pi^2} (2y_e^2) \right) - 4\frac{A^e_{33}A^{e\dagger}_{32}}{16\pi^2} \right) \log \left( \frac{M_U}{M_{\text{SUSY}}} \right)$$

where $M_U$ is the unification scale, $M_{\text{SUSY}}$ 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 $A_{13}^e/A_0 \sim (Y_\nu)_{13} \sim \mathcal{O}(\bar{\epsilon}^3)$, $A_{23}^e/A_0 \sim (Y_\nu)_{23} \sim \mathcal{O}(\bar{\epsilon}^2)$, $A_{33}^e/A_0 \sim m_\tau/(v \cos \beta)$ 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 sizeable correction from the running only in the $\tau - \mu$ and $\tau - e$ sectors, i.e. where the third generation is involved.

### III. EDMs and Flavour Phases

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

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

being related to the imaginary part of a chirality-changing, flavour-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 [56]. This makes EDMs very convenient observables where to look for new physics related to CP violation.

As we have seen above, if CP is spontaneously broken in the flavour sector, the usual flavour-independent phases coming from $\mu$ and $A_f$ are approximately zero and we expect EDMs to be under control\(^5\). Nevertheless, flavour-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 on 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

\(^5\) 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.
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 one for the whole rephasing invariant combination, using the expected size of the offdiagonal elements in the flavour 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 flavour 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 \( \mu \), we can neglect the chargino contribution to \( d_e \), and concentrate on neutralinos completely. Neutralino contributions to \( d_e \) are:

\[
\Delta \chi^0_e = \left( \frac{\epsilon}{16\pi^2} \right) \sum m \left( A_{eij}^L A_{eij}^R \right) \frac{1}{m_{\chi_j^0}} F \left( \frac{m_{\chi_j^0}^2}{m_{\tilde{l}_i}^2} \right) 
\]

\[
A_{eij}^L = Y_e N_{3j}^* R_{\tilde{e}_R i} - \frac{g'}{\sqrt{2}} N_{1j}^* R_{\tilde{e}_L i} - \frac{g}{\sqrt{2}} N_{2j}^* R_{\tilde{e}_L i}
\]

\[
A_{eij}^R = Y_e N_{3j}^* R_{\tilde{e}_L i} + \sqrt{2} g' N_{1j}^* R_{\tilde{e}_R i}
\]

with \( R_{\tilde{e}_i} \) and \( N_{kj} \) being elements of the charged slepton and neutralino mixing matrices \[77\], respectively and the loop function \( F(x) = \frac{x}{2(x-1)^2} (x^2 - 1 - 2x \log x) \).

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

\[
\frac{d_e}{\epsilon} = \frac{\alpha M_1}{8\pi \cos^2 \theta_W M_\tilde{e}^2} \sum \left[ (\delta^e_{LL})_{1i} (\delta^e_{LR})_{i1} f_1 + (\delta^e_{LR})_{1i} (\delta^e_{RR})_{i1} f_2 + (\delta^e_{LL})_{1i} (\delta^e_{LR})_{ij} (\delta^e_{RR})_{j1} f_3 \right]
\]

where \( M_\tilde{e}^2 \) is the average slepton mass and the loop functions, \( f_i \), can be derived from \[70\].

Let us explain briefly why these are the relevant contributions, and identify the dominant ones. All phases in this \( SU(3) \) flavour model are contained within the sfermion mass matrices and thus we shall need at least one mass insertion on the slepton line. Since all flavour-conserving insertions \((\delta^e_{LR})_{ii}\) are real to leading order, we will need to combine at least two flavour-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 through two binos (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 binos 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 \((\delta^e_{LR})_{1i}(\delta^e_{RR})_{i1}\) or \((\delta^e_{LL})_{1i}(\delta^e_{LR})_{i1}\). With three mass insertions, we consider only contributions with a single LR insertion. It is well-known that each \(\delta^e_{LR}\) insertion is suppressed by a cumulative factor \(m_\tau/M_\tilde{e}\). Therefore the dominant contribution comes from \((\delta^e_{LL})_{1i}(\delta^e_{LR})_{ij}(\delta^e_{RR})_{j1}\). Obviously, the largest contribution is the one that involves a central \((\delta^e_{LR})_{33}\), due to the \(m_\tau\tan\beta\) enhancement. Thus, the most important contribution to \(d_e\) with three mass insertions is the pure bino \((\delta^e_{LL})_{13}(\delta^e_{LR})_{33}(\delta^e_{RR})_{31}\).

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-insertions contribution shall dominate no matter the size of \((\delta^e_{LR})_{ij}\), 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:

\[
(\delta^e_{LR})_{1i}(\delta^e_{RR})_{i1} \approx A_0 \epsilon^{\tilde{e}} \frac{m_\tau}{M_\tilde{e}}
\]

\[
(\delta^e_{LL})_{1i}(\delta^e_{LR})_{ij}(\delta^e_{RR})_{j1} \approx \epsilon^6 \frac{m_\tau m_\nu}{M_\tilde{e}} \tan\beta,
\]

and therefore for \(A_0 \simeq M_\tilde{e}\) the triple mass insertion is dominant except for small values of \(\epsilon y_{\tilde{e}3}\tan\beta\).

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

The situation for \(d_\tau\) 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 $(\delta_{LL}^e)_{32}(\delta_{LR}^e)_{22}(\delta_{RR}^e)_{23}$. This would be smaller than $d_\mu$ by a factor $m_\mu/m_\tau$, almost two orders of magnitude, and thus $d_\tau$ clearly violates the naive scaling relation. In contrast, when $A_0 \neq 0$, the main contributions from the double MIs with a flavour changing $\delta_{LR}$ insertion are identical in magnitude to those for $d_\mu$ and so, we expect $d_\tau$ to be of size comparable to $d_\mu$. In this case, we should take into account the possible presence of subdominant phases in the SCKM basis in flavour diagonal trilinear couplings [39]. In any case, this breaks again the mass scaling relation, even though not so drastically as in the previous situation. The observation of such a bizarre behavior would be a very clear signal favoring these type of flavour models.

IV. LEPTON FLAVOUR VIOLATING DECAYS

As discussed in the previous sections, supersymmetric flavour models are characterized by non-universal 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 determines the arising of potentially large mixing among flavours. In particular, in the lepton sector, the same mass insertions which induce EDMs are sources of lepton flavour 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 $BR(l_i \rightarrow l_j \gamma)$.

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

$$BR(l_i \rightarrow l_j \gamma) = \frac{48\pi^2\alpha}{G_F^2}(|A_L^{ij}|^2 + |A_R^{ij}|^2),$$

(14)

with the SUSY contribution to each amplitude given by the sum of two terms $A_{L,R} = A_{L,R}^N + A_{L,R}^C$, where $A_{L,R}^N$ and $A_{L,R}^C$ 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:

$$A_L^{ij} = \frac{\alpha_2}{4\pi} \left( \frac{\delta_{LL}^e}_{ij} \right) \left[ \frac{\mu M_2 \tan \beta}{(M_2^2 - \mu^2)} F_{2LL}(a_2, b) + \tan^2 \theta_W \frac{\mu M_1 \tan \beta}{(M_1^2 - \mu^2)} F_{1LL}(a_1, b) \right]$$

(15)

and

$$A_R^{ij} = \frac{\alpha_1}{4\pi} \left( \frac{\delta_{RR}^e}_{ij} \right) \left( \frac{M_1}{m_{l_i}} \right) F_{1LR}(a_1),$$

(16)
where $\theta_W$ is the weak mixing angle, $a_{1,2} = M_{1,2}^2/\bar{m}^2$, $b = \mu^2/m_{\tilde{\tau}}^2$ and the loop functions $F_{1\text{LL}}, F_{2\text{LL}}, F_{1\text{RR}}$ and $F_{1\text{LR}}$ can be obtained from the expressions in Refs. [80, 81].

Here we can see that $(\delta^e_{\text{LL}})_{ij}$ and $(\delta^e_{\text{RR}})_{ij}$ contributions are $\tan\beta$-enhanced. In contrast to MFV models with RH neutrinos, in our model the largest contribution to $\mu \to e\gamma$ comes from the RR sector, being $(\delta^e_{\text{RR}})_{12} \simeq \bar{e}^3$ while $(\delta^e_{\text{LL}})_{12} \simeq \bar{e}^2 \bar{e}$ (see Section 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 $(\delta^e_{\text{LR}})_{ij}$ arises from pure $\tilde{B}$ exchange and it is completely independent of $\tan\beta$. However, in these flavour models LR mass insertions can be still 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 $(\delta^e_{\text{RR}})_{12} m_\mu \tan\beta$ with $(\delta^e_{\text{LR}})_{12} M_1$, as we can see from Eq. (16). Using the expression for the mass insertions of Eq. (5c), we see that in these models:

$$\frac{(\delta^e_{\text{RR}})_{12} m_\mu \tan\beta}{(\delta^e_{\text{LR}})_{12} M_1} \simeq \frac{\bar{e}^3/3}{\bar{e}^3 A_0} \frac{m_\mu \tan\beta}{(m_\tau/M_2^2)} M_1 = m_\mu \tan\beta \frac{M_2^2}{3 m_\tau A_0 M_1}. \quad (17)$$

Therefore, we can see that if $M_2^2/(A_0 M_1) \sim O(1)$, 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 $BR(\mu \to e\gamma)$ 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 $BR(\mu \to e\gamma)$. 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 $m_\tau/m_\mu$ 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 $(\delta^e_{\text{LL}})_{32}$, which are $\tan\beta$-enhanced. Therefore, we expect the $BR(\tau \to \mu\gamma)$ 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 $\delta^e_{\text{LL}}$, which contributes both to chargino and neutralino diagrams. Thus a rough estimation for the relative sizes of the branching ratios can be:

$$\frac{BR(\tau \to e\gamma)}{BR(\mu \to e\gamma)} \approx \left(\frac{m_\tau}{m_\mu}\right) 5 \frac{\Gamma_\mu (\delta^e_{\text{LL}})_{13}^2}{\Gamma_\tau (\delta^e_{\text{LL}})_{12}^2} \approx O(1) \quad (18)$$

$$\frac{BR(\tau \to \mu\gamma)}{BR(\mu \to e\gamma)} \approx \left(\frac{m_\tau}{m_\mu}\right) 5 \frac{\Gamma_\mu (\delta^e_{\text{LL}})_{23}^2}{\Gamma_\tau (\delta^e_{\text{LL}})_{12}^2} \approx O(10^3) \quad (19)$$

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

It is important also to clarify the dependence of our results on the chosen value of $Y_\nu$. Notice that, in our $SU(3)$ model, the value of the neutrino Yukawa couplings are fixed by the symmetry following \cite{18}. However, different values of $Y_\nu$ could be possible in other examples while respecting 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 $m_{\tilde{L}}^2$ depend on the value of $y_{33}^\nu$ at 1 loop. Therefore only the predictions on $\tau \rightarrow l_i\gamma$ can be affected by a change on $y_{33}^\nu$ and even in this case the contributions from $m_{\tilde{e}}^2$, independent on $y_{33}^\nu$, 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 SUGRA 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 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 $M_{\text{flav}} = 2 \times 10^{16}$ GeV, as explained in Section III and explicitly shown in the Appendix. Then we evolve the different flavour matrices to the electroweak scale using 1-loop RGEs \cite{87}. $O(1)$ coefficients in the Yukawas matrices are determined by requiring a good fit on the fermion masses and quark mixings at $M_Z$ \cite{37}. 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 $(a_u^3/a_d^3)$, that 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 column in all soft mass matrices and in the case of left-handed sleptons generates also the dominant contribution to $(M_{\tilde{L}}^2)_{23}$. At this scale, we apply the LEP bounds on the lightest sparticle and Higgs masses and we require a neutral 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
TABLE II: Applied constraints coming from EDMs and LFV (left) and neutral meson sector (right).

| Observable | Bound | Observable | Bound |
|------------|-------|------------|-------|
| $d_e$      | $< 1.4 \times 10^{-27}$ e cm | $\Delta m_K$ | $< 3.48 \times 10^{-15}$ GeV |
| $BR(\mu \to e\gamma)$ | $< 1.2 \times 10^{-11}$ | $\Delta m_D$ | $< 4.61 \times 10^{-14}$ GeV |
| $BR(\tau \to \mu\gamma)$ | $< 4.5 \times 10^{-8}$ | $\Delta m_B$ | $< 3.34 \times 10^{-13}$ GeV |
| $BR(\tau \to e\gamma)$ | $< 1.1 \times 10^{-7}$ | $\Delta m_{B_s}$ | $< 1.17 \times 10^{-11}$ GeV |
| $BR(b \to s\gamma)^{\text{SUSY}}$ | $< 0.88 \times 10^{-4}$ | $\epsilon_K$ | $< 2.239 \times 10^{-3}$ |

current bounds from $d_e$, LFV processes and FCNC measurements on the hadronic sector: $\Delta m_K$, $\epsilon_K$, $\Delta m_D$, $\Delta m_B$, $\Delta m_{B_s}$, and $b \to s\gamma$ [88], as shown in Table II. Nonetheless, only $\mu \to e\gamma$, $\tau \to \mu\gamma$ and $\epsilon_K$ 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 (grey).

Initially, we present the results in a generic $SU(3)$ model with phases $O(1)$ in the flavour soft terms as shown in the Appendix, Eq. (A1). Then, we take the RRV model as an explicit example with a well-defined phase structure as shown in Eq. (A1)\(^6\).

A. Generic $SU(3)$ Model

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

As concluded on Section III, the most important contributions to $d_e$ come first from the $(\delta^e_{\text{LL}})_{13}(\delta^e_{\text{LR}})_{33}(\delta^e_{\text{RR}})_{31}$ insertion, and then from $(\delta^e_{\text{LR}})_{1i}(\delta^e_{\text{RR}})_{i1}$. How significant is the latter depends on the value 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 $(\delta_{\text{RR}}^e)_{31}$. This allows us to estimate the largest area in the $m_0 - M_{1/2}$ plane into which EDM experiments could probe.

In Figures 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 grey, medium grey and light grey) regions show contours for $|d_e|$ equal to $1 \times 10^{-28}$, $5 \times 10^{-29}$ and $1 \times 10^{-29}$ e cm, respectively.

Figure 1 takes $A_0 = 0$ and $\tan \beta = 10, 30$. In this case all off-diagonal $\delta_{\text{LR}}$ terms are generated

\(^6\)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.
by the running, and are thus small. $d_e$ is then basically due to $(\delta_e^{\text{LL}})_{13}(\delta_e^{\text{LR}})_{33}(\delta_e^{\text{RR}})_{31}$ 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 $5 \times 10^{-29}$ 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 flavour models in the presence of large phases would be of the order of $1.1 \times 10^{-28}$ e cm for $\tan\beta = 10$ and $3.0 \times 10^{-28}$ 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 the 900 GeV. Thus, if large flavour 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 Figure 2 we set $A_0 = m_0$ ($a_0 = 1$). As expected, the $(\delta_e^{\text{LR}})_{13}(\delta_e^{\text{RR}})_{31}$ 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 values of $d_e$ are now similar to the values found in the $A_0 = 0$ case, although 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

FIG. 1: Contours of $|d_e| = 1 \times 10^{-28}$ e cm (red/dark grey), $|d_e| = 5 \times 10^{-29}$ e cm (orange/medium grey) and $|d_e| = 1 \times 10^{-29}$ e cm (yellow/light grey) 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 (grey) and dark brown (black) respectively.
FIG. 2: 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 Figure 1 for the meaning of different regions.

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 $\epsilon_K$, which were shown previously in green (grey). We show the bounds of $\mu \rightarrow e\gamma$ (green/medium grey), $\tau \rightarrow \mu\gamma$ (yellow/light grey) and $\epsilon_K$ (red/dark grey). 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 Figure 3, dotted in green (medium grey). We assume it capable of reaching a sensitivity to the $\mu \rightarrow e\gamma$ branching ratio of $10^{-13}$ [84]. We also show the reach of $\tau \rightarrow \mu\gamma$ experiments at the Super Flavour Factory, hatched in yellow (light grey). We take the experiment to be able to measure the branching ratio down to $2 \times 10^{-9}$ [89].

The impact of MEG in these SU(3) flavour 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, $\mu \rightarrow e\gamma$ decay should be seen at MEG. The same can be said for $\tau \rightarrow \mu\gamma$ 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 Figure 4 in the decay $\mu \rightarrow e\gamma$. The appearance of a ($\delta^e_{LR})_{12}$ term implies a considerable new neutralino contribution. This new contribution can
FIG. 3: Current constraints due to $\mu \rightarrow e\gamma$ (green/medium grey), $\tau \rightarrow \mu\gamma$ (yellow/light grey) and $\epsilon_K$ (red/dark grey) in the $m_0$-$M_{1/2}$ plane for $\tan \beta = 10, 30$ and $A_0 = 0$. The green (medium grey) dotted region corresponds to the reach of $\mu \rightarrow e\gamma$ at the MEG experiment, while the yellow (light grey) hatched region is the reach of $\tau \rightarrow \mu\gamma$ at the Super Flavour Factory.

FIG. 4: Current constraints due to $\mu \rightarrow e\gamma$ (green/medium grey), $\tau \rightarrow \mu\gamma$ (yellow/light grey) and $\epsilon_K$ (red/dark grey) in the $m_0$-$M_{1/2}$ plane for $\tan \beta = 10, 30$ and $A_0 = 1$. The green (medium grey) dotted region corresponds to the reach of $\mu \rightarrow e\gamma$ at the MEG experiment, while the yellow (light grey) hatched region is the reach of $\tau \rightarrow \mu\gamma$ at the Super Flavour Factory.
then interfere with the previous neutralino-chargino diagrams. Positive or negative interference
depends on the relative sign (phase) between $(\delta_{LR})_{12}$ and $(\delta_{LL})_{12}$ 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 Figure 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 flavour violating trilinear terms. The
reason for this is that the other dominant insertions are proportional to $m_\tau \tan \beta$, as explained in
Section IV. In contrast, $d_\mu$ and $d_\tau$, even though larger than $d_e$, can not be probed by the upcoming
experiments. As explained in section III in these flavour models, the muon EDM is typically two
orders of magnitude larger than the electron EDM. In this case this would imply that $d_\mu$ 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 Eq. (A1), 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 $\theta_{23} i \bar{\theta}_3 + h.c.$ contribute to the soft mass matrices.

With $A_0 = 0$, we can see that the leading phases in the $(\delta^e_{LL})_{13}$ and $(\delta^e_{RR})_{31}$ are equal and
therefore they cancel in $(\delta^e_{LL})_{13}(\delta^e_{LR})_{33}(\delta^e_{RR})_{31}$. The most significant contribution to $d_e$ shall
depend on the phase of the subleading term $L_1$, $2(\beta_3 - \chi)$. Due to 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 Figure 5.

On the other hand, for $A_0 \neq 0$, the phase $2(\beta_3 - \chi)$ appears at leading order in the combination $(\delta^e_{LR})_{13}(\delta^e_{RR})_{31}$. As we have seen in the previous section, with phases $O(1)$ in all MIs, this

\footnote{Notice that these phases are observable and will contribute to other CP violation observables.}
contribution is comparable to the \((\delta^e_{LL})_{13}(\delta^e_{LR})_{33}(\delta^e_{RR})_{31}\) contribution in the large \(m_0\) (and thus large \(A_0\)) region. Therefore, in this model, where \((\delta^e_{LL})_{13}(\delta^e_{LR})_{33}(\delta^e_{RR})_{31}\) 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 similar \(d_e\) values 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 Figure 6.

The discussion 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 sizeable 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 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 roughly a factor of 5 with respect to the generic 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 two smaller than in the generic model.

VI. CONCLUSIONS

We have shown that the flavour and CP problems in the supersymmetric extensions of the SM are deeply related to the origin of flavour (and CP) in the Yukawa matrices. It is natural to think that the same mechanism generating the flavour 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 flavour symmetry with spontaneous CP violation in the flavour sector can simultaneously solve both problems.

In this paper, we have analyzed the phenomenology of a non-Abelian $SU(3)$ flavour symmetry. In this model, flavour-independent phases are naturally zero and only flavour-dependent phases are present in the soft-breaking terms. We have studied the contributions to the leptonic electric dipole moments from these flavour 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 Flavour
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 A: SOFT MATRICES IN THE SCKM BASIS

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

In order to make a mass insertion analysis of a given process, one must have all soft matrices at 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 flavour 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
TABLE III: Charges required to build satisfactory Yukawa matrices in the RVV Model.

| Field | $\psi$ | $\psi^c$ | $H$ | $\Sigma$ | $\theta_3$ | $\theta_{23}$ | $\bar{\theta}_3$ | $\bar{\theta}_{23}$ |
|-------|--------|--------|-----|--------|---------|---------|---------|---------|
| SU(3) | 3      | 3      | 1   | 1      | 3       | 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       |

have:

$$Y_e = \begin{pmatrix} \frac{\epsilon^4}{3} & 0 & 0 \\ 0 & 3\bar{\epsilon}^2 & 0 \\ 0 & 0 & 1 \end{pmatrix} y_{33}^e$$  \hspace{1cm} (A1a)

$$\frac{A_{e}}{A_0} = \begin{pmatrix} \frac{\epsilon^4}{3} & \bar{\epsilon}^3 e^{i\alpha_1} & \bar{\epsilon}^3 e^{i\beta_1} \\ \bar{\epsilon}^3 e^{i\alpha_1} & 3\bar{\epsilon}^2 & 3\bar{\epsilon}^2 e^{i\gamma_1} \\ \bar{\epsilon}^3 e^{i\beta_1} & 3\bar{\epsilon}^2 e^{i\gamma_1} & 1 \end{pmatrix} y_{33}^e$$  \hspace{1cm} (A1b)

$$\frac{(m_{\tilde{e}_R}^2)^T}{m_0^2} = \begin{pmatrix} 1 + \bar{\epsilon}^2 y_{33}^e & \frac{1}{3}\bar{\epsilon}^3 e^{-i\alpha_2} & 1 + \bar{\epsilon}^2 e^{i\gamma_2} \\ \frac{1}{3}\bar{\epsilon}^3 e^{-i\beta_2} & \bar{\epsilon}^2 e^{-i\gamma_2} & 1 + y_{33}^e \end{pmatrix}$$  \hspace{1cm} (A1c)

$$\frac{m_{\tilde{e}_L}^2}{m_0^2} = \begin{pmatrix} 1 + \bar{\epsilon}^2 y_{33}^e & \frac{1}{3}\bar{\epsilon}^3 e^{-i\alpha_3} & \bar{\epsilon}^3 y_{33}^e e^{i\beta_3} \\ \frac{1}{3}\bar{\epsilon}^2 \bar{\epsilon} e^{-i\alpha_3} & 1 + \bar{\epsilon}^2 & \bar{\epsilon}^2 y_{33}^e e^{i\gamma_3} \\ \bar{\epsilon}^3 y_{33}^e e^{-i\beta_3} & \bar{\epsilon}^2 y_{33}^e e^{-i\gamma_3} & 1 + y_{33}^e \end{pmatrix}$$  \hspace{1cm} (A1d)

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

In the following we will present an explicit example of an $SU(3)$ flavour 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.

After spontaneous breaking of the flavour symmetry (and CP symmetry) the vevs of the different
fields are:

$$\langle \theta_3 \rangle = \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix} \otimes \begin{pmatrix} a_u^u \\ a_d^d e^{i\chi} \\ 0 \end{pmatrix}; \quad \langle \bar{\theta}_3 \rangle = \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix} \otimes \begin{pmatrix} a_u^u e^{i\alpha_u} \\ a_d^d e^{i\alpha_d} \\ 0 \end{pmatrix};$$

$$\langle \theta_{23} \rangle = \begin{pmatrix} 0 \\ b_{23} \\ b_{23} e^{i\beta_3} \end{pmatrix}; \quad \langle \bar{\theta}_{23} \rangle = \begin{pmatrix} 0 \\ b_{23} e^{i\beta_2} \\ b_{23} e^{(\beta_2-\beta_3)} \end{pmatrix};$$

where we require the following relations:

$$\left( \frac{a_u^u}{M_u} \right)^2 = \nu_{33}^u, \quad \left( \frac{a_d^d}{M_d} \right)^2 = \nu_{33}^d,$$

$$\frac{b_{23}}{M_u} = \varepsilon, \quad \frac{b_{23}}{M_d} = \bar{\varepsilon}.$$

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

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

$$\frac{m_{\nu_R}^2}{m_0^2} = \frac{1 + \varepsilon^2}{\frac{1}{3}\varepsilon^3} \begin{pmatrix} \frac{1}{3}\varepsilon^3 & \frac{1}{3}\varepsilon^3 e^{-i(\beta_3-\chi)} \\ 1 + \varepsilon^2 e^{i(\beta_3-\chi)} & 1 + \varepsilon^2 e^{-i(\beta_3-\chi)} \end{pmatrix} \quad (A4b)$$

$$\frac{m_{\tau}^2}{m_0^2} = \begin{pmatrix} 1 + \varepsilon^2 y_{33} & \frac{1}{3}\varepsilon^2 \varepsilon & L_1^* \varepsilon^3 y_{33} e^{-i(\beta_3-\chi)} \\ \frac{1}{3}\varepsilon^2 \varepsilon & 1 + \varepsilon^2 & 3L_2^* \varepsilon^2 y_{33} e^{-i(\beta_3-\chi)} \\ L_1 \varepsilon^3 y_{33} e^{-i(\beta_3-\chi)} & 3L_2 \varepsilon^2 y_{33} e^{-i(\beta_3-\chi)} & 1 + y_{33}' \end{pmatrix} \quad (A4c)$$

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):

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

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

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

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

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