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Higgs-mediated $e \to \mu$ transitions in II Higgs doublet model and supersymmetry

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Abstract: We study the phenomenology of the $e-\mu$ lepton flavor violation (LFV) in a general two Higgs Doublet Model (2HDM) including the supersymmetric case. We compute the decay rate expressions of $\mu \to e\gamma$, $\mu \to eee$, and $\mu \to e$ conversion in nuclei at two loop level. In particular, it is shown that $\mu \to e\gamma$ is generally the most sensitive channel to probe Higgs-mediated LFV. The correlations among the decay rates of the above processes are also discussed.

Keywords: Rare Decays, Supersymmetry Phenomenology.

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Contents

1. Introduction

The Supersymmetric (SUSY) extension of the Standard Model (SM) is one of the most promising candidate for physics beyond the SM. Besides direct searches for SUSY particles, it is also important to analyze implications of such a theory in the low-energy phenomena, through virtual effects of SUSY particles. Lepton flavor violating (LFV) processes are excellent candidates to explore such virtual SUSY effects from low energy experiments [1]. In fact, in a SM framework with massive neutrinos, all the LFV transitions are expected at a very suppressed level, very far from any future and reasonable experimental resolutions. On the other hand, the observation of neutrino oscillation have established the existence of lepton family number violation.

This mixing is expected to be manifested also in the charged lepton sector through the observation of rare decay processes such as $\mu \to e\gamma$, $\tau \to \mu\gamma$ etc. Any experimental signals of such a process would clearly indicate the presence of a non standard mechanism. In a supersymmetric (SUSY) framework, new direct sources of flavor violation appear, provided the presence of off-diagonal soft terms in the slepton mass matrices and in the trilinear couplings [2]. In practice, flavor violation would originate from any misalignment between fermion and sfermion mass eigenstates. LFV processes arise at one loop level through the exchange of neutralinos (charginos) and charged sleptons (sneutrinos). The amount of the LFV is regulated by a Super-GIM mechanism that can be much less severe than in the non supersymmetric case [3–5].[1]

Another potential source of LFV in models such as the minimal supersymmetric standard model (MSSM) is the Higgs sector.

[1] As recently shown in ref. [6], some of these effects are common to many extensions of the SM, even to non-susy scenarios, and can be described in a general way in terms of an effective field theory. Moreover, in the context of general $SU(2)_L \times U(1)_Y$ seesaw scenarios, large LFV effects can be induced by the exchange of left-handed and/or right-handed neutral singlets [7].

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In fact, extensions of the Standard Model containing more than one Higgs doublet generally allow flavor-violating couplings of the neutral Higgs bosons with fermions. Such couplings, if unsuppressed, will lead to large flavor-changing neutral currents in direct opposition to experiments.

The MSSM avoids these dangerous couplings at the tree level segregating the quark and Higgs fields so that one Higgs (H_u) can couple only to up-type quarks while the other (H_d) couples only to d-type. Within unbroken supersymmetry this division is completely natural, in fact, it is required by the holomorphy of the superpotential. However, after supersymmetry is broken, couplings of the form QU_cH_d and QD_cH_u are generated at one loop [8].

As shown in ref. [9, 10], the presence of these loop-induced non holomorphic couplings also leads to the appearance of flavor-changing couplings of the neutral Higgs bosons that are particularly relevant at large values of $\tan\beta$. As a natural consequence, a variety of flavor-changing processes such as $B^0 \to \mu^+\mu^-$ [10], $\bar{B^0} - B^0$ [11], $K \to \pi\nu\bar{\nu}$ [13] etc. is generated.[2] Higgs-mediated FCNC can have sizable effects also in the lepton sector [14]: given a source of nonholomorphic couplings, and LFV among the sleptons, Higgs-mediated LFV is unavoidable.

These effects have been widely discussed in the recent literature both in a generic 2HDM [15 − 17] and in supersymmetry [17 − 21] frameworks. In particular, it has been shown that a tree level Higgs-exchange leads to $\tau \to l_j l_k l_k$ [14], $\tau \to l_j\eta$ [18], $B^0 \to l_j\tau$ [20] and $\mu \to e$ conversion in Nuclei [21].

Recently, it was pointed out that Higgs-mediated LFV effects can also generate violations of lepton universality at the 1% level in the $R = \Gamma(K \to e\nu)/\Gamma(K \to \mu\nu)$ ratio [22].

Moreover, Higgs-mediated FCNC can have a sizable impact also in loop-induced processes, such as $\tau \to l_j\gamma$ [17].

In this letter, we investigate the effects of Higgs mediated LFV in the $e - \mu$ transitions both in a generic two Higgs Doublet Model (2HDM) and in Supersymmetry. We evaluate analytical expressions and correlations for the rates of $\mu \to e\gamma$, $\mu \to eee$ and $\mu \to e$ conversion in nuclei at two loop level, establishing which the most promising channels to detect LFV signals are.

2. LFV in the Higgs sector

SM extensions containing more than one Higgs doublet generally allow flavor-violating couplings of the neutral Higgs bosons with fermions. Such couplings, if unsuppressed, will lead to large flavor-changing neutral currents in direct opposition to experiments. The possible solution to this problem involves an assumption about the Yukawa structure of the model. A discrete symmetry can be invoked to allow a given fermion type to couple to a single Higgs doublet, and in such case FCNC's are absent at tree level. In particular, when a single Higgs field gives masses to both types of fermions the resulting model is referred as 2HDM-I. On the other hand, when each type of fermion couples to a different Higgs doublet the model is said 2HDM-II.

[2]For a recent and detailed analysis of the B physics phenomenology at large $\tan\beta$ within the Minimal Flavor Violating framework, see [12].

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In the following, we will assume a scenario where the type-II 2HDM structure is not protected by any symmetry and is broken by loop effects (this occurs, for instance, in the MSSM).

Let us consider the Yukawa interactions for charged leptons, including the radiatively induced LFV terms:

$$-\mathcal{L} \simeq \bar{l}_{\mathrm{Ri}} Y_{l_i} H_1 \overline{L_i} + \bar{l}_{\mathrm{Ri}} \left(Y_{l_i} \Delta_L^{ij} + Y_{l_j} \Delta_R^{ij} \right) H_2 \overline{L_j} + h.c. \tag{2.1}$$

where H_1 and H_2 are the scalar doublets, l_{Ri} are lepton singlet for right handed fermions, L_k denote the lepton doublets and Y_{l_k} are the Yukawa couplings.

In the mass-eigenstate basis for both leptons and Higgs bosons, the effective flavor-violating interactions are described by the four dimension operators:

$$\begin{aligned} -\mathcal{L} \simeq\ & (2G_F^2)^{\frac{1}{4}} \frac{m_{l_i}}{c_\beta^2} \left(\Delta_L^{ij} \bar{l}_R^i l_L^j + \Delta_R^{ij} \bar{l}_L^i l_R^j \right) \left(c_{\beta-\alpha} h^0 - s_{\beta-\alpha} H^0 - i A^0 \right) \\ & + (8G_F^2)^{\frac{1}{4}} \frac{m_{l_i}}{c_\beta^2} \left(\Delta_L^{ij} \bar{l}_R^i \nu_L^j + \Delta_R^{ij} \nu_L^i \bar{l}_R^j \right) H^{\pm} + h.c. \end{aligned} \tag{2.2}$$

where α is the mixing angle between the CP-even Higgs bosons h_0 and H_0, A_0 is the physical CP-odd boson, $H^{\pm}$ are the physical charged Higgs-bosons and t_β is the ratio of the vacuum expectation value for the two Higgs (where we adopt the notation, $c_x, s_x = \cos x, \sin x$ and $t_x = \tan x$). Irrespective to the mechanism of the high energy theories generating the LFV, we treat the $\Delta_{L,R}^{ij}$ terms in a model independent way. In order to constrain the $\Delta_{L,R}^{ij}$ parameters, we impose that their contributions to LFV processes do not exceed the experimental bounds [17].

On the other hand, there are several models with a specific ansatz about the flavor-changing couplings. For instance, the famous multi-Higgs-doublet models proposed by Cheng and Sher [23] predict that the LFV couplings of all the neutral Higgs bosons with the fermions have the form $H f_i f_j \sim \sqrt{m_i m_j}$.

In Supersymmetry, the Δ^{ij} terms are induced at one loop level by the exchange of gauginos and sleptons, provided a source of slepton mixing. In the so called MI approximation, the expressions of $\Delta_{L,R}^{ij}$ are given by

$$\begin{aligned} \Delta_L^{ij} =& - \frac{\alpha_1}{4\pi} \mu M_1 \delta_{LL}^{ij} m_L^2 \left[I'(M_1^2, m_R^2, m_L^2) + \frac{1}{2} I'(M_1^2, \mu^2, m_L^2) \right] + \\ & + \frac{3}{2} \frac{\alpha_2}{4\pi} \mu M_2 \delta_{LL}^{ij} m_L^2 I'(M_2^2, \mu^2, m_L^2) \,, \end{aligned} \tag{2.3}$$

$$\Delta_R^{ij} = \frac{\alpha_1}{4\pi} \mu M_1 m_R^2 \delta_{RR}^{ij} \left[I'(M_1^2, \mu^2, m_R^2) - (\mu \leftrightarrow m_L) \right] \tag{2.4}$$

respectively, where μ is the the Higgs mixing parameter, $M_{1,2}$ are the gaugino masses and $m_{L(R)}^2$ stands for the left-left (right-right) slepton mass matrix entry. The LFV mass insertions (MIs), i.e. $\delta_{XX}^{3\ell} = (\tilde{m}_\ell^2)_{XX}^{3\ell} / m_X^2$ $(X = L, R)$, are the off-diagonal flavor changing entries of the slepton mass matrix. The loop function $I'(x, y, z)$ is such that $I'(x, y, z) = dI(x, y, z)/dz$, where $I(x, y, z)$ refers to the standard three point one-loop integral which has mass dimension -2

$$I_3(x,y,z) = \frac{xy\log(x/y) + yz\log(y/z) + zx\log(z/x)}{(x-y)(z-y)(z-x)}\,. \tag{2.5}$$

The above expressions, i.e. the eqs. (2.3), (2.4), depend only on the ratio of the susy mass scales and they do not decouple for large $m_{\rm SUSY}$. As first shown in ref. [19], both Δ_R^{ij} and Δ_L^{ij} couplings suffer from strong cancellations in certain regions of the parameter space due to destructive interferences among various contributions. For instance, from eq. (2.4) it is clear that, in the Δ_R^{ij} case, such cancellations happen if $\mu = m_L$.

In the SUSY seesaw model, the MIs of the slepton mass matrix appear in the left-handed sleptons through the neutrino Yukawa interactions. The superpotential of the lepton sector is given by $W = Y_e H_1 l_R^c L + Y_\nu H_2 N^c L + (1/2) M_N N^c N^c$, where N^c is the superfields corresponding to the right-handed neutrinos. The neutrino mass matrix is obtained by integrating out the heavy right-handed neutrinos as $m_\nu = (Y_\nu^T M_N^{-1} Y_\nu) v^2 \sin^2\beta/2$, where v is the vacuum expectation value (VEV) of the Higgs field ($v = 246\,\mathrm{GeV}$). The correct size of the neutrino masses is obtained for right-handed neutrinos as heavy as $10^{14}\,\mathrm{GeV}$ for $f_\nu \sim O(1)$. The Yukawa coupling Y_ν violates the lepton flavor conservation and this violation is communicated to the slepton mass matrix at low-energy. The renormalization group equation (RGE) running effect induces the following off-diagonal components in the left-handed slepton mass matrix

$$(\tilde{m}^2_{\ell_L})_{ij} \simeq -\frac{1}{8\pi^2} m_0^2 (3 + a^2) \left(Y_\nu^\dagger \log\frac{M_{\rm GUT}}{M_N} Y_\nu \right)_{ij}, \tag{2.6}$$

where the SUSY breaking parameters m_0 and a stand for the scalar mass and the trilinear scalar coupling at the GUT scale, respectively.

Given our ignorance about the mixings in Y_ν, we consider two extremal benchmark cases as discussed, within the SO(10) framework, in [5]. As a minimal mixing case we take the one in which the neutrino and the up-quark Yukawa unify at the high scale, so that the mixing is given by the CKM matrix; this case is named 'CKM-case'. As a maximal mixing scenario we take the one in which the observed neutrino mixing is coming entirely from the neutrino Yukawa matrix, so that $Y_\nu = U_{PMNS} \cdot Y_u^{\rm diag}$, where U_{PMNS} is the Pontecorvo-Maki-Nakagawa-Sakata matrix; in this case the unknown U_{e3} PMNS matrix element turns out to be crucial in evaluating the size of LFV effects. The maximal case is named 'PMNS-case'. As regards the δ_{LL}^{21} MI one obtains [5]

$$\delta_{LL}^{21} = -\frac{3}{8\pi^2}\, Y_t^2 V_{td} V_{ts} \ln\frac{M_X}{M_{R_3}} \qquad \mathrm{CKM-case} \tag{2.7}$$

$$\delta_{LL}^{21} = -\frac{3}{8\pi^2}\, Y_t^2 U_{e3} U_{\mu 3} \ln\frac{M_X}{M_{R_3}} \qquad \mathrm{PMNS-case}\,. \tag{2.8}$$

So, in the CKM-case, it turns out that $\delta_{LL}^{21} \simeq 3\cdot 10^{-5}$ while in the PMNS-case, taking $U_{e3} = 0.07$ at about half of the current CHOOZ bound, we get $\delta_{LL}^{21} \simeq 10^{-2}$.

3. $e-\mu$ transitions in the non-decoupling limit

In this section, we will analyze $e-\mu$ transitions through the study of $\mu\to e\gamma$, $\mu\to eee$ and $\mu\to e$ conversion in nuclei in the non-decoupling limit of a 2HDM, where $s_{\beta-\alpha}=0$ and t_β is large. In particular, we derive analytical expressions and correlations for the examined branching ratios in order to establish which the most promising channels to detect Higgs mediated LFV are. $\mu\to e\gamma$ is generated by a dipole operator arising, at least, from one loop Higgs exchange. However, higgs mediated dipole transitions imply three chirality flips: two in the Yukawa vertices and one in the lepton propagator. This strong suppression can be overcome at higher order level. Going to two loop level, one has to pay the typical $g^2/16\pi^2$ price but one can replace light fermion masses from yukawa vertices with heavy fermion (boson) masses circulating in the second loop [15, 24]. In this case, the virtual higgs boson couples only once to the lepton line, inducing the needed chirality flip. As a result, two loop amplitudes provide the major effects and we find that $Br(\mu\to e\gamma)$ is given by

$$Br(\mu\to e\gamma)\simeq\frac{3}{8}\frac{\alpha_{\rm el}}{\pi}\frac{m_\mu^4}{M_{h,A}^4}\Delta_{21}^2t_\beta^6\Bigg[\pm\log\frac{m_\mu^2}{M_{h,A}^2}-\frac{2\alpha_{\rm el}}{\pi}\left(\frac{m_W^2}{m_\mu^2}\right)\frac{F(a_W)}{t_\beta}+ \pm\frac{\alpha_{\rm el}}{\pi}\sum_{f=b,\tau}N_fq_f^2\left(\frac{m_f^2}{m_\mu^2}\right)\left(\log\frac{m_f^2}{M_{h,A}^2}\right)^2\Bigg]^2 \tag{3.1}$$

where $N_{\tau,b}=1,3$, q_f is the electric charge of the fermion f and $a_W=m_W^2/m_h^2$. The term proportional to $F(a_W)$ arises from two loop effects induced by Barr-Zee type diagrams [24] with a W boson exchange. The loop function $F(z)$ is given by

$$F(z)\simeq 3f(z)+\frac{23}{4}g(z)+\frac{f(z)-g(z)}{2z} \tag{3.2}$$

with the Barr-Zee loop integrals given by:

$$g(z)=\frac{1}{4}\int_0^1dx\frac{\log\left(z/x(1-x)\right)}{z-x(1-x)}\,, \tag{3.3}$$

$$f(z)=\frac{1}{4}\int_0^1dx\frac{1-2x(1-x)\log\left(z/x(1-x)\right)}{z-x(1-x)}. \tag{3.4}$$

For $z\ll 1$ it turns out that:

$$F(z)\sim\frac{35}{16}(\log z)^2+\frac{\log z+2}{4z}. \tag{3.5}$$

The first term of eq. (3.1) refers to one loop contributions while the last term arises from two loop effects induced by fermionic loops. In the computation, we retained only the $h^0\gamma\gamma$ effective vertex neglecting the $(1-4\sin^2_W)$ suppressed contributions arising from the $h^0Z\gamma$ vertex.

To get a feeling on the relative size among different contributions, we note that two loop fermionic (bosonic) amplitudes are enhanced by an m_f^2/m_μ^2 ($m_W^2\cot\beta/m_\mu^2$) factor with respect to the one loop amplitude. In fact, one gets heavy fermionic (bosonic) masses both from the fermionic (bosonic) propagators and from the $H\bar{f}f \sim m_f t_\beta$ ($HWW \sim m_W$) couplings. Two loop effects generated by the top quark are generally subdominant. In fact, bearing in mind that any Model II 2HDM predicts that $H\bar{t}t \sim m_t/t_\beta$ and noting that the top amplitude isn't enhanced by large logarithm factors one finds naively that

$$\frac{A^b}{A^t} \sim \frac{q_b^2}{q_t^2}\frac{m_b^2 t_\beta^2}{m_{\rm top}^2}\left(\log\frac{m_b^2}{m_h^2}\right)^2,$$

where $A_{t,b}$ stands for the top and bottom two loop amplitudes. Since the Higgs mediated LFV is relevant only at large $t_\beta \geq 30$, it is clear that τ and b contributions are dominant.

Moreover, from eqs. (3.1)–(3.5) it is straightforward to check that two loop effects are largely dominated by the W exchange instead of the exchange of heavy fermions. A possible exception arises only if $m_A \ll m_h$. In fact, bearing in mind that pseudoscalar bosons do not couple to a W pair, it turns out that two loop W effects are sensitive only to scalar mediation, in contrast to the fermionic case. At this point, we proceed to consider the contributions to $\mu \to eee$ and $\mu Al \to eAl$. We find that $\mu \to eee$ is completely dominated by the photonic $\mu \to e\gamma^*$ dipole amplitude so that $Br(\mu \to eee) \simeq \alpha_{em} Br(\mu \to e\gamma)$. On the other hand, $\mu \to e$ conversion in Nuclei gets the major effects by the scalar operator through the tree level Higgs exchange that leads to the following expression for $Br(\mu Al \to eAl)$:

$$Br(\mu Al \to eAl) \simeq 1.8 \times 10^{-4}\,\frac{m_\mu^7 m_p^2}{v^4 m_h^4 \omega_{\rm capt}^{\rm Al}}\,\Delta_{21}^2 t_\beta^6\,, \tag{3.6}$$

where $\omega_{\rm capt}^{\rm Al} \simeq 0.7054 \cdot 10^6 sec^{-1}$. In fact, in contrast to $\mu \to 3e$, that is suppressed by the electron mass through the $H(A)\bar{e}e \sim m_e$ coupling, $\mu N \to eN$ is not suppressed by the light constituent quark m_u and m_d but only by the nucleon masses, because the Higgs-boson coupling to the nucleon is shown to be characterized by the nucleon mass using the conformal anomaly relation [25]. In particular, the most important contribution turns out to come from the exchange of the scalar Higgs boson h and H which couples to the strange quark [26].[3] Moreover, from a previous analysis [17], we know that $\mu \to e\gamma^*$ (chirality conserving) monopole amplitudes are generally subdominant compared to (chirality flipping) dipole effects. In addition, the enhancement mechanism induced by Barr-Zee type diagrams is effective only for chirality flipping operators so, in the following, we will disregard chirality conserving one loop effects. Let us derive now the approximate relations among $\mu Al \to eAl$, $\mu \to e\gamma$ and $\mu \to eee$ branching ratios

$$\frac{Br(\mu \to e\gamma)}{Br(\mu Al \to eAl)} \simeq 10^2\left(\frac{F(a_W)}{\tan\beta}\right)^2, \quad \frac{Br(\mu \to eee)}{Br(\mu \to e\gamma)} \simeq \alpha_{\rm el} \tag{3.7}$$

[3] As discussed in [21], the coherent $\mu - e$ conversion process, where the initial and final nuclei are in the ground state, is expected to be enhanced by a factor of $O(Z)$ (where Z is the atomic number) compared to incoherent transition processes. Since the initial and final states are the same, the elements $\langle N|\bar{p}p|N\rangle$ and $\langle N|\bar{n}n|N\rangle$ are nothing but the proton and the neutron densities in a nucleus in the non-relativistic limit of nucleons. In this limit, the other matrix elements $\langle N|\bar{p}\gamma_5 p|N\rangle$ and $\langle N|\bar{n}\gamma_5 n|N\rangle$ vanish. Therefore, in the coherent $\mu - e$ conversion process, the dominant contributions come from the exchange of h and H, not A.

where the approximate expression for $F(a_W)$ is given by eq. (3.5). In the above equations we retained only dominant two loop effects arising from W exchange. The exact behavior for the examined processes is reported in figure 1 where we can see that $\mu \to e\gamma$ gets the largest branching ratio except for a region around $m_H \sim 700$Gev where strong cancellations among two loop effects sink its size. For a detailed discussion about the origin of these cancellations and their connection with non-decoupling properties of two loop W amplitude, see ref. [15]. On the other hand, $\mu \to e\gamma$ amplitude can receive large one loop contributions by a double LFV source, namely by $(\Delta^{21})_{eff.} = \Delta^{23}\Delta^{31}$ and therefore, the resulting $Br(\mu \to e\gamma)$ is:

$$
\begin{aligned}
Br(\mu \to e\gamma) \simeq \frac{3}{8}\frac{\alpha_{\rm el}}{\pi}\left(\frac{m_\tau^2}{M_{h,A}^2}\right)^2 t_\beta^8 \Bigg[& \left(\pm \log\frac{m_\tau^2}{M_{h,A}^2} + \frac{4(-5)}{3}\right)\Delta_L^{23}\Delta_L^{31} + \\
& \pm\left(\frac{m_\tau}{m_\mu}\right)\left(\log\frac{m_\tau^2}{M_{h,A}^2} + \frac{3}{2}\right)\Delta_R^{23}\Delta_L^{31}\Bigg]^2 + (L \leftrightarrow R).
\end{aligned} \tag{3.8}
$$

Now, in contrast to one loop contributions with a single LFV coupling (see the first term of eq. (3.1)), it is always possible to pick up m_τ instead of m_μ both at the LFV Yukawa vertices and at the fermion propagator. However, if the LFV couplings are generated radiatively (as it happens for instance in a Susy framework), the above enhancement is modulated by the loop suppression. In practice, the dominance of one loop effects (with two LFV couplings) over two loop effects (with one LFV coupling) depends on the specific model we are treating, namely on the size of Δ_{ij} terms. Assuming that the contributions with a double source of LFV (see eq. (3.8)) dominate over those with a single LFV source (see eq. (3.1)), the following ratios are expected:

$$
\frac{Br(\mu Al \to eAl)}{Br(\mu \to e\gamma)} \simeq \frac{Br(\mu \to eee)}{Br(\mu \to e\gamma)} \simeq \alpha_{\rm el}. \tag{3.9}
$$

The $\mu \to e\gamma^*$ dominance in $Br(\mu Al \to eAl)$ and $Br(\mu \to eee)$ is the reason of the above correlations. On the other hand, the same correlations are expected, for instance, in a Susy framework with gaugino mediated LFV and then, the predictions of eq. (3.9) prevent us from distinguishing between the two scenarios. Possible deviations from eq. (3.9) can arise only through tree level Higgs exchange effects to $\mu Al \to eAl$.

4. $e - \mu$ transitions in the decoupling limit

The decoupling limit of a 2HDM is a particularly appealing scenario in that it is achieved by Supersymmetry. In this context, the higgs bosons masses are nearly degenerate $m_A \simeq m_H \simeq m_{H^\pm}$ being the mass splitting of order $\mathcal{O}(m_Z^2/m_A)$ and, in addition, it turns out that $c_{\beta-\alpha} = 0$ and $m_Z/m_A \to 0$. In particular, the couplings of the light Higgs boson h are nearly equal to those of the SM Higgs boson. In a Supersymmetric framework, besides the higgs mediated LFV transitions, we have also LFV effects mediated by the gauginos through loops of neutralinos (charginos)- charged sleptons (sneutrinos). On the other hand, the above contributions have different decoupling properties regulated by the mass of the heaviest scalar mass (m_H) or by the heaviest mass in the slepton gaugino loops ($m_{\rm SUSY}$).

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However, in both cases, the effective operator for $l_i \to l_j \gamma$ is

$$\frac{m_{l_i}}{m^2_{H,\mathrm{SUSY}}}\;\bar{l}_i \sigma^{\mu\nu} l_j F_{\mu\nu}.$$

In principle, the m_{SUSY} and m_H masses can be unrelated so, we can always proceed by considering only the Higgs mediated effects (assuming a relatively light m_H and an heavy m_{SUSY}) or only the gaugino mediated contributions (if m_H is heavy). So, taking into account the only Higgs-mediated effects, we get the following branching ratio for $\mu \to e\gamma$:[4]

$$\begin{aligned} Br(\mu \to e\gamma) &\simeq \frac{3}{2}\frac{\alpha^3_{\mathrm{el}}}{\pi^3}\Delta^2_{21} t^6_\beta \Bigg[\sum_{f=b,\tau} N_f q^2_f \frac{m^2_f}{M^2_H}\left(\log\frac{m^2_f}{M^2_H} + 2\right) - \frac{m^2_W}{M^2_H}\frac{F(a_W)}{t_\beta} + \\ &\quad + \frac{N_c}{4}\left(q^2_{\tilde t}\,\frac{m_t \mu}{t_\beta M^2_H}\,\sin 2\theta_{\tilde t}\, h(x_{\tilde t H}) - q^2_{\tilde b}\,\frac{m_b A_b}{M^2_H}\,\sin 2\theta_{\tilde b}\, h(x_{\tilde b H})\right)\Bigg]^2 \\ &\simeq \frac{3}{2}\frac{\alpha^3_{\mathrm{el}}}{\pi^3}\,\Delta^2_{21}\, t^4_\beta \left(\frac{m^4_W}{M^4_H}\right)\Big(F(a_W)\Big)^2 \end{aligned} \tag{4.1}$$

where $a_W = m^2_W/m^2_H$, $x_{\tilde f H} = m^2_{\tilde f}/m^2_H$, $\theta_{\tilde t, \tilde b}$ are squarks mixing angles and the loop function $h(z)$ is given by:

$$h(z) = \int_0^1 dx \frac{x(1-x)\log\left(z/x(1-x)\right)}{z - x(1-x)}. \tag{4.2}$$

The asymptotic form of $h(z)$, which may be useful for an easy understanding of the results, is given by:

$$h(z) = \begin{cases} -\left(\log z + 2\right) & z \ll 1 \\ 0.344 & z = 1 \\ \frac{1}{6z}(\log z + \frac{5}{3}) & z \gg 1. \end{cases} \tag{4.3}$$

The first two terms of eq. (4.1) refer to two loop effects induced by fermionic and W loops, respectively, while the last term appears only in a supersymmetric framework and it is relative to squark loops [27]. In the previous section, we have seen that W effects dominate over the fermionic ones. Moreover, being the H and A masses almost degenerate in the decoupling limit, the H and A contributions partially cancel themselves in the fermionic amplitude because of their opposite signs. This is in contrast to the W amplitude that turns out to be sensitive only to H effects.

As regards the squark loop effects, it is very easy to realize that they are negligible compared to W effects. In fact, it is well known that Higgs mediated LFV can play a relevant or even a dominant role compared to gaugino mediated LFV provided that slepton masses are not below the TeV scale while maintaining the Higgs masses at the electroweak scale (and assuming large t_β values). In this context, it is natural to assume squark masses

[4]In a SUSY framework, the couplings between the scalar and the fermions are given by $-i(\sqrt{2}G_F)^{1/2}\tan\beta H \xi_f m_f \overline{f} f$ where the parameters ξ_f are equal to one at tree level but they can get large corrections from higher order effects. For instance, ξ_b gets contributions from gluino-squark loops (proportional to $\alpha_s t_\beta$) that enhance or suppress significantly the tree level value of ξ_b [8]. In the ξ_τ case the leading one loop effects induced by chargino-sneutrino contributions (proportional to $\alpha_w t_\beta$) do not affect ξ_τ so significantly. For simplicity's sake, we disregard the above factors in the following.

at least of the same order as the slepton masses (at the TeV scale) so that $x_{\tilde{f}H} \gg 1$. So, even for maximum squark mixings, namely for $\sin 2\theta_{\tilde{t},\tilde{b}} \simeq 1$, and large A_b and μ terms, two loop squark effects remain much below the W effects, as it is straightforward to check by eqs. (4.1), (4.2), (4.3). In the decoupling limit, $\mu \to eee$ and $\mu Al \to eAl$ are still dominated by the two loop $\mu \to e\gamma^*$ amplitude and by a three level Higgs exchange, respectively. Finally one gets the following relations:

$$\frac{Br(\mu \to e\gamma)}{Br(\mu Al \to eAl)}\bigg|_{\rm Higgs} \simeq 10^2 \left(\frac{F(a_W)}{\tan\beta}\right)^2 \; , \quad \frac{Br(\mu \to eee)}{Br(\mu \to e\gamma)}\bigg|_{\rm Higgs} \simeq \alpha_{\rm el}. \tag{4.4}$$

As we can note, the above predictions are exactly the same ones we found in the non-decoupling limit and the numerical results are reported in figure 1. However, this property is no longer true for processes associated to $e-\tau$ and $\mu-\tau$ transitions [17]. Let us now consider the one loop contributions to $\mu - e$ transitions arising from $(\Delta^{21})_{eff.} = \Delta^{23}\Delta^{31}$. The corresponding $Br(\mu \to e\gamma)$ is:

$$\begin{aligned} Br(\mu \to e\gamma) \simeq \frac{3}{2}\frac{\alpha_{\rm el}}{\pi}\left(\frac{m_\tau^2}{m_A^2}\right)^2 t_\beta^8 \Bigg[\bigg(\frac{\delta m}{m_A}\log\frac{m_\tau^2}{m_A^2}+\frac{1}{6}\bigg)\Delta_L^{23}\Delta_L^{31} + \\ +\left(\frac{m_\tau}{m_\mu}\right)\left(\frac{\delta m}{m_A}\right)\left(\log\frac{m_\tau^2}{m_A^2}+\frac{3}{2}\right)\Delta_R^{23}\Delta_L^{31}\Bigg]^2 + (L \leftrightarrow R) \end{aligned} \tag{4.5}$$

where $\delta m = m_A - m_H$. The proportional term to $\Delta_R^{23}\Delta_L^{31} \sim \delta_{RR}^{23}\delta_{LL}^{31}$ in eq.16 is enhanced by an m_τ/m_μ factor compared to the proportional term to $\Delta_L^{23}\Delta_L^{31} \sim \delta_{LL}^{23}\delta_{LL}^{31}$. On the other hand, this enhancement is not effective in a Susy framework. In fact, the upper bounds on $\delta_{RR}^{23}\delta_{LL}^{31}$ imposed by the gaugino mediated effects to $Br(\mu \to e\gamma)$ are stronger than those relative to $\delta_{LL}^{23}\delta_{LL}^{31}$ of the same m_τ/m_μ factor [28], as we will discuss. Differently from the non-decoupling limit case, now one loop effects with two LFV couplings are suppressed by the mass splitting $\delta m/m_A$. In a SUSY framework, if $\delta m/m_A \simeq 10\%$, $\Delta^{21} \sim 10^{-3}\delta^{21}$ and $\delta^{21} \sim \delta^{23}\delta^{31}$ we get $Br^{\mu\to e\gamma}_{1-loop}$, roughly two-three order of magnitude below the $Br^{\mu\to e\gamma}_{2-loop}$ obtained from two loop effects with a single LFV coupling. However, in a generic model II 2HDM, one loop effects can still provide the major effects depending on the size of the Δ_{ij} terms. In the following, we are interested to make a comparison between Higgs and gaugino mediated LFV effects. To this purpose let us report the branching ratio of $l_i \to l_j\gamma$ induced by the one loop exchange of neutralinos, charginos and sleptons

$$\frac{BR(\mu \to e\gamma)}{BR(\mu \to e\nu_\mu\bar{\nu}_e)} = \frac{48\pi^3\alpha}{G_F^2}(|A_L|^2 + |A_R|^2)\,,$$

where the $A_{L,R}$ amplitudes are given by

$$\begin{aligned} A_L = \frac{\alpha_2}{4\pi}\delta_{LL}^{21} t_\beta \; \Bigg[\; & \mu M_2 \frac{(f_{2n}(a_{2L},b_L)+f_{2c}(a_{2L},b_L))}{m_L^4(M_2^2-\mu^2)} \\ & + \tan^2\theta_W \mu M_1 \left(\frac{-f_{2n}(a_{1L},b_L)}{m_L^4(M_1^2-\mu^2)} + \frac{1}{(m_R^2-m_L^2)}\cdot \right. \\ & \left.\left(\frac{2f_{2n}(a_{1L})}{m_L^4}+\frac{1}{(m_R^2-m_L^2)}\left(\frac{f_{3n}(a_{1R})}{m_R^2}-\frac{f_{3n}(a_{1L})}{m_L^2}\right)\right)\right)\Bigg], \end{aligned} \tag{4.6}$$

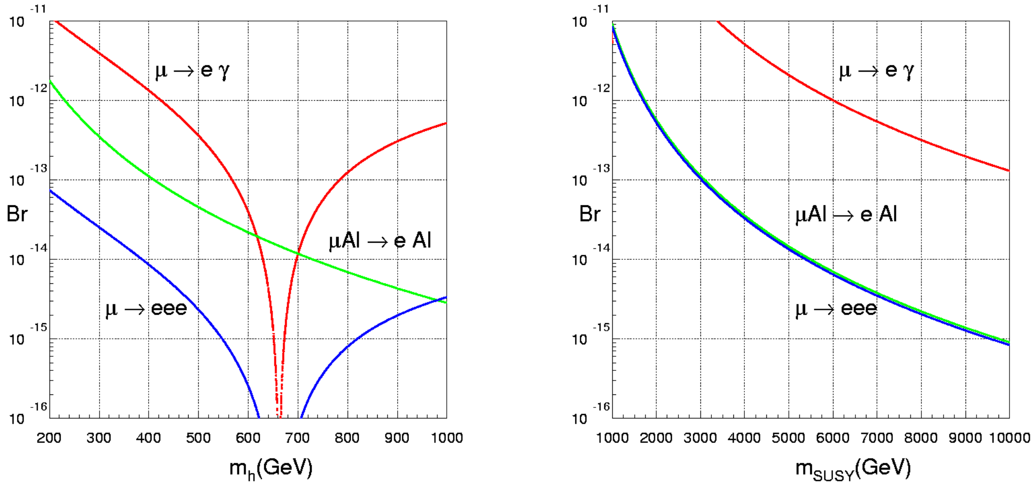

Figure 1: Left: Branching ratios of $\mu \to e\gamma$, $\mu \to eee$ and $\mu Al \to eAl$ in the Higgs mediated LFV case vs the Higgs boson mass m_h. In the decoupling (non-decoupling) limit m_h refers to the heaviest (lightest) Higgs boson mass. Right: Branching ratios of $\mu \to e\gamma$, $\mu \to eee$ and $\mu Al \to eAl$ in the gaugino mediated LFV case vs a common SUSY mass $m_{\rm SUSY}$. In both of figures we set $t_\beta = 50$ and $\delta^{21}_{LL} = 10^{-2}$ (that corresponds, in a generic 2HDM, to $\Delta^{21}_L \simeq 5 \cdot 10^{-6}$).

$$A_R = \frac{\alpha_1}{4\pi}\delta^{21}_{RR} t_\beta \mu M_1 \left[\frac{2 f_{2n}(a_{1R}, b_R)}{m_R^4 (M_1^2 - \mu^2)} + \frac{1}{(m_L^2 - m_R^2)} \cdot \left(\frac{2 f_{2n}(a_{1R})}{m_R^4} + \frac{1}{(m_L^2 - m_R^2)} \left(\frac{f_{3n}(a_{1L})}{m_L^2} - \frac{f_{3n}(a_{1R})}{m_R^2} \right) \right) \right], \qquad (4.7)$$

respectively, and $a_{1L,2L} = M^2_{1,2}/m^2_L$, $a_{1R} = M^2_1/m^2_R$ and $b_{L,R} = \mu^2/m^2_{L,R}$. The loop functions $f_{i(c,n)}(x)$'s are such that $f_{i(c,n)}(x,y) = f_{i(c,n)}(x) - f_{i(c,n)}(y)$ with.

$$f_{2n}(a) = \frac{-5a^2 + 4a + 1 + 2a(a+2)\ln a}{4(1-a)^4} \qquad f_{3n}(a) = \frac{1 + 2a\ln a - a^2}{2(1-a)^3}$$

$$f_{2c}(a) = \frac{-a^2 - 4a + 5 + 2(2a+1)\ln a}{2(1-a)^4}\,.$$

As we can see from eq. (4.6), the A_L amplitude includes both U(1) (the terms proportional to $\tan^2\theta_W$) and SU(2) type contributions. The U(1) contributions correspond to pure $\tilde{B}$ exchange, with chirality-flip in the internal fermion line or to $\tilde{B} - \tilde{H}^0$ exchange with chirality flip realized at the Yukawa vertex. For the SU(2) case, we have $\tilde{W} - \tilde{H}$ exchange both for charginos and for neutralinos. However, given that $\tilde{W}$ fields do not couple to right-handed fields, pure $\tilde{W}$ exchange can not mediate any contribution with internal sfermion line chirality flip in contrast to the U(1) case. On the contrary, the A_R amplitude receives only U(1) contributions. As regard the A_R amplitude, we observe that it suffers from some cancellations among different contributions in regions of the parameter space. The origin of these cancellations is the destructive interference between the contributions coming from the $\tilde{B}$ (with internal chirality flip) and $\tilde{B}\tilde{H}^0$ exchange. It is easy to check numerically that these contributions, have opposite sign in all the parameter space. On the other hand the same type of contributions in the δ_{LL} case have the same sign. The reason for this difference is the opposite sign in the hypercharge of SU(2) doublets and singlets.

Finally, we observe that, if the SUSY model contains both δ^{23}_{LL} and δ^{31}_{RR} MI types, we get an additional contribution so that $A^{\rm tot}_L = A_L + A^{'}_L$ with $A^{'}_L$ given by:

$$
\begin{aligned}
A^{'}_L = &- \frac{\alpha_1}{2\pi}\left(\frac{m_\tau}{m_\mu}\right)\mu M_1 t_\beta \frac{\delta^{23}_{LL}\delta^{31}_{RR}}{(m^2_L - m^2_R)^2}\cdot \\
&\cdot\left[\frac{f_{2n}(a_L)}{m^4_L} + \frac{f_{2n}(a_R)}{m^4_R} + \frac{1}{(m^2_R - m^2_L)}\left(\frac{f_{3n}(a_R)}{m^2_R} - \frac{f_{3n}(a_L)}{m^2_L}\right)\right].
\end{aligned}
\tag{4.8}
$$

A particularly interesting feature of the above amplitude is the m_τ/m_μ enhancement with respect to the usual Bino-like mediated processes. This is due to the implementation of the chirality flip in the internal sfermion line through $\delta^{LR}_{33} \sim m_\tau \mu \tan\beta$ and not by $\delta^{LR}_{22} \sim m_\mu \mu \tan\beta$, as usual. The $A^{'}_R$ amplitude, relative to $\delta^{RR}_{23}\delta^{LL}_{31}$, is simply obtained by $A^{'}_R = A^{'}_L(L \leftrightarrow R)$. The contribution reported in eq. (4.8) has to be compared with the second term of eq. (4.5) that is the analog contribution in the Higgs mediated LFV case.

We stress that, in eq. (4.8), we have not included contributions proportional to $\delta^{RR,LL}_{23}\delta^{RR,LL}_{31}$ because they are generally suppressed (or at most comparable) compared to those proportional to $\delta^{RR,LL}_{21}$. On the contrary, in eq. (4.5), terms proportional to $\Delta^{RR,LL}_{23}\Delta^{RR,LL}_{31}$ were retained because enhanced by a $(m_\tau/m_\mu)^2$ factor compared to the corresponding effects proportional to $\Delta^{RR,LL}_{21}$.

The processes $\mu \to eee$ and μ–e conversion in Nuclei get contributions not only from penguin-type diagrams (both with photon or Z-boson exchange) but also from box-type diagrams. In fact, the dipole $\mu \to e\gamma^*$ contribution in these processes is also given by eqs. (4.6), (4.7) and therefore is enhanced by a $\tan\beta$ factor. On the other hand the other contributions, monopole or boxes, are not proportional to $\tan\beta$. So the dipole contributions usually dominate specially in the large $\tan\beta$ regime and one can find the simple theoretical relations

$$
\left.\frac{Br(\mu\to eee)}{Br(\mu\to e\gamma)}\right|_{\rm Gauge} \simeq \left.\frac{Br(\mu - e \text{ in Ti})}{Br(\mu\to e\gamma)}\right|_{\rm Gauge} \simeq \alpha_{\rm el}
\tag{4.9}
$$

In order to make the comparison between Higgs and Chargino mediated LFV effects as simple as possible, let us consider the simple case where all the susy particles are degenerate. In this case, it turns out that

$$
\Delta^{21}_L \sim \frac{\alpha_2}{24\pi}\delta^{21}_{LL}\,,
$$

$$
\begin{aligned}
BR(\mu\to e\gamma)\Big|_{\rm Gauge} &= \frac{2\alpha_{\rm el}}{75\pi}\left(1 + \frac{5}{4}\tan^2\theta_W\right)^2\left(\frac{m^4_W}{m^4_{\rm SUSY}}\right)\left(\delta^{21}_{LL}\right)^2 t^2_\beta\,, \\
Br(\mu\to e\gamma)\Big|_{\rm Higgs} &\simeq \frac{3}{2}\frac{\alpha^3_{\rm el}}{\pi^3}\left(\frac{\alpha_2}{24\pi}\right)^2\left(\frac{m^4_W}{M^4_H}\right)\left(F(a_W)\right)^2\left(\delta^{21}_{LL}\right)^2 t^4_\beta\,.
\end{aligned}
\tag{4.10}
$$

In figure 1 we report the branching ratios of the examined processes as a function of the heaviest Higgs boson mass m_H (in the Higgs LFV mediated case) or of the common susy

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mass $m_{\rm SUSY}$ (in the gaugino LFV mediated case). We set $t_\beta = 50$ and we consider the PMNS scenario as discussed in section 2 so that $(\delta^{21}_{LL})_{PMNS} \simeq 10^{-2}$. Subleading contributions proportional to $(\delta^{23}_{LL(RR)}\delta^{31}_{RR(LL)})_{PMNS}$ (see eqs. (4.5), (4.8)) were neglected since, in the PMNS scenario, it turns out that $(\delta^{23}_{LL(RR)}\delta^{31}_{RR(LL)})_{PMNS}/(\delta^{21}_{LL})_{PMNS} \simeq 10^{-3}$ [5]. As we can see from figure 1, Higgs mediated effects start being competitive with the gaugino mediated ones when $m_{\rm SUSY}$ is roughly one order of magnitude larger then the Higgs mass m_H. Moreover, we stress that, both in the gaugino and in the Higgs mediated cases, $\mu \to e\gamma$ gets the largest effects. In particular, within the PMNS scenario, it turns out that Higgs mediated $Br(\mu \to e\gamma) \sim 10^{-11}$ when $m_H \sim 200$GeV and $t_\beta = 50$, that is just closed to the present experimental resolution.

5. Conclusions

In this letter we have studied Higgs-mediated LFV $e - \mu$ transitions in 2HDM and Supersymmetry frameworks. The sources of LFV were parametrized in a model independent way in order to be as general as possible. In particular, we have analyzed $\mu \to e\gamma$, $\mu \to eee$ and $e - \mu$ conversion in nuclei finding that $\mu \to e\gamma$ is generally the most sensitive channel to probe Higgs-mediated LFV. Analytical expressions for the rates of the above processes and their correlations have been established up to two loop level. Particular emphasis was given to the correlations among the processes as an important signature of the theory. In fact, while it is rather difficult to predict the absolute branching ratio value for any given process (depending on the amount of LFV sources and on the mass spectrum), possible correlations with other processes seem to be a more powerful tool to disentangle different scenarios. In this respect, experimental improvements in all the examined $e-\mu$ transitions would be very welcome. On the other hand, we have shown that the Higgs-mediated contributions to LFV processes can be within the present or upcoming experimental resolutions and provide an important chance to detect new physics beyond the Standard Model.

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