Published by Institute of Physics Publishing for SISSA

Received: *September 26, 2005*
Accepted: *November 17, 2005*
Published: *February 17, 2006*

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

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Abstract: We study the phenomenology of the $\mu - \tau$ and $e - \tau$ lepton flavour violation (LFV) in a general two Higgs Doublet Model (2HDM) including the supersymmetric case. We consider several LFV decay modes of the charged fermion τ, namely $\tau \to l_j \gamma$, $\tau \to l_j l_k l_k$ and $\tau \to l_j \eta$. The predictions and the correlations among the rates of the above processes are computed. In particular, it is shown that $\tau \to l_j \gamma$ processes are the most sensitive channels to Higgs-mediated LFV specially if the splitting among the neutral Higgs bosons masses is not below the 10% level.

Keywords: Supersymmetry Phenomenology.

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Contents

1. Introduction

The observation of neutrino oscillation have established the existence of lepton family number violation. As a natural consequence of this phenomenon, one would expect flavour mixing to appear also in the charged leptons sector. This mixing can be manifested in rare decay processes such as $\mu \to e\gamma$, $\tau \to \mu\gamma$ etc. In the Standard Model with massive neutrinos these processes are mediated, at one loop level, by the exchange of the W bosons and neutrinos; however, in analogy to the quark sector, the resulting rates are GIM suppressed and turn out to be proportional to the ratio of masses of neutrinos over the masses of the W bosons. In addition, if neutrinos are massive, we would expect LFV transitions also in the Higgs sector through the decay modes $H^0 \to l_i l_j$ mediated at one loop level by the exchange of the W bosons and neutrinos. However, as for the $\mu \to e\gamma$ and the $\tau \to \mu\gamma$ case, also the $H^0 \to l_i l_j$ rates are GIM suppressed.

In a supersymmetric (SUSY) framework the situation is completely different. Besides the previous contributions, supersymmetry provides new direct sources of flavour violation, namely the possible presence of off-diagonal soft terms in the slepton mass matrices and in the trilinear couplings [1]. In practice, flavour 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 [2, 3].[1]

Another potential source of LFV in models such as the minimal supersymmetric standard model (MSSM) could be the Higgs sector, in fact, extensions of the Standard Model containing more than one Higgs doublet generally allow flavor-violating couplings of the neutral Higgs bosons. Such couplings, if unsuppressed, will lead to large flavor-changing neutral currents in direct opposition to experiments. The MSSM avoid these dangerous couplings at the tree level segregating the quark and Higgs fields so that one Higgs (H_u)

[1] As recently shown in ref. [4], 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.

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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 [5]. In particular, the presence of a non zero μ term, coupled with SUSY breaking, is enough to induce non-holomorphic Yukawa interactions for quarks and leptons. For large $\tan\beta$ values the contributions to d-quark masses coming from non-holomorphic operator QD_cH_u can be equal in size to those coming from the usual holomorphic operator QD_cH_d despite the loop suppression suffered by the former. This is because the operator itself gets an additional enhancement of $\tan\beta$.

As shown in reference [6] the presence of these loop-induced non-holomorphic couplings also leads to the appearance of flavor-changing couplings of the neutral Higgs bosons. These new couplings generate a variety of flavor-changing processes such as $B^0 \to \mu^+\mu^-$, $\bar{B}^0 - B^0$ etc.[7].

Higgs-mediated FCNC can have sizable effects also in the lepton sector [8]: given a source of non-holomorphic 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 [9, 10] and in supersymmetry [11] frameworks. However, so far most of the attention has been devoted to the tree level effects and in particular to the $\tau \to l_j l_k l_k$ and $\tau \to \mu\eta$ processes. On the other hand, the Higgs-mediated FCNC can have a sizable impact also in loop-induced processes, such as $\tau \to l_j\gamma$. The main purpose of this letter is a detailed investigation of these effects (a comprehensive analysis of the $e - \mu$ transitions will be presented in an upcoming letter [12]). We consider, in particular, the additional dipole and monopole operators induced by the Higgs exchange. As a consequence, $\tau \to l_j\gamma$ processes are generated and $\tau \to l_j l_k l_k$ decay rates get additional contributions by the monopole and dipole operators. We perform the analysis both in a general and in a supersymmetric two Higgs Doublet Models.

2. LFV in the Higgs sector

As it is well known, Standard Model extensions containing more than one Higgs doublet generally allow flavor-violating couplings of the neutral Higgs bosons which arise as a consequence of the fact that each fermion type can couple to both Higgs doublets. Such couplings, if unsuppressed, will lead to large flavor-changing neutral currents in direct opposition to experiments. The possible solution to this problem involve 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.

When each fermion type couple to both Higgs doublets, FCNC could be kept under control if there exists a hierarchy among the Yukawa matrices. For instance, it is possible to assume that the model has a flavor symmetry able to reproduce the observed fermion

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masses and mixing angles. Another possibility is that each type of fermion couples to a different Higgs doublet at the tree level, and the coupling with the other Higgs doublet arises only as a radiative effect. In the following we will assume the last scenario. This occurs, for instance, in the MSSM where the type-II 2HDM structure is not protected by any symmetry and is broken by loop effects.

We consider the following generic Yukawa interactions for charged leptons, including the radiatively induced LFV terms:

$$-\mathcal{L} \simeq Y_\mu H_1^0 \overline{\mu}_R \mu_L - Y_\tau H_1^0 \overline{\tau}_R \tau_L + Y_\tau H_2^0 \Delta_L^{3j} \overline{\tau}_R l_L^j + Y_\tau H_2^0 \Delta_R^{3j} \overline{l_R^j} \tau_L + \text{h.c.}\,, \tag{2.1}$$

where the $\Delta_{L,R}^{3j}$ parameters are the source of LFV (for instance, in the MSSM, they are generated at one loop level by the slepton mixing).
In the mass-eigenstate basis for both leptons and Higgs bosons, the effective flavor-violating interactions are described by the four dimension operators:

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

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 $\tan\beta$ is the ratio of the vacuum expectation value for the two Higgs. 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.[2] In order to constrain the $\Delta_{L,R}^{ij}$ parameters, we impose that their contributions to LFV processes as $l_i \to l_j l_k l_k$ and $l_i \to l_j \gamma$ do not exceed the experimental bounds. At tree level Higgs exchange contribute only to $l_i \to l_j l_k l_k$. On the other hand, at the one loop level, also the dipole operators arise and the LFV radiative decays $l_i \to l_j \gamma$ are allowed.

However, the one loop Higgs mediated dipole transition implies three chirality flips: two in the Yukawa vertices and one in the lepton propagator. This strong suppression can be overcome at hiher order level. Going to two loop level, one has to pay the typical price of $g^2/16\pi$ but one can replace the light fermion masss from Yukawa vertices with the heavy fermion (boson) masses circulating in the second loop. In this case, the virtual Higgs boson couple only once to the lepton line, inducing the needed chirality flip. As a result, the two loop amplitude can provide the major effects. Naively, the ratio between the two loop fermionic amplitude and the one loop amplitude is:

$$\frac{A_{l_i \to l_j \gamma}^{(2-loop)_f}}{A_{l_i \to l_j \gamma}^{1-loop}} \sim \frac{\alpha_{em}}{4\pi} \frac{m_f^2}{m_{l_i}^2} \log\left(\frac{m_f^2}{m_H^2} \right),$$

where $m_f = m_b, m_\tau$ is the mass of the heavy fermion circulating in the loop. We remind that in a model II 2HDM the Yukawa couplings between neutral Higgs boson and quarks

[2] On the other hand, there are several models with a specific ansatz about the flavour-changing couplings. For instance, the famous multi-Higgs-doublet models proposed by Cheng and Sher [13] predict that the LFV couplings of all the neutral Higgs bosons with the fermions have the form $Hf_if_j \sim \sqrt{m_i m_j}$.

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are $H\bar{t}t \sim m_t/\tan\beta$ and $H\bar{b}b \sim m_b/\tan\beta$. Since the Higgs mediated LFV is relevant only at large $\tan\beta \geq 30$, it i clear that the main contributions arise from the $_\tau$ and b fermionsand not from the top quark. So, in this framework, $\tau \to l_j\gamma$ do not receives sizable two loop effects by an heavy fermionic loop differently from the $\mu \to e\gamma$ case. However,the situaion can drastically change when a W boson circulates in the loop Barr-Zee diagrams. Bearing in mind that $HW^+W^- \sim m_W$ and that pseudoscalar bosons do not couple to a W pair, it turns out that $A^{(2-loop)_f}_{l_i\to l_j\gamma}/A^{1-loop}_{l_i\to l_j\gamma} \sim m_W^2/(m_f^2\tan\beta)$ thus, two loop W effects are expected to dominate, as it is confirmed numerically [9].

Moreover, up to one loop level, the $l_i \to l_jl_kl_k$ processes get additional contributions induced by $l_i \to l_j\gamma^*$ amplitudes. It is worth noting that the Higgs mediated monopole and dipole amplitudes have the same $\tan^3\beta$ dependence. This has to be contrasted to the non-Higgs contributions. For instance, within susy, the gaugino mediated dipole amplitude is proportional to $\tan\beta$ while the monopole amplitude is $\tan\beta$ independent.

The general expression for the Higgs mediated $l_i \to l_jl_kl_k$ and $l_i \to l_j\gamma$ rates read:

$$
\begin{aligned}
\frac{Br(l_i \to l_jl_kl_k)}{Br(l_i \to l_j\nu_i\nu_j)} &= \frac{1}{8G_F^2}\left[(3+5\delta_{jk})\frac{|S_L^+|^2}{8} + (3+\delta_{jk})\frac{|S_L^-|^2}{4} + e^4\left|M_L\right|^2(2+\delta_{jk}) - \right.\\
&\left. - 4e^4D_L^\gamma M_L(2+\delta_{jk}) + 8e^4\left|D_L^\gamma\right|^2\left(\log\frac{m_{l_i}^2}{m_{l_k}^2} - 3\right) + (L\leftrightarrow R)\right]\\
\frac{Br(l_i \to l_j\gamma)}{Br(l_i \to l_j\nu_i\nu_j)} &= \frac{48\pi^3\alpha_{em}}{G_F^2}\left[\left|D_L^\gamma\right|^2 + \left|D_R^\gamma\right|^2\right],
\end{aligned}
$$

where the scalar $S_{L,R}$, the monopole $M_{L,R}$ and the dipole $D_{L,R}$ amplitudes read:

$$
S^+_{L,R} = \frac{G_F}{\sqrt{2}}\frac{m_{l_i}m_{l_k}}{M_H^2}\frac{1}{c_\beta^3}\left[\frac{c_\alpha s_{\beta-\alpha}}{m_H^2} - \frac{s_\alpha c_{\beta-\alpha}}{m_h^2} + \frac{s_\beta}{m_A^2}\right]\Delta_{L,R} \tag{2.2}
$$

$$
S^-_{L,R} = \frac{G_F}{2}\frac{m_{l_i}m_{l_k}}{M_H^2}\frac{1}{c_\beta^3}\left[\frac{c_\alpha s_{\beta-\alpha}}{m_H^2} - \frac{s_\alpha c_{\beta-\alpha}}{m_h^2} - \frac{s_\beta}{m_A^2}\right]\Delta_{L,R} \tag{2.3}
$$

$$
\begin{aligned}
M_{L,R} &= \frac{G_F}{48\sqrt{2}\pi^2}\frac{m_{l_i}^2}{c_\beta^3}\left[\frac{c_\alpha s_{\beta-\alpha}}{m_H^2}\left(\log\frac{m_{l_i}^2}{M_H^2} + \frac{5}{6}\right) - \frac{s_\alpha c_{\beta-\alpha}}{m_h^2}\left(\log\frac{m_{l_i}^2}{M_h^2} + \frac{5}{6}\right) + \right.\\
&\left. + \frac{s_\beta}{m_A^2}\left(\log\frac{m_{l_i}^2}{M_A^2} + \frac{5}{6}\right)\right]\Delta_{L,R}
\end{aligned} \tag{2.4}
$$

$$
\begin{aligned}
D_L &= -\frac{G_F}{8\sqrt{2}\pi^2}\frac{m_{l_i}^2}{c_\beta^3}\left[\frac{c_\alpha s_{\beta-\alpha}}{m_H^2}\left(\log\frac{m_{l_i}^2}{M_H^2} + \frac{4}{3} - \frac{\alpha_{el}}{\pi}\frac{m_W^2}{m_\tau^2}\frac{F(a_W)}{\tan\beta}\right) - \right.\\
&- \frac{s_\alpha c_{\beta-\alpha}}{m_h^2}\left(\log\frac{m_{l_i}^2}{M_h^2} + \frac{4}{3} - \frac{\alpha_{el}}{\pi}\frac{m_W^2}{m_\tau^2}\frac{F(a_W)}{\tan\beta}\right) - \\
&\left. - \frac{s_\beta}{m_A^2}\left(\log\frac{m_{l_i}^2}{M_A^2} + \frac{5}{3}\right)\right]\Delta_L
\end{aligned} \tag{2.5}
$$

$$
D_R = D_L(L\leftrightarrow R) + \frac{G_F}{48\sqrt{2}\pi^2}\frac{m_{l_i}^2}{M_{H^\pm}^2}\frac{\Delta_R}{c_\beta^3}\,. \tag{2.6}
$$

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where $a_W = m_W^2/m_H^2$. The terms proportional to $F(a_W)$ arise from loop effects induced by Barr-Zee type diagrams with a W boson exchange. The loop function $F(z)$ is given by

$$F(z) \simeq 3f(z) + \frac{23}{4}g(z) + \frac{f(z)-g(z)}{2z} \tag{2.7}$$

with the Barr-Zee loop integrals given by:

$$g(z) = \frac{1}{4}\int_0^1 dx \frac{\log(z/x(1-x))}{z-x(1-x)} \tag{2.8}$$

$$f(z) = \frac{1}{4}\int_0^1 dx \frac{1-2x(1-x)\log(z/x(1-x))}{z-x(1-x)} \tag{2.9}$$

for $z \ll 1$ it turns out that

$$F(z) \sim \frac{35}{16}(\log z)^2 + \frac{\log z + 2}{4z}\,. \tag{2.10}$$

The $\tau \to \mu(e)\eta$ process receives the only contribution from the pseudoscalar A and the resulting branching ratio is:

$$\frac{Br(\tau \to l_j\eta)}{Br(\tau \to l_j\bar{\nu}_j\nu_\tau)} \simeq 9\pi^2\left(\frac{f_\eta^8 m_\eta^2}{m_A^2 m_\tau}\right)^2\left(1-\frac{m_\eta^2}{m_\tau^2}\right)^2\left[\xi_s+\frac{\xi_b}{3}\left(1+\sqrt{2}\frac{f_\eta^0}{f_\eta^8}\right)\right]^2\Delta_{3j}^2\tan^6\beta\,,$$

where $m_\eta^2/m_\tau^2 \simeq 9.5\times10^{-2}$ and the relevant decay constants are $f_\eta^0 \sim 0.2f_\pi$, $f_\eta^8 \sim 1.2f_\pi$ and $f_\pi \sim 92\,\mathrm{MeV}$ [14]. The parameters ξ_f appear in the couplings between the scalar and the fermions $-i(\sqrt{2}G_F)^{1/2}\tan\beta H\xi_f m_f\overline{f}f$. Although they are equal to one at tree level they can get large corrections from higher order effects. This is the case, for instance, of Susy where contributions arising from gluino-squark loops (proportional to $\alpha_s\tan\beta$) can enhance or suppress significantly the tree level value of ξ_b [5–7].

2.1 Non-decoupling limit: $\sin(\beta-\alpha)=0$

In this section we will derive the expressions and the correlations among the rates of the above processes in the limiting case where $\sin(\beta-\alpha)=0$ and $\tan\beta$ is large. In particular, we will establish which the most promising channels to detect Higgs mediated LFV are.

For $\tau \to l_j\gamma$ and $\tau \to l_j l_k l_k$ branching ratios we get, respectively

$$\frac{Br(\tau \to l_j\gamma)}{Br(\tau \to l_j\bar{\nu}_j\nu_\tau)} \simeq \frac{3\alpha_{el}}{8\pi}\left(\frac{m_\tau^2}{m_{h,A}^2}\right)^2\left(\log\frac{m_\tau^2}{m_{h,A}^2}+\frac{4}{3}\right)^2\tan^6\beta\Delta_{3j}^2 \tag{2.11}$$

$$\frac{Br(\tau \to l_j l_k l_k)}{Br(\tau \to l_j\bar{\nu}_j\nu_\tau)} \simeq \frac{m_\tau^2 m_{l_k}^2}{8m_{h,A}^4}\Delta_{3j}^2\tan^6\beta\left[\frac{3}{8}(1+\delta_{jk})+\frac{\alpha_{el}^2}{\pi^2}\frac{m_\tau^2}{m_{l_k}^2}\times\right.$$
$$\left.\times\left(\log\frac{m_\tau^2}{m_{l_k}^2}-3\right)\left(\log\frac{m_\tau^2}{m_{h,A}^2}+\frac{4}{3}\right)^2\right], \tag{2.12}$$

where we have retained only the dominant contribution from the lightest h or A Higgs bosons.

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In the above expressions we disregarded subleading two loop effects although they are retained in the numerical analysis. On the other hand, two loop effects provide a sizavle reduction of $Br(\tau \to l_j\gamma)$ and $Br(\tau \to l_jee)$ in the large m_h regime as it is shown in figure 1. Such effects are not visible in $Br(\tau \to l_j\mu\mu)$ because it is dominated by the tree level Higgs exchange contributions. We note that, while $\tau \to \mu(e)\eta$ rates decouple in the heavy pseudoscalar limit, the $Br(\tau \to l_jl_kl_k)$ and $Br(\tau \to l_j\gamma)$ branching ratios can get additional contributions by the h scalar. The $\tau \to l_jl_kl_k$ rates contain two terms: the first comes from the tree level Higgs exchange, the second from the dipole operator neglecting subdominant contributions by the monopole operator.

In general the one loop induced Higgs contributions have both advantages and disadvantages. The disadvantages consist in the additional α_{el}^2 factor, the advantages consist in the possibility to replace light lepton masses with the mass of the decaying particles; in addition we get an extra large $\log(m_\tau^2/m_h^2)$ factor from the loop functions. We remark that the scalar contributions to $\tau \to l_jee$ are very suppressed compared to the dipole contributions while they are of the same order in the $\tau \to l_j\mu\mu$ cases.

In order to understand which the best candidate to detect LFV among $\tau \to l_jl_kl_k$, $\tau \to l_j\gamma$ or $\tau \to l_j\eta$ is, we derive the following relations:

$$\frac{Br(\tau\to l_j\gamma)}{Br(\tau\to l_j\eta)} \simeq \frac{1}{5}\left(\log\frac{m_\tau^2}{m_A^2}+\frac{4}{3}\right)^2 \sim 10 \tag{2.13}$$

$$\frac{Br(\tau\to l_jee)}{Br(\tau\to l_j\gamma)} \simeq \frac{\alpha_{el}}{3\pi}\left(\log\frac{m_\tau^2}{m_e^2}-3\right) \sim 10^{-2} \tag{2.14}$$

$$\begin{aligned}\frac{Br(\tau\to l_j\mu\mu)}{Br(\tau\to l_j\gamma)} &\simeq \frac{\alpha_{el}}{3\pi}\left(\log\frac{m_\tau^2}{m_\mu^2}-3\right)+\frac{\pi}{\alpha_{el}}\frac{(1+\delta_{j\mu})}{8}\left(\frac{m_\mu^2}{m_\tau^2}\right)\left(\log\frac{m_\tau^2}{m_A^2}+\frac{4}{3}\right)^{-2}\\ &\sim [2+3(1+\delta_{j\mu})]\cdot 10^{-3}\,,\end{aligned} \tag{2.15}$$

where the last equalities in eqs. (2.13), (2.15) are obtained by setting $m_A = 150\,\text{GeV}$. In general, the above equations imply that $\tau \to l_j\gamma$ is dominant with respect to $\tau \to l_jl_kl_k$ or $\tau \to l_j\eta$ in the not decoupling limit. In addition, we stress that a tree level Higgs exchange predicts that $Br(\tau \to l_j\mu\mu)/Br(\tau \to l_jee) \sim m_\mu^2/m_e^2$ while at the one loop level one gets:

$$\begin{aligned}\frac{Br(\tau\to l_j\mu\mu)}{Br(\tau\to l_jee)} &\simeq 0.2+15(1+\delta_{j\mu})\left(\log\frac{m_\tau^2}{m_A^2}+\frac{4}{3}\right)^{-2}\\ &\sim \Big[2+3(1+\delta_{j\mu})\Big]\cdot 10^{-1}\,,\end{aligned} \tag{2.16}$$

where the last relation in eq. (2.16) holds for $m_A = 150\,\text{GeV}$. In particular, eq. (2.16) allow us to conclude that, in the not decoupling limit, $\tau \to l_jee$ is more sensitive to Higgs mediated LFV than $\tau \to l_j\mu\mu$, as it is reproduced by figure 1.

2.2 Decoupling limit: $\cos(\beta-\alpha) = 0$

In the decoupling limit, where $\cos(\beta-\alpha) = 0$ and $m_Z/m_{A^0} \to 0$, the couplings of the light Higgs boson h^0 are nearly equal to those of the SM Higgs boson. This is a particularly

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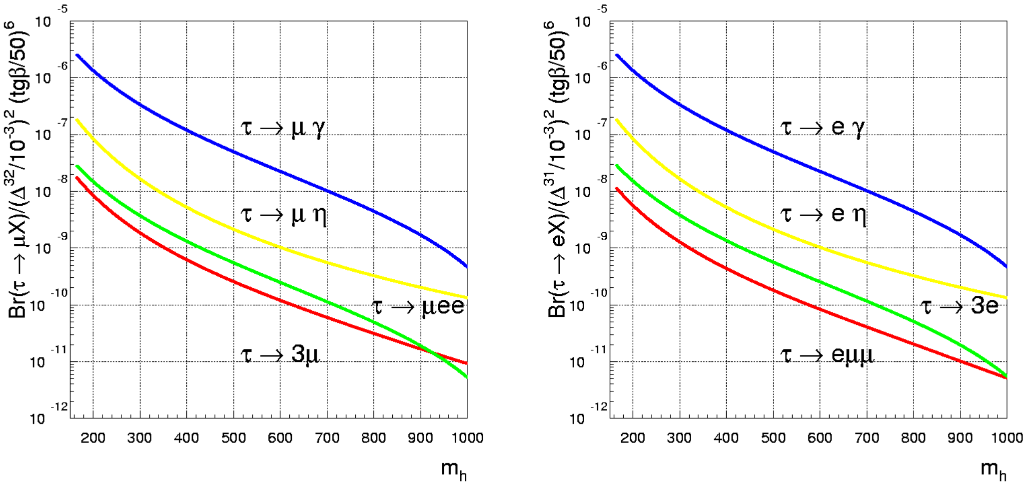

Figure 1: Branching ratios of various $\tau \to \mu$ and $\tau \to e$ LFV processes vs the lightest Higgs boson mass m_A in the non decoupling limit. In the figures we assume $X = \gamma, \mu\mu, ee, \eta$.

interesting limit being that achieved in the Susy framework. In the decoupling limit, $m_{A^0} \simeq m_{H^0} \simeq m_{H^\pm}$ (the mass differences are of order $\mathcal{O}(m_Z^2/m_{A^0})$) and, in particular, the MSSM predicts [15]:

$$m_A^2 - m_H^2 = \frac{\alpha_2 N_c \mu^2}{24\pi M_{\rm susy}^4} \left(\frac{A_t^2 m_t^4}{s_\beta^4 m_W^2} + \frac{A_b^2 m_b^4}{c_\beta^4 m_W^2} \right), \tag{2.17}$$

where $A_{t,b}$ are parameters appearing in the trilinear scalar couplings, μ is the mixing mass between the two Higgs in the superpotential and $M_{\rm susy}$ is a typical susy scalar mass. It turns out that pseudoscalar and scalar one loop amplitudes have opposite signs so, being $m_A \simeq m_H$, they cancel each other to a very large extent. Since these cancellations occur, two loop effects can become important or even dominant in contrast to the non-decoupling limit case. As final result, we find the following approximate expressions:

$$\frac{Br(\tau \to l_j \gamma)}{Br(\tau \to l_j \bar{\nu}_j \nu_\tau)} \simeq \frac{3\alpha_{el}}{2\pi} \left(\frac{m_\tau^2}{m_A^2} \right)^2 \tan^6\beta \Delta_{\tau j}^2 \Bigg(\frac{\delta m}{m_A} \log\frac{m_\tau^2}{m_A^2} + \frac{1}{6} + \\ + \frac{\alpha_{el}}{\pi} \left(\frac{m_W^2}{m_\tau^2} \right) \frac{F(a_W)}{\tan\beta} \Bigg)^2 \tag{2.18}$$

$$\frac{Br(\tau \to l_j l_k l_k)}{Br(\tau \to l_j \bar{\nu}_j \nu_\tau)} \simeq \frac{m_\tau^2 m_{l_k}^2}{32 m_A^4} \Delta_{\tau j}^2 \tan^6\beta \, [3 + 5\delta_{jk}] + \\ + \frac{\alpha_{el}}{3\pi} \left(\log\frac{m_\tau^2}{m_{l_k}^2} - 3 \right) \frac{Br(\tau \to l_j \gamma)}{Br(\tau \to l_j \bar{\nu}_j \nu_\tau)}, \tag{2.19}$$

where $\delta m = m_H - m_A$. It is noteworthy that one and to loop amplitudes have the same signs. In addition, two loops effects dominate in large positions of the parameter space, specially for large m_H values, where the mass splitting $\delta m = m_H - m_A$ decreases to 0. In

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figure 2 we scan over the $\delta m/m_A$ range allowed by the A_t, A_b, μ, $M_{\rm susy}$ parameters within $A_t{=}A_b{=}0$ (degenerate case) and $A_t{=}A_b{=}\mu{=}2M_{\rm susy}$. This choice of the parameter space is phenomenologically available and, in particular, is compatible with the experimental bounds on the lightest stop and Higgs boson masses.

To get a feeling of the allowed rates for Higgs-mediated LFV decays in Supersymmetry it is useful to specify the Δ_{3j} expressions in terms of the susy parameters. We remind that the Δ_{3j} terms are induced at one loop level by the exchange of gauginos and sleptons. Assuming that all the susy particles are of the same order of magnitude but μ (μ being the Higgs mixing parameter), it turns out that

$$\Delta_{3j} \sim \frac{\alpha_2}{24\pi}\frac{\mu}{m_{\rm SUSY}}\,\delta_{3j},$$

where δ_{3j} is the LFV insertions in the slepton mass matrices. The above expression depends only on the ratio of the susy mass scales and it does not decouple for large $m_{\rm SUSY}$.

The unknown δ_{3j} parameters can be determined only if we specify completely the LFV susy model. In figure 2 we have taken the normalization $\Delta_{3j} = 10^{-3}$ that requires, in general, large $\delta_{3j} \sim 1$. The amount of the δ_{3j} mass insertions is constrained by the gaugino mediated LFV and, in general, $\delta_{3j} \sim 1$ requires $m_{\rm SUSY} \sim 1\,{\rm TeV}$ to not exceed the experimental bounds [16].

The numerical results shown in figure 2 allow us to draw several interesting observations:

- $\tau \to l_j\gamma$ has the largest branching ratios except for a region around $m_H \sim 700\,{\rm GeV}$ where strong cancellations among two loop effects sink their size.[3]. The following approximate relations are found:

  $$\frac{Br(\tau \to l_j\gamma)}{Br(\tau \to l_j\eta)} \simeq \left(\frac{\delta m}{m_A}\log\frac{m_\tau^2}{m_A^2} + \frac{1}{6} + \frac{\alpha_{el}}{\pi}\left(\frac{m_W^2}{m_\tau^2}\right)\frac{F(a_W)}{\tan\beta}\right)^2 \geq 1,$$

  where the last relation is easily obtained by using the approimation for $F(z)$ given in eq. (2.10). If two loop effects were disregarded, then we would obtain $Br(\tau \to l_j\gamma)/Br(\tau \to l_j\eta) \in (1/36, 1)$ for $\delta m/m_A \in (0, 10\%)$. Two loop contributions signiicantly enhance $Br(\tau \to l_j\gamma)$ specially for $\delta m/m_A \to 0$.

- In figure 2, non neglegible mass splitting $\delta m/m_A$ effects can be visible at low m_H regime through the bands of the $\tau \to l_j\gamma$ and $\tau \to l_jee$ process. These effects tend to anish with increasing m_H as it is correctly reproduced in figure 2. $\tau \to l_j\mu\mu$ does not receive visible effects by $\delta m/m_A$terms being dominated by the tree level Higgs exchange.

- As it is shown in figure 2, $Br(\tau \to l_j\gamma)$ is generally larger than $Br(\tau \to l_j\mu\mu)$; thei ratio is regulated by the following approximante relation:

  $$\frac{Br(\tau \to l_j\gamma)}{Br(\tau \to l_j\mu\mu)} \simeq \frac{36}{3+5\delta_{j\mu}}\frac{Br(\tau \to l_j\gamma)}{Br(\tau \to l_j\eta)} \geq \frac{36}{3+5\delta_{j\mu}}\,,$$

[3] For a detailed discussion about the origin of these cancellations and their connectio with non-decoupling properties f two loop W amplitude, see ref. [9]

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where the last relation is valid onlyout of the cancellation region. Moreover, from the above relation it turns out that:

$$\frac{Br(\tau \to l_j \eta)}{Br(\tau \to l_j \mu\mu)} \simeq \frac{36}{3+5\delta_{j\mu}}\,.$$

If we relax the condition $\xi_{s,b} = 1$, $Br(\tau \to l_j \eta)$ can get values few times smaller or bigger than those in figure 2.

- It is noteworthy that a tree level Higgs exchange predicts that $Br(\tau \to l_j ee)/Br(\tau \to l_j \mu\mu) \sim m_e^2/m_\mu^2$ while at two loop level, we obtain (out of cancellation region):

$$\frac{Br(\tau \to l_j ee)}{Br(\tau \to l_j \mu\mu)} \simeq \frac{0.4}{3+5\delta_{j\mu}} \frac{Br(\tau \to l_j \gamma)}{Br(\tau \to l_j \eta)} \geq \frac{0.4}{3+5\delta_{j\mu}}\,.$$

  Let us underline that, in the cancelation region, the lower bound of $Br(\tau \to l_j ee)$ is given by the monopole contributions. So, in this region, $Br(\tau \to l_j ee)$ is much less suppressed than $Br(\tau \to l_j \gamma)$.

The correlations among the rates of the above processes are an important signature of the Higgs-mediated LFV and allow us to discriminate between the gaugino mediated LFV and Higgs-mediated LFV. In fact, in the gaugino mediated case, $Br(\tau \to l_j l_k l_k)$ get the largest contributions by the dipole amplitudes that are $\tan\beta$ enhanced with respect to all the other amplitudes resulting in a precise ratio with $Br(\tau \to l_j \gamma)$, namely $BR(\tau \to l_j l_k l_k)/BR(\tau \to l_j \gamma) \simeq \frac{\alpha_{el}}{3\pi}(\log(m_\tau^2/m_{l_k}^2) - 3)$. Moreover, the gaugino-mediated LFV predicts $BR(\tau \to l_j \mu\mu)/BR(\tau \to l_j ee) \simeq (\log(m_\tau^2/m_\mu^2) - 3)/(\log(m_\tau^2/m_e^2) - 3) \simeq 0.2$.

If some ratios different from the above are discovered, then this would be clear evidence that some new process is generating the $\tau \to l_j$ transition, with Higgs mediation being a leading candidate.

3. Conclusions

In this letter we have studied the allowed rates for Higgs-mediated LFV decays both in a general two Higgs Model and in Supersymmetry.

In particular, we have analyzed the decay modes of the τ lepton, namely $\tau \to l_j l_k l_k$, $\tau \to l_j \gamma$ and $\tau \to l_j \eta$. Analytical relations and correlations among the rates of the above processes have been established at two loop level in the Higgs Boson exchange.

The correlations among the processes are a precise signature of the theory. In this respect experimental improvements in all the decay channels of the τ lepton would be very welcome. We have parametrized the source of LFV in a model independent way in order to be as general as possible. We found that $\tau \to l_j \gamma$ processes are generally the most sensitive channels to probe Higgs-mediated LFV specially if the splitting among the neutral Higgs bosons masses is not below 10%. This condition can be fulfilled if $M_{A(H)} \sim M_W$, that is just the situation in which the Higgs LFV effects are more effective. We have also shown that $\tau \to l_j \eta$ and $\tau \to l_j l_k l_k$ are very usefull probes of this scenario. In conclusion, we can say 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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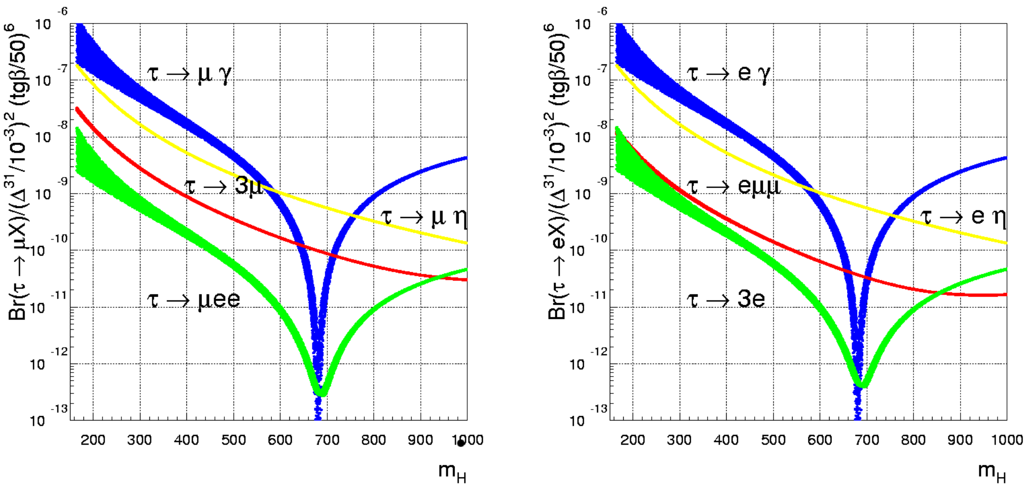

Figure 2: Branching ratios of various $\tau \to \mu$ and $\tau \to e$ LFV processes vs the Higgs boson mass m_A in the decoupling limit. In the figures we assume $X = \gamma, \mu\mu, ee, \eta$. The bands correspond to the allowed $\delta m/m_A$ values as explained in the text.

Acknowledgments

I thank A. Brignole, G. Isidori and A. Masiero for useful discussions.

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