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LFV and EDMs in SUSY with flavour symmetries

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Summary. — After more than four decades of impressive and frustrating theoretical and experimental efforts to reveal signals of the presence of TeV new physics through its effects in CP-conserving and CP-violating flavour-changing neutral-current processes, the main response seems to lie in an effective flavour blindness of the new physics at the electroweak scale (Minimal Flavour Violation). This perspective keeps still open the door for main surprises in the sector of lepton flavor violation. In this talk I focus on an alternative road where both the flavour puzzle within the Standard Model (*i.e.*, a rationale for the smallness of Yukawa couplings and fermion mixings) and the flavour problem of TeV new physics are simultaneously tackled in a supersymmetric extension of the Standard Model where the flavour structure is dictated by a (spontaneously broken) flavour or horizontal symmetry.

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PACS `12.10.-g` – Unified field theories and models.
PACS `12.60.Jv` – Supersymmetric models.
PACS `13.35.Bv` – Decays of muons.

1. – General status of flavour in our search for new physics

At least 40 years of efforts (and successes) in probing Flavor-Changing Neutral-Current (FCNC) phenomena have lead to the following conclusion: the Cabibbo-Kobayashi-Maskawa (CKM) flavour structure of the Standard Model (SM) represents the main bulk of flavour and (flavour-violating) CP violation in the hadronic sector at least down to distances of the order of $(100\,\text{GeV})^{-1}$. I think that, first of all, the relevance of this result should not be underestimated: this understanding represents a major breakthrough in our knowledge of the fundamental properties of Nature. As for the leptonic sector, the SM brilliantly succeeds to highly suppress charged lepton flavour violations (LFV) linking such suppression to the smallness of neutrino masses.

Unfortunately, the giant progress in our knowledge of flavour was not matched by a major breakthrough in our search for New Physics (NP) signals in rare FCNC and CPV processes. To be sure, we now know much more on the relation between the flavour structure and TeV NP, namely such low-energy NP should be flavour blind (*i.e.* the only flavour source would be given by the SM Yukawa couplings—the so-called MFV, Minimal Flavour Violation, assumption), or, if it possesses new flavour sources in addition to the SM Yukawa couplings, they should contribute to FCNC processes by no more than 10–20% of what the SM contributes. In both cases, it is clear that the TeV NP should be far from the generic case where it introduces new sources of flavour by its own without any specific suppression characteristic of the SM. Indeed, quite the opposite has to occur: the NP should enjoy a very stringent "flavour protection", either a total one forbidding any new source of flavour connected to the presence of NP (MFV framework) or, in any case, a very efficient suppression of the NP intrinsic FCNC contributions. As for the latter situation, this could obviously arise if the NP instead of being at the TeV scale should appear at a multi-TeV scale, but, in that case, the fine-tuning needed to ensure the correct energy scale for the electroweak breaking would become more and more severe. Alternatively, as usual in particle physics, the concept of "protection" immediately recalls the concept of "symmetry", with the possibility of a slightly broken symmetry to guarantee an adequate suppression of the FCNC NP contributions.

Before tackling this latter issue of the "symmetry protection" with a specific example, let me comment on the other possibility to reconcile TeV NP and FCNC suppression, namely the MFV case. From the theoretical point of view, it is not so simple to obtain a purely MFV situation. For instance, if the TeV NP is represented by a supersymmetric (SUSY) extension of the SM, then, in the case of supergravity (SUGRA), MFV is more an exception than a viable solution: indeed, barring the case of a purely dilaton-mediated SUGRA breaking (which seems hardly achievable, anyway), in all other cases where moduli take part in the process of SUGRA breaking we expect the scalar fermion masses not to be flavour universal, hence inducing a new source of flavour in FCNC SUSY contributions. On the other hand, if other mechanisms are adopted for SUSY breaking, for instance gauge (GMSB)—or anomaly (AMSB)—mediation, then strict MFV can be enforced. If MFV is present, then the chances to observe NP signals in hadronic FCNC processes become quite slim. There are possible exceptions to this grim scenario: for instance, in the case where two Higgs doublets are present, the rate of the process $B_s \to \mu^+\mu^-$ can be extraordinarily enhanced given the very large dependence of the process on the ratio of the two Higgs VEVs (the so-called $\tan\beta$ parameter). Alternatively, one can still try to search for departures from the SM expectations in hadronic FCNC processes, but to have concrete hopes to see something significant one has to go to Super-Flavour machines where accuracies at the percent level may be achieved (at the same time, one has to improve our theoretical accuracy reaching again the percent level, in particular in the calculation of the hadronic matrix elements).

In the case of MFV, the situation appears much more promising when we consider the leptonic sector, more specifically lepton flavour violation (LFV) in the charged lepton sector ($\mu \to e\gamma$, $\mu - e$ conversion in nuclei, $\tau \to \mu\gamma$, etc.). Here two facts play a crucial role and both are related to neutrinos: i) we have to provide a mass to neutrinos, hence we have to extend the particle spectrum of the SM, with the possible introduction of new flavour sources (for instance, new Yukawa couplings related to Yukawa terms where left- and right- handed neutrinos are put into communication); ii) we witness a large LFV in the neutrino sector—neutrino oscillations—with the possibility that such flavour changes may have implications in other sectors of the theory.

The links between the issues of TeV NP and FCNC on one side and neutrino flavour properties and implications on the other side have been since long exploited. More than twenty years ago, Francesca Borzumati and myself pointed out that in supergravities where neutrinos acquire a mass through the see-saw mechanism, even if the SUSY breaking entails universal sfermion masses, it is possible to (largely) misalign slepton and lepton mass matrices thanks to the influence of the neutrino Yukawa couplings in the running of the slepton masses from the scale at which the soft scalar masses appear down to the right-handed neutrino mass scale(s) [1]. In the 20 years elapsed from that work much (theoretical and experimental) progress has been made and with the new running or upcoming experiments in LFV we are certainly covering very interesting areas of parameters space in several TeV NP cases (see, for instance, the recent review by Junji Hisano [2] and the references quoted therein).

As for the above point ii), the relevance of the large neutrino mixings in other sectors of the theory appears in a striking way when we consider grand-unified theories (GUTs) in the context of SUGRA extensions of the SM. In those frameworks it can happen that the large LFV present in the neutrino sector can be "transferred" to some hadronic sector, typically the masses of the right-handed scalar quarks during their running down to the scale of the right-handed neutrino masses. I take this opportunity to remind here the work in this field of a collaborator of mine, Darwin Chang [3], an enthusiastic researcher in looking for NP signals, who has prematurely left us. The potentialities for NP FCNC contributions in SUGRA GUTs have been thoroughly investigated in a couple of recent papers, in particular emphasizing the intriguing possibility to constrain hadronic (leptonic) FCNC SUSY contributions making use of FCNC leptonic (hadronic) processes [4].

Coming back to the possibility that a symmetry may be the source of the "flavour protection" to solve the NP flavour problem, I am going to briefly report here on a recent work on this issue in collaboration with Lorenzo Calibbi, Joel Jones, Jae-hyeon Park, Werner Porod and Oscar Vives [5].

We were driven by the idea that the SM flavour puzzle (namely, the search for a rationale for fermion masses and mixings) and the NP flavour problem could find a simultaneous answer once the flavour properties of the SM and of the NP beyond it could emerge from a "flavour symmetry". In other words, once we have a theory of flavour this should be able to simultaneously account for the smallness of (some) Yukawa couplings and mixing angles in the SM as well as for the smallness of the FCNC contributions where NP particles run in the loops. The key for such solution can be an enlargement of the SM and NP symmetries with the presence of a flavour or horizontal symmetry. In the limit of exact flavour symmetry we would have a complete degeneracy of three fermion families together with their scalar partners in a SUSY extension of the SM. The (spontaneous) breaking of such symmetry originates the hierarchical structure in fermion families as we observe it as well as it gives rise to a specific pattern of non-universality in the masses of sfermions belonging to different generations.

We consider an $SU(3)$ flavour model. Under the $SU(3)$ flavor symmetry, the three generations of SM fields, both $SU(2)_L$-doublets and singlets, are triplets $\mathbf{3}$ and the Higgs fields are singlets. As flavons, we have θ_3, θ_{23} (anti-triplets $\bar{\mathbf{3}}$), $\bar{\theta}_3$ and $\bar{\theta}_{23}$ (triplets $\mathbf{3}$). The full superpotential is determined by $SU(3)$, and several global symmetries which forbid unwanted terms that would spoil the observed structure of the Yukawa couplings. Using an appropriately chosen set of charges [5], the leading terms in the superpotential are

$$(1)\quad W_{\rm Y} = H\psi_i\psi_j^c \Big[\theta_3^i\theta_3^j + \theta_{23}^i\theta_{23}^j\Sigma + \left(\epsilon^{ikl}\overline{\theta}_{23,k}\overline{\theta}_{3,l}\theta_{23}^j + \epsilon^{jkl}\overline{\theta}_{23,k}\overline{\theta}_{3,l}\theta_{23}^i\right)\left(\theta_{23}\overline{\theta}_3\right) + \epsilon^{ijl}\overline{\theta}_{23,l}\left(\theta_{23}\overline{\theta}_3\right)^2 + \epsilon^{ijl}\overline{\theta}_{3,l}\left(\theta_{23}\overline{\theta}_3\right)\left(\theta_{23}\overline{\theta}_{23}\right) + \dots\Big],$$

where the flavon fields have been normalized to the corresponding mediator mass, which means that all the flavon fields in this equation should be understood as θ_i/M_f. The field Σ is a Georgi-Jarlskog field that gets a vev in the $B-L$ direction, distinguishing leptons and quarks. Furthermore, this model is embedded in a $SO(10)$ grand-unified structure at high scales, which allow us to relate quark and lepton (including neutrino) Yukawa couplings. However, the $SU(2)_R$ subgroup of $SO(10)$ must be broken as we need different mediator masses for the up and down sector, and θ_3 and $\overline{\theta}_3$ are $\mathbf{3}\oplus\mathbf{1}$ representations of $SU(2)_R$ which is broken by their vev's [6-8].

The flavon fields get the following vev's:

$$\langle\theta_3\rangle = \begin{pmatrix}0\\0\\1\end{pmatrix}\otimes\begin{pmatrix}a_3^u & 0\\ 0 & a_3^d\,e^{i\chi}\end{pmatrix}; \qquad \langle\bar\theta_3\rangle = \begin{pmatrix}0\\0\\1\end{pmatrix}\otimes\begin{pmatrix}a_3^u\,e^{i\alpha_u} & 0\\ 0 & a_3^d\,e^{i\alpha_d}\end{pmatrix}; \tag{2}$$
$$\langle\theta_{23}\rangle = \begin{pmatrix}0\\ b_{23}\\ b_{23}\,e^{i\beta_3}\end{pmatrix}; \qquad \langle\bar\theta_{23}\rangle = \begin{pmatrix}0\\ b_{23}\,e^{i\beta_2'}\\ b_{23}\,e^{i(\beta_2'-\beta_3)}\end{pmatrix};$$

where we require the following relations:

$$\left(\frac{a_3^u}{M_u}\right)^2 = y_t, \quad \left(\frac{a_3^d}{M_d}\right)^2 = y_b, \quad \frac{b_{23}}{M_u} = \varepsilon, \quad \frac{b_{23}}{M_d} = \bar\varepsilon, \tag{3}$$

where $\bar\varepsilon \simeq 0.15$, $\varepsilon \simeq 0.05$. These relations are valid at the flavour breaking scale, that we take as the GUT scale in the numerical evaluation.

It is straightforward to see that this superpotential reproduces correctly the required Yukawa structure,

$$Y_d \propto \begin{pmatrix}0 & x_{12}^d\,\bar\varepsilon^3 & x_{13}^d\,\bar\varepsilon^3\\ x_{12}^d\,\bar\varepsilon^3 & \bar\varepsilon^2 & x_{23}^d\,\bar\varepsilon^2\\ x_{13}^d\,\bar\varepsilon^3 & x_{23}^d\,\bar\varepsilon^2 & 1\end{pmatrix}, \qquad Y_u \propto \begin{pmatrix}0 & x_{12}^u\,\varepsilon^3 & x_{13}^u\,\varepsilon^3\\ x_{12}^u\,\varepsilon^3 & \varepsilon^2 & x_{23}^u\,\varepsilon^2\\ x_{13}^u\,\varepsilon^3 & x_{23}^u\,\varepsilon^2 & 1\end{pmatrix}, \tag{4}$$

where x_{ij}^a are $O(1)$ coefficients fixed by the observed values of fermion masses and mixings.

We can now turn to the soft breaking terms. A universal, flavour diagonal mass term will always be allowed. Moreover, in a SUSY theory, the same messenger fields as in the Yukawas will couple the flavons to the scalar fields in the soft terms. Thus, the ε and $\bar\varepsilon$ parameters still act as expansion parameters, and represent important corrections to the soft terms.

Clearly any coupling involving a flavon field and its Hermitian conjugate (*i.e.* $\theta_3^{i\dagger}\theta_3^j$) is invariant under the flavour symmetry. From this we can deduce that the soft mass terms get a minimum structure determined uniquely by the flavon content of the model and their vev's. This minimum structure is obtained from the following effective terms:

$$(M_{\tilde f}^2)_i^j = m_0^2\Bigg(\delta_i^j + \left[\theta_{3i}^\dagger\theta_3^j + \overline{\theta}_{3,i}\overline{\theta}_3^{\dagger\,j} + \theta_{23i}^\dagger\theta_{23}^j + \overline{\theta}_{23,i}\overline{\theta}_{23}^{\dagger\,j}\right] \\ + \left(\epsilon_{ikl}\theta_3^k\theta_{23}^l\right)\left(\epsilon^{jmn}\theta_{3m}^\dagger\theta_{23n}^\dagger\right) + \left(\epsilon_{ikl}\overline{\theta}_3^{\dagger\,k}\overline{\theta}_{23}^{\dagger\,l}\right)\left(\epsilon^{jmn}\overline{\theta}_{3,m}\overline{\theta}_{23,n}\right) + \ldots\Bigg). \tag{5}$$

One can find a choice of global charges that reproduces the correct Yukawa structure and does not allow other terms at leading order in the Kähler potential (soft masses).

In the squark sector, after rephasing the fields such that the CKM matrix elements V_{ud}, V_{us}, V_{cb} and V_{tb} are real, the soft masses in the SCKM basis are

$$\text{(6a)} \qquad \left(M^2_{\tilde{u}^c_R}\right)^T = \begin{pmatrix} 1+\varepsilon^2\, y_t & -\varepsilon^3\, e^{i\omega'} & -\varepsilon^3\, e^{i(\omega'-2\chi)} \\ -\varepsilon^3\, e^{-i\omega'} & 1+\varepsilon^2 & \varepsilon^2\, e^{-2i\chi} \\ -\varepsilon^3\, e^{-i(\omega'-2\chi)} & \varepsilon^2\, e^{2i\chi} & 1+y_t \end{pmatrix} m_0^2,$$

$$\text{(6b)} \qquad \left(M^2_{\tilde{d}^c_R}\right)^T = \begin{pmatrix} 1+\bar{\varepsilon}^2\, y_b & -\bar{\varepsilon}^3\, e^{i\omega_{us}} & -\bar{\varepsilon}^3\, e^{i\omega_{us}} \\ -\bar{\varepsilon}^3\, e^{-i\omega_{us}} & 1+\bar{\varepsilon}^2 & \bar{\varepsilon}^2 \\ -\bar{\varepsilon}^3\, e^{-i\omega_{us}} & \bar{\varepsilon}^2 & 1+y_b \end{pmatrix} m_0^2,$$

$$\text{(6c)} \qquad M^2_{\tilde{Q}} = \begin{pmatrix} 1+\varepsilon^2\, y_t & -\varepsilon^2\bar{\varepsilon}\, e^{i\omega_{us}} & -\bar{\varepsilon}^3\, y_t e^{i\omega_{us}} \\ -\varepsilon^2\bar{\varepsilon}\, e^{-i\omega_{us}} & 1+\varepsilon^2 & \bar{\varepsilon}^2\, y_t \\ -\bar{\varepsilon}^3\, y_t e^{-i\omega_{us}} & \bar{\varepsilon}^2\, y_t & 1+y_t \end{pmatrix} m_0^2,$$

where $M^2_{\tilde{Q}}$ is in the basis where Y_d is diagonal. The phases ω_{us}, ω', and δ_i can be found in ref. [5]. The structure of $M^2_{\tilde{Q}}$ in the basis where Y_u is diagonal is similar to $M^2_{\tilde{u}^c_R}$. We have omitted $O(1)$ constants in front of each term, and subdominant terms which can have other phases as β_3 and χ. The slepton soft masses have the same structure, but can be numerically different, since they have a different vev for the Georgi-Jarlskog field $\langle\Sigma_e\rangle = 3\langle\Sigma_d\rangle$.

Although eq. (5) is the minimal structure (RVV1) present for all possible models, it is possible to build other symmetry-dependent soft-mass structures for particular choices of the global symmetries and charges. The observed structure in the Yukawa couplings does not fix completely the introduced global charges and it is possible to add new invariant combinations of flavon fields to the Kähler potential without modifying the Yukawas. The first example of these new combinations of flavon fields in the Kähler is achieved by allowing a $\theta_3^i\bar{\theta}_{23}^j$ term (RVV2). A second possibility is to allow a $(\epsilon^{ikl}\theta_3^k\theta_{23}^l)\theta_3^j$ term (RVV3) in the Kähler. The required charges for each of these two possibilities can be found in ref. [5]. The structure of the Yukawa couplings in the superpotential remains unchanged. This is due to the fact that the superpotential is a holomorphic function of the fields while the Kähler is only a real function.

The trilinear couplings follow the same symmetries as the Yukawas. Thus, they have the same flavon structure in RVV1, RVV2 and RVV3. Although they have the same structure, they do not have the same $O(1)$ constants, which means that the rotation into the SCKM basis does not diagonalize them.

In the quark sector, the misalignment of the Y_u and Y_d matrices gives sizeable contributions to the LL and LR sectors. In the lepton sector with RH neutrinos, the same happens due to the misalignment of Y_ν and Y_e [1, 9, 10]. The Y_ν contribution is highly model dependent.

Flavour models based on $SU(3)$ give rise to potentially large rates of LFV processes, such that positive signals of LFV can be found in the currently running or near-future experiments, at least for SUSY masses within the reach of the LHC [11]. The presence of large mixing among flavours relies on the features of the above $SU(3)$ model: the presence of nonuniversal scalar masses already at the scale where the SUSY breaking terms appear, and the fact that the trilinear A_f matrices are in general not aligned with the corresponding Yukawa matrices. Let us start considering the case $A_0 = 0$, where

TABLE I. – *Order of magnitude of LFV mass insertions, for the three models.*

	$\|(\delta^e_{LL})_{12}\|$	$\|(\delta^e_{LL})_{13}\|$	$\|(\delta^e_{LL})_{23}\|$	$\|(\delta^e_{RR})_{12}\|$	$\|(\delta^e_{RR})_{13}\|$	$\|(\delta^e_{RR})_{23}\|$
RVV1	$\frac{1}{3}\varepsilon^2\bar{\varepsilon}$	$y_t\bar{\varepsilon}^3$	$3y_t\bar{\varepsilon}^2$	$\frac{1}{3}\bar{\varepsilon}^3$	$\frac{1}{3}\bar{\varepsilon}^3$	$\bar{\varepsilon}^2$
RVV2	$\frac{1}{3}\varepsilon^2\bar{\varepsilon}$	$\frac{1}{3}\sqrt{y_t}\varepsilon\bar{\varepsilon}$	$\sqrt{y_t}\varepsilon$	$\frac{1}{3}\bar{\varepsilon}^3$	$\frac{1}{3}\sqrt{y_b}\bar{\varepsilon}^2$	$\sqrt{y_b}\bar{\varepsilon}$
RVV3	$3y_t\varepsilon\bar{\varepsilon}^2$	$y_t\varepsilon$	$3y_t\bar{\varepsilon}^2$	$\frac{1}{3}\bar{\varepsilon}^3$	$y_b\bar{\varepsilon}$	$\bar{\varepsilon}^2$

the latter effect is strongly reduced so that, in terms of mass insertions, BR$(l_i \to l_j\gamma)$ mainly depends on $|(\delta^e_{LL})_{ij}|^2$ and $|(\delta^e_{RR})_{ij}|^2$. Looking at the structure of the slepton soft mass matrices in the three versions of the model (table I), we see that RVV1 and RVV2 are expected to give similar predictions for BR$(\mu \to e\gamma)$ and BR$(\tau \to \mu\gamma)$. In the case of RVV3, the prediction for BR$(\tau \to \mu\gamma)$ will be also similar to the previous two cases, while we expect BR$(\mu \to e\gamma)$ to be strongly enhanced. For RVV3, the LL mass insertion is larger by a factor $9\, y_t\, \bar{\varepsilon}/\varepsilon = \mathcal{O}(10)$ with respect to RVV1 and RVV2, and the BR$(\mu \to e\gamma)$ is consequently increased by two orders of magnitude.

To summarize, let us compare the expectations for the different LFV processes. In the case $A_0 = 0$, considering for simplicity only the contribution from δ^e_{LL}, we have

$$\frac{\mathrm{BR}(\tau \to e\,\gamma)}{\mathrm{BR}(\mu \to e\,\gamma)} \approx \left(\frac{m_\tau}{m_\mu}\right)^5 \frac{\Gamma_\mu}{\Gamma_\tau}\frac{(\delta^e_{\mathrm{LL}})^2_{13}}{(\delta^e_{\mathrm{LL}})^2_{12}} \approx \mathcal{O}(1)\,(\mathrm{RVV1}, \mathrm{RVV2}, \mathrm{RVV3}), \tag{7}$$

$$\frac{\mathrm{BR}(\tau \to \mu\,\gamma)}{\mathrm{BR}(\mu \to e\,\gamma)} \approx \left(\frac{m_\tau}{m_\mu}\right)^5 \frac{\Gamma_\mu}{\Gamma_\tau}\frac{(\delta^e_{\mathrm{LL}})^2_{23}}{(\delta^e_{\mathrm{LL}})^2_{12}} \approx \mathcal{O}(10^3)\,(\mathrm{RVV1}, 2), \mathcal{O}(10)\,(\mathrm{RVV3}), \tag{8}$$

where Γ_μ (Γ_τ) is the μ (τ) full width. Given the fact that the upper bound on BR$(\tau \to e\,\gamma)$ is 4 orders of magnitude higher than that on BR$(\mu \to e\,\gamma)$, we see that BR$(\tau \to e\gamma)$ is not able to constrain the parameter space better than BR$(\mu \to e\gamma)$ in none of the three models. On the other hand, we expect from eq. (8) that the present constraints given by $\mu \to e\gamma$ and $\tau \to \mu\gamma$, that differ by three orders of magnitude, are comparable for RVV1 and RVV2, while $\mu \to e\gamma$ should give the strongest constraint in the case of RVV3.

In the case $A_0 \neq 0$, generally large δ^e_{LR} insertions arise as a consequence of the misalignment between A_f and the corresponding Yukawa matrix Y_f. In this case, the neutralino contribution to BR$(\mu \to e\gamma)$ gets strongly enhanced [11] and the present (or future) bound requires heavier SUSY masses to be fulfilled, specially in the region where the gaugino mass is much larger than the common sfermion mass. Nevertheless, we expect this effect to be visible only in the case of RVV1 and RVV2, while for RVV3 the very large $(\delta^e_{LL})_{12}$ should still give the dominant contribution.

For the numerical analysis for the LFV decays, we fix the unknown $O(1)$ parameters to random values. The presently allowed region on the m_0-$M_{1/2}$ plane is approximately $(m_0, M_{1/2}) \gtrsim (700, 300)$ GeV. In the case of RVV3, $\mu \to e\gamma$ already gives a strong constraint, $(m_0, M_{1/2}) \gtrsim (1400, 800)$ GeV, which is much more stringent than the one provided by $\tau \to \mu\gamma$. As a consequence, for SUSY masses lying within the LHC reach, RVV3 results already rather disfavoured, while RVV1 and RVV2 are not strongly constrained. Considering the sensitivity expected at the MEG experiment for BR$(\mu \to e\gamma)$, $\mathcal{O}(10^{-13})$, we see that also RVV1 and RVV2 will be tested in most of the parameter space accessible to the LHC, while RVV3 will be completely probed well beyond the LHC reach.

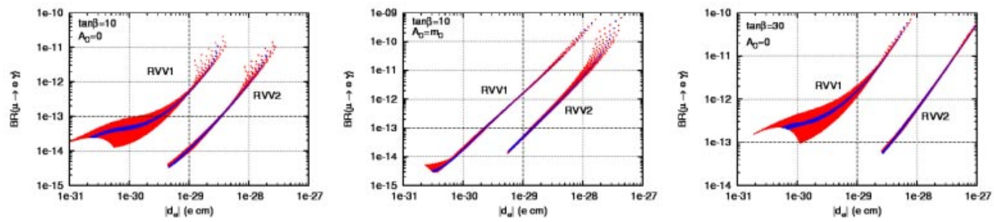

Fig. 1. – (Colour on-line) BR($\mu \to e\gamma$) *vs.* $|d_e|$ for different scenarios. See the text for details.

Moreover, in case of larger values of $\tan\beta$ (*e.g.*, $\tan\beta = 30$), since BR($\mu \to e\gamma$) $\propto \tan^2\beta$, MEG will be able to test all the parameter space accessible to the LHC also for RVV1 and RVV2.

As RVV3 is heavily constrained by LFV, in the following we shall exclude it from our analysis, and concentrate exclusively on RVV1 and RVV2.

The EDMs of fermions, such as the electron and the neutron, are highly suppressed in the SM, and thus they are excellent observables where to look for CP-violation in new physics.

The electron EDM was studied in ref. [11] within the context of RVV1. In these models CP is spontaneously broken in the flavour sector. Therefore, the phases in the μ parameter and diagonal A_f terms are highly suppressed and can be neglected. In such a case, the imaginary parts required for EDMs only appear from flavour-changing mass insertions.

Electron EDM predictions are large enough to be probed at future EDM experiments. For relatively light SUSY masses we obtain $d_e \sim 10^{-29}\,e\,\text{cm}^{-1}$ and $d_e \sim 10^{-28}\,e\,\text{cm}^{-1}$, for RVV1 and RVV2, respectively. The latter predicts a value of d_e about one order of magnitude larger than the former for any particular value of m_0 and $M_{1/2}$ due to the larger ε suppresion. This means that by reaching $d_e \sim 10^{-29}\,e\,\text{cm}^{-1}$ one could probe a much larger part of the evaluated parameter space, with $m_0 \lesssim 1500\,\text{GeV}$, $M_{1/2} \lesssim 2000\,\text{GeV}$. In particular, for RVV2, observation of SUSY at the LHC and solving the ϵ_K tension [12] would force d_e to be larger than $10^{-29}\,e\,\text{cm}^{-1}$. However, we have to take into account that these values will vary by factors $O(1)$ because of the unknown $O(1)$ coefficients to the different MIs.

If we require in addition that the $(g-2)_\mu$ discrepancy between SM and data is explained by SUSY, we are restricted in a region of rather light SUSY masses, where most of the observables are expected to be close to the present experimental bounds.

In fig. 1, we compare the discovery potential of the two most promising leptonic observables, $\mu \to e\gamma$ and the electron EDM. The correlation of BR($\mu \to e\gamma$) *vs.* $|d_e|$ is plotted for both RVV1 and RVV2, in the case $\tan\beta = 10$, $A_0 = 0$ (left), $\tan\beta = 10$, $A_0 = 1$ (center) and $\tan\beta = 30$, $A_0 = 0$ (right). We study the mass range: $0 < m_0 < 2.5\,\text{TeV}$, $0 < M_{1/2} < 1.5\,\text{TeV}$. In the figures, only the "ϵ_K-favoured" region with negative $(\delta^d_{RR})_{12}$ has been plotted with blue and red colours corresponding to two different implementations of the constraint of having SUSY contributions to account for the "SM-deficit" in reproducing the correct ϵ_K value [12]. The horizontal line corresponds to the final sensitivity of MEG, the vertical line to the sensitivity on $|d_e|$ of the running Yale-PdO experiment. We see that, for RVV1, $\mu \to e\gamma$ should be able to constrain the parameter space more strongly than eEDM, while for RVV2 it is $|d_e|$ the most sensitive observable (except for the large $\tan\beta$ case). These features could be useful in the future, in order to discriminate among different models and, more in general, shed light on the structure of mixings and phases in the slepton mass matrix.

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