Enhancement of $\mu \to e\gamma$ in the supersymmetric SU(5) GUT at large $\tan \beta$

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Abstract

Branching ratio of $\mu \to e\gamma$ is evaluated in the supersymmetric SU(5) grand unified theory taking account of higher dimensional operators, which reproduce the realistic quark and lepton masses and the mixing parameters. It is shown by numerical calculation that at large $\tan \beta$ the branching ratio is enhanced by a few orders of magnitude since the flavor mixings for both the left- and the right-handed sleptons are induced by the effect of the higher dimensional operators.
The supersymmetric grand unified theory (SUSY GUT), unifying three gauge groups in the Standard Model (SM), is one of the most interesting models from experimental and theoretical points of view. This model explains the electric-charge quantization automatically. The weak mixing angle predicted in this model has been experimentally justified at the 1% level of accuracy. It is, therefore, important to search for the experimental signatures for the SUSY GUT.

It is known that the interaction at the GUT scale ($M_{\text{GUT}} \sim 10^{16}\text{GeV}$) can give effects on the lepton flavor violation (LFV) processes, such as $\mu \rightarrow e\gamma$, through the LFV slepton masses induced by the radiative correction [1]. In particular, the large top-quark Yukawa coupling can be a source of sizable LFV in the slepton masses [3]. The precise values of the event rates depend on the detail of the models [2-6]. In the minimal SU(5) SUSY GUT the Yukawa interaction of the colored Higgs multiplet gives non-negligible LFV masses to only the right-handed sleptons, but not to the left-handed sleptons, and the LFV event rates are significantly reduced by destructive interference among the diagrams [5]. On the other hand, in the SO(10) SUSY GUT two diagrams enhanced by ($m_\tau/m_\mu$) (Fig. 1), which do not interfere with each other, almost dominate in the LFV event processes since both the left-handed and the right-handed sleptons can have LFV masses, and the branching ratio of $\mu \rightarrow e\gamma$ may reach to the present experimental bound [3].

It is known that the mass ratios of down-type quarks and charged leptons in the first and the second generations cannot be explained in the minimal SU(5) SUSY GUT, while the bottom-tau ratio is justified in regions $\tan\beta \simeq 2$ or $\gtrsim 30$ [7]. One of the possible solutions is to introduce higher dimensional operators suppressed by the gravitational scale ($M_G \sim 10^{18}\text{GeV}$). Since the GUT scale is near to the gravitational scale, it is conceivable that the higher dimensional operators may give sizable contribution to the fermion masses in the first and the second generations. In Ref. [4] the contribution from the higher dimensional operators to the LFV processes is considered. It is noted that since both the left-handed and the right-handed sleptons can have sizable LFV masses at large $\tan\beta$, the branching ratios are expected to be enhanced in the general SU(5) SUSY GUT compared with the minimal case. In this paper, we evaluate numerically the branching ratio of $\mu \rightarrow e\gamma$ in the SU(5) SUSY GUT, introducing the higher dimensional operators suppressed by the gravitational scale. We show that for large $\tan\beta$ it is enhanced by a
few orders of magnitude compared with that in the minimal case, and that the destructive interference among the diagrams disappears.

Let us first discuss the minimal SU(5) SUSY GUT to show that the LFV masses for the left-handed sleptons are not induced in this case. Both quarks and leptons are embedded in $\phi(5^*)$ and $\psi(10)$ as follows,

$$
\psi = \frac{1}{\sqrt{2}} \begin{pmatrix}
0 & U & -U & U & D \\
0 & U & U & U & D \\
0 & U & D & 0 & \bar{E} \\
0 & \bar{E} & & & 0
\end{pmatrix},
$$

$$
\phi = \begin{pmatrix} D \\
D \\
\bar{D} \\
E \end{pmatrix} - \begin{pmatrix} N \\
N \\
\bar{N} \end{pmatrix},
$$

where $Q(\equiv (U, D))$ and $L(\equiv (N, E))$ are SU(2)$_L$ doublet (left-handed) quarks and leptons, and $\bar{U}, \bar{D},$ and $\bar{E}$ are SU(2)$_L$ singlet (right-handed) quarks and leptons. The doublet Higgs supermultiplets $H_f$ and $\bar{H}_f$ in the minimal SUSY standard model (MSSM) are embedded in $H(5)$ and $\bar{H}(5^*)$ with the colored Higgs supermultiplets $H_c$ and $\bar{H}_c$,

$$
H = \begin{pmatrix} H_c & H_c & H_f^+ & H_f^0 \end{pmatrix},
$$

$$
\bar{H} = \begin{pmatrix} \bar{H}_c & \bar{H}_c & \bar{H}_f^- & \bar{H}_f^0 \end{pmatrix}.
$$

In order to break the SU(5) gauge symmetry to those of the SM, we introduce an adjoint Higgs supermultiplet $\Sigma(24)$, whose vacuum expectation value is given as

$$
\langle \Sigma \rangle = \begin{pmatrix} 2 & 2 \\
2 & -3 \\
-3 & -3 \end{pmatrix} V.
$$

The renormalizable superpotential leading to the fermion masses is given as

$$
W_R = \frac{1}{4} f_u^{(0)ij} \psi_i^{AB} \psi_j^{CD} H^E \epsilon_{ABCDE} + \sqrt{2} f_d^{(0)ij} \psi_i^{AB} \phi_{Aj} \bar{H}_B
$$

where $i$ and $j$ ($= 1 - 3$) are generation indices and $A, B, \cdots (= 1 - 5)$ are SU(5) ones.

After removing unphysical degrees of freedom in the Yukawa coupling constants, they are given as

$$
f_u^{(0)ij} = V_{km}^{k_i} f_{uk} e^{i\theta_k} V_{km}^{kJ},
$$

$$
f_d^{(0)ij} = f_{d_j} \delta^{ij},
$$

where $\theta_k$ are phases of the CKM matrix.
where $V_{KM}$ corresponds to the Kobayashi-Maskawa (KM) matrix at the GUT scale and $\theta_1$’s ($i = 1 - 3$) are additional phases which satisfy $\theta_1 + \theta_2 + \theta_3 = 0$. After taking a basis of the SM fields as

$$\psi_i \supset \{ Q_i, e^{-i\theta_j}V_{KM}^{ji\ast}U_j, \bar{E}_i \},$$
$$\phi_i \supset \{ \bar{D}_i, L_i \},$$

we get

$$W_R = f_d \bar{E}_i L_i \overline{H}_f + f_d Q_i \bar{D}_i \overline{H}_f + V_{KM}^{ji} f_u Q_i \bar{U}_j H_f + V_{KM}^{ji\ast} f_d e^{-i\theta_j} U_i \overline{D}_j \overline{H}_c,$$

$$+ \frac{1}{2} V_{KM}^{ki} f_u e^{i\theta_j} V_{KM}^{kj\ast} f_d Q_i \bar{Q}_j H_c + V_{KM}^{ij\ast} f_d Q_i \bar{U}_j H_c.$$  (7)

Terms in the first line represent the superpotential of the MSSM. In the fourth term of the right-handed side the right-handed leptons have flavor-violating interaction with the colored Higgs multiplet, which is controlled by the KM matrix. This interaction leads to non-negligible LFV masses for the right-handed sleptons through the radiative correction at one-loop level as,

$$[m^2_{\bar{e}i}]_{ij} = -\frac{3}{8\pi^2} f^2_{u3} V_{KM}^{3i} V_{KM}^{3j\ast}(3 + a_0^2)m_0^2 \log \frac{M_G}{M_{GUT}},$$  (8)

with $i \neq j$. In this paper we assume the minimal supergravity scenario for keeping the universality of SUSY breaking parameters at tree level, and $m_0$ and $a_0$ are the SUSY breaking scalar mass and the trilinear scalar coupling parameter at the tree level \[3\].

Here, we ignored the up-type Yukawa coupling constants except for that of top quark. On the other hand, the Yukawa coupling of the left-handed sleptons to the colored Higgs multiplet is diagonal. Though the off-diagonal components are induced at the higher orders, the effect on the LFV left-handed slepton masses is negligibly small.

Both the Yukawa coupling constants for down-type quarks and those for leptons are given by $f_{di}$ at the GUT scale in the minimal case. As explained in introduction, however, this is not justified at least for the first and the second generations. Since the result that only the right-handed sleptons have LFV masses depends on this unrealistic assumption for the Yukawa coupling, we investigate the LFV processes taking into account higher dimensional operators with $\Sigma$ in the superpotential.
We consider the following superpotential including higher dimensional operators up to the dimension five,

\[ W = W_R + \frac{1}{4M_G} f^{(1)ij}_{u} \left[ \psi_i^{AB} \psi_j^{CD} + \psi_j^{AB} \psi_i^{CD} \right] \Sigma^E_A H^F \epsilon_{BCDEF} \]

\[ + \frac{1}{4M_G} f^{(2)ij}_{d} \left[ \psi_i^{AB} \psi_j^{CD} - \psi_j^{AB} \psi_i^{CD} \right] \Sigma^E_A H^F \epsilon_{BCDEF} \]

\[ + \frac{\sqrt{2}}{M_G} f^{(1)ij}_{d} \left[ \Sigma_A^B \psi_i^B + \Sigma_A^B \psi_i^A \right] \phi_{A \bar{A}} \mathcal{H}_C \]

\[ + \frac{\sqrt{2}}{M_G} f^{(2)ij}_{d} \left[ \Sigma_A^B \psi_i^B - \Sigma_A^B \psi_i^A \right] \phi_{A \bar{A}} \mathcal{H}_C, \]

where \( W_R \) is given by Eq. (5). In order to see explicitly the interaction of the colored Higgs multiplets after inserting the vacuum expectation value into \( \Sigma \), we decompose \( W \) into

\[ W = f^{ij}_{l} E_i L_j \mathcal{H}_f + f^{ij}_{d} Q_i D_j \mathcal{H}_f + f^{ij}_{u} Q_i U_j \mathcal{H}_f \]

\[ + f^{ij}_{cR} L_i \bar{U}_j H_c + f^{ij}_{cL} Q_i \bar{Q}_j H_c \]

\[ + f^{ij}_{cR} U_i \bar{D}_j H_c + f^{ij}_{cL} Q_i \bar{Q}_j H_c. \]

Here, the Yukawa coupling constants of the doublet Higgs multiplets are given as

\[ f^{ij}_{l} = f^{(0)ij}_{d} - \frac{6V}{M_G} f^{(1)ij}_{d}, \]

\[ f^{ij}_{d} = f^{(0)ij}_{d} - \frac{V}{M_G} f^{(1)ij}_{d} + \frac{5V}{M_G} f^{(2)ij}_{d}, \]

\[ f^{ij}_{u} = f^{(0)ij}_{u} - \frac{3V}{2M_G} f^{(1)ij}_{u} + \frac{5V}{2M_G} f^{(2)ij}_{u}, \]

and those of the colored Higgs multiplets are

\[ f^{ij}_{cR} = f^{(0)ij}_{u} + \frac{V}{M_G} f^{(1)ij}_{u} + \frac{5V}{M_G} f^{(2)ij}_{u}, \]

\[ f^{ij}_{cL} = -f^{(0)ij}_{d} + \frac{V}{M_G} f^{(1)ij}_{d} + \frac{5V}{M_G} f^{(2)ij}_{d}, \]

\[ f^{ij}_{cR} = f^{(0)ij}_{d} + \frac{4V}{M_G} f^{(1)ij}_{d}, \]

\[ f^{ij}_{cL} = -\frac{1}{2} f^{(0)ij}_{u} - \frac{V}{2M_G} f^{(1)ij}_{u}. \]
The contribution from the higher dimensional operators to the Yukawa coupling constants of the doublet Higgs multiplets can be as large as $V/M_G \sim 10^{-2}$ without violation of unitarity. Then, from Eq. (11) the fermion masses for the first and the second generations are determined by both the renormalizable and the higher dimensional terms, while those of the third generation are almost determined by the renormalizable terms. In this case, the Yukawa coupling constants of the colored Higgs multiplet to the right-handed leptons are not necessarily diagonalized by the KM matrix unlike in the minimal case given in Eq. (7). More remarkably, we can not diagonalize both the Yukawa coupling constants of the left-handed leptons to the doublet and the colored Higgs multiplets simultaneously unless both $f_d^{(1)}$ and $f_d^{(2)}$ are vanishing. This means that the left-handed sleptons may also receive LFV masses from the radiative correction.

Let us evaluate the LFV event rate of $\mu \to e\gamma$. Since degrees of freedom in the higher dimensional operators are huge, we parameterize the Yukawa couplings as follows,

\begin{align}
  f_l^{ij} &= f_l \delta^{ij}, \\
  f_d^{ij} &= f_d \delta^{ij}, \\
  f_u^{ij} &= V_{KM}^{ji} f_u^{ij},
\end{align}

(13)

and

\begin{align}
  f_{eR}^{ij} &= V_{eR}^{ki} f_{eR} f_{U1}^{kj}, \\
  f_{eL}^{ij} &= V_{eL}^{ki} f_{eL} f_{U1}^{kj}, \\
  f_{\tau R}^{ij} &= V_{\tau R}^{ki} f_{\tau R} f_{U1}^{kj}, \\
  f_{\tau L}^{ij} &= V_{\tau L}^{ki} f_{\tau L} f_{U1}^{kj},
\end{align}

(14)

The LFV SUSY breaking parameters of slepton are given at one-loop level keeping the logarithmic terms as

\begin{align}
  [m_{\tilde{e}}^2]^{i}_{j} &= -\frac{3}{8\pi^2} f_{eR}^{2} V_{eR}^{3i} V_{e}^{3j} (3 + a_0^2) m_0^2 \log \frac{M_G}{M_{GUT}}, \\
  [m_{\tilde{\tau}}^2]^{i}_{j} &= -\frac{3}{8\pi^2} f_{\tau L}^{2} V_{\tau L}^{3i} V_{\tau}^{3j} (3 + a_0^2) m_0^2 \log \frac{M_G}{M_{GUT}},
\end{align}

(15)
\[ A^{ij}_l = -\frac{9}{16\pi^2} \left( f^2_{\tau L3} V^i_{1^3i} V^j_{1^3j^*} f^i_l + f^i_l f^2_{c R3} V^i_{e^3} V^j_{e^3} \right) a_0 m_0 \log \frac{M_G}{M_{\text{GUT}}}, \]  

with \( i \neq j \). Here, they are defined as

\[ -\mathcal{L}_{\text{SB}} = [m^2_{e^3_i}] e^{z^3_i} e_j^* + [m^2_{e^3_i}] l^{i*} l_j + (A_{e}^{i} e_{j} l_{i} h_{f} + \text{h.c.}) \]  

where \( \tilde{e} \) and \( \tilde{l} \) are the right-handed and the left-handed sleptons, and \( h_f \) is one of the doublet Higgs bosons. Also, we assume the hierarchical structure of \( f_{c Rk} \) and \( f_{\tau Lk} \). From the above equations, the LFV masses depend on the magnitude of \( f_{c R3} \) and \( f_{\tau L3} \), and the mixing matrices \( V_{e} \) and \( V_{l} \). Notice that when \( \tan \beta \) is large, the off-diagonal terms in the left-handed slepton mass matrix can be as large as those of the right-handed sleptons provided that \( V^3 V^i_{1^3i} \sim V^3 V^j_{e^3} \). In such a case the branching ratio of \( \mu \rightarrow e\gamma \) is almost determined by two diagrams (Fig. 1), which are proportional to the mass of tau lepton. They do not generally interfere with each other. If only the right-handed sleptons have the sizable LFV masses, every diagram contributing to \( \mu \rightarrow e\gamma \) is proportional to the muon mass and destructive interference tends to occur. Therefore, while at small \( \tan \beta \) the effect from the higher dimensional operators is small, at large \( \tan \beta \) the destructive interference disappears and the branching ratio can be larger than in the minimal case.

In Fig. 2, we show dependence of the branching ratio of \( \mu \rightarrow e\gamma \) on the right-handed selectron mass \( m_{e_R} \) for tan \( \beta \) = 6 and 30. As an illustration, we assume \( V_{e} = V_{l} = V_{\text{KM}}, f_{l3} = f_{\tau R3} = -f_{\tau L3}, \) and \( f_{u3} = f_{c R3} = -f_{c L3}/2 \) at the GUT scale. In this and next figures, the lines denoted as "nonminimal" are calculated on this assumption. This is realized, for example, at \( \tan \beta = 30 \) by taking the following choice,

\[ f^{(0)} = \begin{pmatrix} 9.8 \times 10^{-5} & -4.5 \times 10^{-4} & 3.2 \times 10^{-3} \\ -4.5 \times 10^{-4} & 2.3 \times 10^{-3} & -2.3 \times 10^{-2} \\ 3.2 \times 10^{-3} & -2.3 \times 10^{-2} & 0.65 \end{pmatrix}, \]  

\[ f^{(0)} = \begin{pmatrix} 1.6 \times 10^{-4} & 8.3 \times 10^{-6} & 1.1 \times 10^{-7} \\ -1.8 \times 10^{-3} & 7.7 \times 10^{-3} & 2.8 \times 10^{-4} \\ 6.9 \times 10^{-4} & -4.8 \times 10^{-3} & 0.21 \end{pmatrix}, \]  

\[ f^{(1)} = \begin{pmatrix} 1.6 \times 10^{-3} & 1.4 \times 10^{-4} & 1.8 \times 10^{-6} \\ -2.9 \times 10^{-2} & -0.09 & 4.6 \times 10^{-3} \\ 1.2 \times 10^{-2} & -8.1 \times 10^{-2} & -0.39 \end{pmatrix}, \]  

\[ f^{(2)} = \begin{pmatrix} 1.6 \times 10^{-3} & -1.4 \times 10^{-4} & -1.8 \times 10^{-6} \\ 2.9 \times 10^{-2} & -8.4 \times 10^{-2} & -4.6 \times 10^{-3} \\ -1.2 \times 10^{-2} & 8.1 \times 10^{-2} & -0.39 \end{pmatrix}. \]
and $f_u^{(1)}$ and $f_u^{(2)}$ are zero. Also, we show the branching ratio in the minimal case where $V_\ell$ is $V_{\text{KM}}$ and $V_t$ is the unit matrix. In this figure we take the bino mass 60GeV, $a_0 = 0$, the higgsino mass positive, and the top quark mass 175GeV. In our actual calculation, we solve the renormalization group equations of the diagonal elements in the sfermion mass matrices and use Eqs. (15,16,17) for the off-diagonal components in the slepton mass matrices. Also, we impose the radiative breaking condition of the SU(2)$_L \times$U(1)$_Y$ symmetry, and the experimental constraints including the anomalous magnetic dipole moment of muon. The detailed formula of $\mu \to e \gamma$ is given in Ref. [9]. We also show the experimental upper bound given in Ref. [10]. For $\tan \beta = 30$ the destructive interference at $m_{\tilde{e}_R} \approx 300\text{GeV}$ in the minimal case disappears in the nonminimal case, and the branching ratio is enhanced by two orders of magnitude for $m_{\tilde{e}_R} > 400\text{GeV}$. On the other hand, for $\tan \beta = 6$ the branching ratio is almost the same as in the minimal case.

In Fig. 3 the branching ratio of $\mu \to e \gamma$ is shown as a function of $\tan \beta$ for the minimal and the nonminimal case. Here, $m_{\tilde{e}_R} = 200\text{GeV}$ and 300GeV and the other input parameters are taken as same as in Fig. 2. At large $\tan \beta$ the branching ratio is enhanced in the nonminimal case, and it can reach to one order of magnitude below the experimental bound. Notice that the branching ratio is proportional to $|V_{\tilde{e}33}V_{\tilde{e}32}^*V_{l33}V_{l31}^*|^2$ or $|V_{\tilde{e}33}V_{\tilde{e}31}^*V_{l33}V_{l32}^*|^2$ at large $\tan \beta$. For example, if the mixing angles $V_{l31}$ or $V_{l32}$ is four times larger than those in the KM matrix, a part of the parameter space is already excluded by the experimental bound.

Comments on other processes which may be important for large $\tan \beta$ are in order. First, the EDM of electron may be induced at one-loop level if both the left-handed and the right-handed sleptons have LFV masses. It is proportional to square root of the branching ratio of $\mu \to e \gamma$ at large $\tan \beta$ as is shown in Ref. [4],

$$\left|\frac{d_e}{e}\right| = 3.5 \times 10^{-26}\text{c.m.} \times \frac{\sqrt{2\text{Im}V_{\tilde{e}31}^*V_{\tilde{e}33}V_{l31}^*V_{l33}^*}}{\sqrt{|V_{\tilde{e}31}V_{\tilde{e}33}^*|^2 + |V_{l31}V_{l33}^*|^2}} \left(\frac{\text{Br}(\mu \to e \gamma)}{4.9 \times 10^{-11}}\right)^{\frac{1}{2}}. \quad (23)$$

This depends on the phase and the mixing matrixes, and the EDM of electron can be as large as the experimental upper bound $|d_e/e| \leq 4 \times 10^{-27}\text{c.m.}$ [11]. Next, the amplitude of $b \to s \gamma$ transition is enhanced when $\tan \beta$ is large. In the SM the $b \to s \gamma$ process is induced by an effective operator $\bar{s}_L i \sigma_{\mu \nu} F^{\mu \nu} b_R$. In the MSSM there are contributions
to the operator from the loop of charged Higgs boson and top quark and the loop of chargino and stop. It is known that the latter contribution interferes constructively or destructively with the SM contribution, depending on the SUSY parameters. Moreover, in the SUSY GUTs there are additional contributions to the above operator as well as another operator $\bar{s}_R i \sigma_{\mu \nu} F^{\mu \nu} b_L$. Since these contributions depend on unknown mixing parameters at the GUT scale, in above analysis we do not impose the $b \to s \gamma$ constraint although we have checked at least the MSSM contributions do not exceed the experimental upper bound. Finally, the exchange of the colored Higgs multiplets induces instability of proton, and no observation of proton decay gives the strongest constraint \[1^2\] on

$$f_{ik}^{ij} f_{kl}^{kl} \frac{E_i U_j U_k D_l}{M_{H_c}^{eff}}$$

In this analysis, we do not include the constraint from proton decay, since it is sensitive to the precise value of the mass parameter $M_{H_c}^{eff}$ which can be different from the colored Higgs mass in some models \[1^3\].

In summary, we have evaluated the branching ratio of $\mu \to e \gamma$ in the supersymmetric SU(5) grand unified theory taking account of the higher dimensional operators, and have shown that it reaches to one order of magnitude below the experimental bound for $\tan \beta > 30$ even if $V_e = V_l = V_{KM}$. If the mixing angles are larger, the branching ratio is further enhanced.
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Figure 1: Feynman diagrams contributing to $\mu \rightarrow e\gamma$, which are proportional to $m_\tau$. In these diagrams, $\tilde{e}_{L(R)}$, $\tilde{\mu}_{L(R)}$, and $\tilde{\tau}_{L(R)}$ are the left-handed (right-handed) selectron, smuon, and stau respectively, and $\tilde{B}$ is bino. The symbol $\mu$ refers to the higgsino mass. The arrows represent the chirality.

Figure 2: Dependence of the branching ratio of $\mu \rightarrow e\gamma$ on the right-handed selectron mass in our model for $\tan \beta = 6$ (dashed lines) and $30$ (solid lines). The thick lines are for the nonminimal case in which $V_{\bar{e}}$ and $V_l$ are the same as $V_{KM}$, and the thin lines are for the minimal case in which $V_{\bar{e}} = V_{KM}$ and $V_l = 1$. In this figure we choose the bino mass $60\text{GeV}$, $a_0 = 0$, the higgsino mass positive, and the top quark mass $175\text{GeV}$. The long-dashed line is the experimental upper bound.

Figure 3: Dependence of the branching ratio of $\mu \rightarrow e\gamma$ on $\tan \beta$ for the right-handed selectron mass $200\text{GeV}$ (dashed lines) and $300\text{GeV}$ (solid lines). The thick lines are for the nonminimal case that $V_{\bar{e}}$ and $V_l$ are the same as $V_{KM}$, and the thin lines are for the minimal case in which $V_{\bar{e}} = V_{KM}$ and $V_l = 1$. In this figure we choose the bino mass $60\text{GeV}$, $a_0 = 0$, the higgsino mass positive. The long-dashed line is the experimental upper bound.
\[\begin{array}{c}
\tilde{B} \\
\mu_L \quad \tilde{\mu}_L \quad \tilde{\tau}_L \quad \tilde{\tau}_R \quad \tilde{e}_R \quad e_R \\
\mu_R \quad \tilde{\mu}_R \quad \tilde{\tau}_R \quad \tilde{\tau}_L \quad \tilde{e}_L \quad e_L
\end{array}\]

\[\begin{array}{c}
[m^2_{\tau}^3] \quad m_\tau \mu \tan \beta \quad [m^2_{\tau}^2] \\
[m^2_{\tau}^2] \quad m_\tau \mu \tan \beta \quad [m^2_{\tau}^1]
\end{array}\]
FIG. 2

$M_1 = 60, \ a_0 = 0, \ \mu > 0$

$\text{Br}(\mu \rightarrow \mu \gamma)$

$\tan \beta = 30$ (nonminimal)

$\tan \beta = 30$ (minimal)

$\tan \beta = 6$ (nonminimal)

$\tan \beta = 6$ (minimal)
FIG. 3

\( M_1 = 60, \ a_0 = 0, \ \mu > 0 \)

Br(\( \mu \rightarrow e\gamma \))

- \(-- m_{\nu_R} = 200\) (nonminimal)
- \(-- m_{\nu_L} = 200\) (minimal)
- \( m_{\nu_R} = 300\) (nonminimal)
- \( m_{\nu_R} = 300\) (minimal)