Decoupling and lepton flavor violation in extra dimensional theory

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Abstract

We discuss the fermion mass hierarchy and the flavor mixings in the fat brane scenario of five dimensional SUSY theory. The decoupling solution of the sfermion mass spectrum can be realized by introducing the vector-like mirror fields in an extra dimension. In this scenario, both the left- and right-handed sleptons can have sizable flavor mixings. We point out that this sizable flavor mixings can induce the suitable magnitude of the muon anomalous magnetic moment ($g_\mu - 2$) within the experimental bounds of lepton flavor violating processes.
1 Introduction

Utilizing extra dimensions has shed new insights into various phenomenological aspects of the physics in four dimensions. Antoniadis [1] proposed the possibility that part of the standard model particles live in TeV extra dimensions, in connection to the problem of supersymmetry breaking. One of the advantages of theories with extra dimensions is that small parameters in the theory can be naturally obtained due to the locality not the symmetry. In particular, Arkani-Hamed and Schmaltz [2] have proposed an interesting mechanism, which is referred to as “fat brane scenario”, in which small parameters are obtained by a small overlap of wave functions, even if the parameters in a fundamental theory are of order unity. This mechanism has been applied to various phenomenological issues so far, such as the fermion mass hierarchy [3, 4, 5, 6, 7, 8, 10, 9], the doublet-triplet splitting [11, 12, 13], and the sfermion mass generation [4, 14, 7, 15, 10].

In the previous paper, two of the present authors (N.H. and N.M.) have discussed the fermion mass hierarchy and the flavor mixings in the fat brane scenario of a five dimensional supersymmetric (SUSY) theory [10]. In our set up, the matter lives in the bulk, its zero mode wave functions are Gaussian and are localized at different points in extra dimensions. On the other hand, Higgs fields are localized on a brane. The fermion mass hierarchy is determined by the values at a brane where Higgs fields are localized. Various types of the matter configurations were found, which yield the fermion mass matrices consistent with experimental data.

As for the sfermion mass spectrum, they are generated by the overlap between the wave functions of matter fields and the chiral superfield with nonzero vacuum expectation value (VEV) of the F-component localized on the SUSY breaking brane. In Ref. [15], we have proposed that if SUSY breaking brane is located between the 1st and the 2nd generations, the sfermion mass spectrum becomes the decoupling solution (sometime called as “effective SUSY”) [16]. In this solution, the squarks and sleptons in the 1st and 2nd generations are heavy enough so that their contributions to the FCNC or CP violating processes are sufficiently suppressed.1

On the other hand, the recent experiments again suggest the discrepancy of the muon anomalous magnetic moment from the standard model (SM) prediction [17]. One of the uncertainties of the SM predictions is coming from the evaluation of the hadronic vacuum

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1On the other hand, the gauginos, higgsinos and the sfermions in the 3rd generation are appropriately light to satisfy the naturalness condition on the Higgs boson mass.
polarization. Davier et al. [18] made careful analysis of this contribution (one based on the cross section of $e^+e^-$ to hadrons and the other one based on $\tau$ decay.). The result based on $e^+e^-$ cross section suggests that the SM prediction is about $3\sigma$ deviation from the experimental data, while the result based on $\tau$ decay shows that the SM prediction is consistent with the experimental data. The origin of this difference has not been clarified. In this paper, we take the $e^+e^-$-based result by Davier et al. for the hadronic contributions. It has been shown in Ref.[19] that the lepton-flavor changing process can induce the suitable magnitude of the muon anomalous magnetic moment ($g_\mu - 2$) in the decoupling solution of the minimal supersymmetric standard model (MSSM) satisfying the current experimental bound of the branching ratio of $\tau \to \mu\gamma$.

In this paper we consider the 5 dimensional theory where SUSY breaking brane is put on the center of the right-handed electron of the 2nd and the 3rd generation. In this setup, the SUSY decoupling solution is obtained and the sizable flavor mixings both in the left- and right-handed sleptons can induce the suitable magnitude of $g_\mu - 2$ within the experimental bounds of the lepton flavor violating processes.

## 2 Decoupling in the extra dimension

Let us start to see how the fermion mass hierarchies and SUSY decoupling solutions come out [10].

Consider the up-type Yukawa coupling, for example,

$$W = \int dy \delta(y) Q_i(x, y) \bar{U}_j(x, y) H_u(x),$$

where $x$ denotes the coordinate of four dimensional Minkowski space-time, $y$ is a fifth spatial coordinate of five dimensions. $i, j$ are the generation indices. The order one coefficient is implicit. $Q_i, \bar{U}_i$ and $H_u$ are the chiral superfield which transform as $(3, 2, 1/6)$, $(\bar{3}, 1, -2/3)$ and $(1, 2, 1/2)$ under the Standard Model (SM) gauge groups, $SU(3)_C \times SU(2)_L \times U(1)_Y$. We assume here that the MSSM matter fields live in the bulk and Higgs fields are localized on a brane at $y = 0$. Integrating out the fifth dimensional degrees of freedom, we obtain the effective Yukawa coupling in four dimensions at the compactification scale as,

$$(y_{\text{eff}})_{ij} \simeq \exp[-a^2(y_{Q_i}^2 + y_{\bar{U}_j}^2)],$$
where we assume the form of the zero mode wave function of the matter superfields to be Gaussian such as $\exp[-a^2(y - y_{\Phi_i})^2]$, where $a$ is the inverse width of the zero mode wave functions, $y_{\Phi_i}$ is the coordinate where the matter superfield $\Phi_i(= Q_i, \bar{U}_i, \bar{D}_i, ...)$ is localized. As is clear from (2), the information of Yukawa hierarchy is interpreted as the “geography” of configuration of the matter fields in the extra dimensions. In Ref.[10], we have found various types of fermion mass matrices well describing the fermion mass hierarchies and the flavor mixings in the fat brane scenario.

We note that the sfermion masses correlate with the fermion masses in extra dimensions by introducing SUSY breaking brane because the sfermion masses are determined by the overlap of wave functions of the matter fermions and the chiral superfields with nonvanishing F-term VEV on SUSY breaking brane. Now that we know various matter configurations consistent with experimental data, the sfermion mass spectrum can be calculated and predicted.

Let us discuss the sfermion mass in more detail. In our previous paper, we have proposed the mechanism to generate the sfermion masses in the fat brane scenario [15]. “SUSY breaking brane”, is introduced at $y = L$, where the chiral superfields $X$ with nonvanishing F-term VEV ($X = \theta^2 F$) is assumed to be localized. The extra vector-like superfields $\Phi', \bar{\Phi}'$ with mass $M < M_*$, where $M_*$ is the five dimensional Planck scale, are also introduced and assumed to be localized on a SUSY breaking brane. We consider here the following superpotential,

$$ W = \int dy \delta(y - L)[\frac{\lambda}{\sqrt{M_*}}X(x)\Phi_i(x, y)\bar{\Phi}'(x) + M\Phi(x)\bar{\Phi}'(x)], $$

(3)

$$ = \frac{\lambda}{\sqrt{M_*}\sqrt{a}}\exp[-a^2(L - y_{\Phi_i})^2]X(x)\Phi_i(x)\bar{\Phi}'(x) + M\Phi(x)\bar{\Phi}'(x), $$

(4)

where $\lambda$ is a dimensionless constant of order unity. Below the scale $M$, we can integrate out the massive superfields $\Phi'$ and $\bar{\Phi}'$, then the superpotential vanishes and the effective Kähler potential are generated at tree level,

$$ \delta K_{\text{eff}} = \frac{1}{\sqrt{M_*\sqrt{a}}}\frac{1}{M^2}\exp[-a^2((L - y_{\Phi_i})^2 + (L - y_{\Phi_j})^2)]X^\dagger X\Phi_i^\dagger \Phi_j. $$

(5)

The sfermion masses coming from (5) at the compactification scale are

$$ \tilde{m}_{ij}^2 \simeq \frac{1}{\sqrt{M_*\sqrt{a}}}\exp[-a^2((L - y_{\Phi_i})^2 + (L - y_{\Phi_j})^2)]|F|^2 \frac{1}{M^2}. $$

(6)

It is crucial that the scale suppressing the Kähler potential is replaced with $M < M_*$ not
so as to be negligibly small. Note that $F < M^2$ is assumed in this argument and also the overall sign of the Kähler potential is assumed to be positive.

We would like to mention the gaugino mass in our scenario. The gaugino masses are generated at tree level since we assume that the gauge supermultiplets live in a thick wall,

$$\delta(y - L) \int d^2 \theta \frac{X^\alpha(x)}{M_s^2} W^\alpha(x, y) W(x, y) \Rightarrow M_\lambda = \frac{F}{M_s^2 L_c},$$

where $W_\alpha$ is the field strength tensor superfield and $L_c$ is the width of the thick wall which should be considered as the compactification length in our framework. For the gaugino masses to be around 100 GeV, so we obtain

$$\frac{F}{M_s} \simeq 100(M_s L_c).$$

A-terms are also induced only if the SUSY breaking brane is located at the point where Higgses are located [10].

Now we consider whether the sizable flavor mixings in the left- and right-handed slepton mass matrices can be obtained in the decoupling solutions. The decoupling spectrum of the slepton masses in five dimensional theory has been discussed in Ref.[10]. However, their flavor mixings have not been discussed, which is the issue to be addressed in this paper.

The strategy is the following. In order to obtain the decoupling spectrum and the large mixing between the 2nd and the 3rd generations both in the left- and right-handed sleptons simultaneously, we set SUSY breaking brane at the point where the distances from $\bar{E}_{2,3}$ are the same.

In this setup, the slepton mass matrices becomes as

$$\tilde{m}_{ij}^2 \sim \exp\left[-a^2 \left\{ (s - y_{\Phi_i})^2 + (s - y_{\Phi_j})^2 \right\} \right] \frac{F^2}{M^2},$$

where $s = (y_{\bar{E}_2} + y_{\bar{E}_3})/2$. We will see explicitly the slepton mass matrices for the “anarchy type” fermion mass matrix and improvement I, II [10]. On the other hand, we do not introduce extra vector-like superfields and the relevant superpotential (3) for the quark sector. Therefore, squark masses are negligibly small, radiatively induced at the weak scale by the gaugino RGE effects and flavor blind.

\footnote{Without introducing the extra vector-like superfields, the sfermion masses are negligibly small due to the exponential suppression [15].}
3 Sfermion mass matrices

The “anarchy type” fermion mass matrices with large tan $\beta$ are obtained when the matters are localized in a five dimensional coordinate as $[10]^3$

$$
\begin{align*}
y^{2}_{Q_1} & \simeq -2\ln \epsilon, \\
y^{2}_{Q_2} & \simeq -\ln \epsilon, \\
y^{2}_{Q_3} & \simeq 0, \\
y^{2}_{U_1} & \simeq -2\ln \epsilon, \\
y^{2}_{U_2} & \simeq -\ln \epsilon, \\
y^{2}_{U_3} & \simeq 0, \\
y^{2}_{D_1} & \simeq 0, \\
y^{2}_{D_2} & \simeq 0, \\
y^{2}_{D_3} & \simeq 0, \\
y^{2}_{L_1} & \simeq 0, \\
y^{2}_{L_2} & \simeq 0, \\
y^{2}_{L_3} & \simeq 0, \\
y^{2}_{E_1} & \simeq -2\ln \epsilon, \\
y^{2}_{E_2} & \simeq -\ln \epsilon, \\
y^{2}_{E_3} & \simeq 0, \\
y^{2}_{N_1} & \simeq 0, \\
y^{2}_{N_2} & \simeq 0, \\
y^{2}_{N_3} & \simeq 0.
\end{align*}
$$

(10)

We set the parameter $\epsilon$ to be of order $\lambda^2$, and $\lambda$ is the Cabibbo angle, $\lambda \simeq 0.2$. The configuration (10) generates the following mass matrices for up, down quark sectors and the charged lepton sector:

$$
\begin{align*}
m_u & \simeq \begin{pmatrix}
\epsilon^4 & \epsilon^3 & \epsilon^2 \\
\epsilon^3 & \epsilon^2 & \epsilon \\
\epsilon^2 & \epsilon & 1
\end{pmatrix} \langle H_u \rangle, \\
m_d & \simeq \begin{pmatrix}
\epsilon^2 & \epsilon^2 & \epsilon^2 \\
\epsilon & \epsilon & \epsilon \\
1 & 1 & 1
\end{pmatrix} \langle H_d \rangle, \\
m_l & \simeq \begin{pmatrix}
\epsilon^2 & \epsilon & 1 \\
\epsilon^2 & \epsilon & 1 \\
\epsilon^2 & \epsilon & 1
\end{pmatrix} \langle H_d \rangle.
\end{align*}
$$

(11)

On the other hand, there are no mass hierarchies in the neutrino mass matrix since both left- and right-handed neutrinos are localized at the same point. The light neutrino mass matrix ($\nu^{(l)}$) through the see-saw mechanism$[20]$ is given by

$$
m^{(l)}_{\nu} \simeq \frac{m^D_{\nu}(m^D_{\nu})^t}{m_N} \simeq \begin{pmatrix}
1 & 1 & 1 \\
1 & 1 & 1 \\
1 & 1 & 1
\end{pmatrix} \frac{\langle H_u \rangle^2}{M_R}.
$$

(12)

where $M_R$ is around $10^{15-16}$ GeV.$^4$ All elements of the above matrices have $O(1)$ coefficients. These fermion mass matrices can naturally explain why the flavor mixing in the quark sector is small while that in the lepton sector is large $[21, 22, 23]$. The above fermion mass hierarchies and flavor mixings are roughly consistent with the experimental data, and explicit values of $O(1)$ coefficients of mass matrices can really induce the suitable magnitudes of fermion masses and flavor mixing angles $[21]$.

Here let us show one explicit example of coefficients. According to the method of determining coefficients in Ref.$[21]$, we suggest fermion mass matrices with $O(1)$ coefficients as

$$
\begin{align*}
m_U & = \begin{pmatrix}
0 & 2de^3 & 0 \\
2de^3 & \frac{4}{5}c^2 & 0 \\
0 & b\epsilon & 1
\end{pmatrix}, \\
m_D & = \begin{pmatrix}
d\epsilon^2 & d\epsilon^2 & d\epsilon^2 \\
-d\epsilon & d\epsilon & d\epsilon \\
\frac{c}{2} & b & 1
\end{pmatrix},
\end{align*}
$$

(13)

$^3$These coordinates are in units of $a^{-1}$.

$^4$\textit{This is naturally obtained from the VEV of the singlet field $[10]$}.
\[
m_L = \begin{pmatrix}
e^2 & 0 & 0 \\
b e^2 & -2ce & 0 \\
0 & -be & 5
\end{pmatrix}, \quad m_{\nu}^{(l)} = \begin{pmatrix}
e & e & 0 \\
e & c & 2.5 \\
0 & 2.5 & 5
\end{pmatrix}.
\] (14)

Here we take \(b = 4, c = 3.6, d = 2,\) and \(e = 1.0,\) then the CKM [24] and the MNS [25] matrices are given by

\[
V_{CKM} = U_d^\dagger U_d = \begin{pmatrix}
0.9984 & -0.05650 & -0.000373 \\
0.05649 & 0.9983 & -0.01197 \\
0.001049 & 0.01193 & 0.9999
\end{pmatrix},
\] (15)

\[
U_{MNS} = U_\nu^\dagger U_\nu = \begin{pmatrix}
0.8417 & -0.5322 & 0.0903 \\
-0.4717 & -0.6437 & 0.6025 \\
-0.2625 & -0.5498 & -0.7929
\end{pmatrix},
\] (16)

respectively. Where \(U_u, U_d, U_l,\) and \(U_\nu\) are defined as \((U_u)^\dagger m_u U_t U_u = (m_t^2)_{\text{diagonal}},\)
\((U_d)^\dagger m D m_D^\dagger U_d = (m^2_D)_{\text{diagonal}},\)
\((U_l)^\dagger m L m_L^\dagger U_l = (m^2_L)_{\text{diagonal}},\) and \((U_\nu)^\dagger m_\nu m_\nu^\dagger U_\nu = (m^2_\nu)_{\text{diagonal}},\)
respectively. The fermion mass hierarchies are given by

\[
m_t : m_c : m_u = 1.012 : 0.0045 : 0.00014, \quad m_b : m_s : m_d = 4.49 : 0.12 : 0.00027,
\] (17)

\[
m_\tau : m_\mu : m_\epsilon = 5.00 : 0.28 : 0.0015, \quad m_\nu_\tau : m_\nu_\mu : m_\nu_\epsilon = 6.92 : 1.93 : 0.87,
\] (18)

They are consistent with today’s experimental results. The neutrino mass spectrum is suitable for the LMA solar solution.

Now we calculate the slepton mass matrices. For the decoupling solution, we consider the case that the SUSY breaking brane is located at the point, \(y = s.\) Then the slepton mass matrices take the form as

\[
\tilde{M}_L^2 \simeq \epsilon^{1/2} \begin{pmatrix}
1 & 1 & 1 \\
1 & 1 & 1 \\
1 & 1 & 1
\end{pmatrix} \left(\frac{F}{M}\right)^2,
\] (19)

\[
\tilde{M}_E^2 \simeq \epsilon^{1/2} \begin{pmatrix}
\epsilon^{4-2\sqrt{2}} & \epsilon^{2-\sqrt{2}} & \epsilon^{2-\sqrt{2}} \\
\epsilon^{2-\sqrt{2}} & 1 & 1 \\
\epsilon^{2-\sqrt{2}} & 1 & 1
\end{pmatrix} \left(\frac{F}{M}\right)^2.
\] (20)

In the wide range of parameter regions of order one coefficients, the 3rd generation sfermions become light. The case of the rank of mass matrices being reduced to be 2 is the typical example. This situation is realized when \(\lambda\) and \(M\) in Eq.(3) are common for the 2nd and 3rd generations. In this case the masses of the 1st and the 2nd generations are heavy enough, which is just the realization of decoupling solution. On the other hand, the masses of the 3rd generation can be at least of order 100 GeV through the gravity
mediated effects, which always generate $O(100)$GeV sfermion masses. This gravity effects will be explained later. As for the right-handed slepton, the 1st generation masses are of order 0.15 times smaller than those of the 2nd and 3rd generations as in Eq.(20). However there are parameter regions where the 3rd generation sfermions become light as in the above situation.

Here we show one example of the slepton mass matrix with $O(1)$ coefficients, which induces the decoupling solution. Denoting $6 \times 6$ sfermion mass matrix as

$$\tilde{M}^2 = \begin{pmatrix} \tilde{M}^2_L & \tilde{M}^2_{LR} \\ \tilde{M}^2_{RL} & \tilde{M}^2_R \end{pmatrix},$$  \hfill (21)

we suggest

$$\tilde{M}^2_L \simeq \epsilon^L \begin{pmatrix} 2 & 1 & 1 \\ 1 & 5.7 & 5.64 \\ 1 & 5.64 & 5.6 \end{pmatrix} \left( \frac{F}{M} \right)^2,$$  \hfill (22)

$$\tilde{M}^2_E \simeq \epsilon^E \begin{pmatrix} 3 \times \epsilon^{4-2\sqrt{2}} & -\epsilon^{-2\sqrt{2}} & -\epsilon^{-2\sqrt{2}} \\ -\epsilon^{-2\sqrt{2}} & 5.7 & 5.64 \\ -\epsilon^{-2\sqrt{2}} & 5.64 & 5.6 \end{pmatrix} \left( \frac{F}{M} \right)^2.$$  \hfill (23)

The mass eigenvalues of $\tilde{M}^2_L$ and $\tilde{M}^2_E$ are given by

$$\tilde{m}_{e_L} : \tilde{m}_{\mu_L} : \tilde{m}_{\tau_L} \simeq 6.69 \text{ TeV} : 17.0 \text{ TeV} : 379 \text{ GeV},$$  \hfill (24)

$$\tilde{m}_{e_R} : \tilde{m}_{\mu_R} : \tilde{m}_{\tau_R} \simeq 1.27 \text{ TeV} : 16.8 \text{ TeV} : 380 \text{ GeV},$$  \hfill (25)

where $F/M \simeq \sqrt{5} \times 5$ TeV. We diagonalize these mass matrices in the basis which the fermion mass matrices are diagonal. In this basis $\tilde{M}^2_L$ and $\tilde{M}^2_E$ are changed as

$$M^2_L = U^\dagger_L \tilde{M}^2_L U_L, \quad M^2_E = U^\dagger_R \tilde{M}^2_E U_R,$$  \hfill (26)

where $U_r$ is defined as $U^\dagger_r m_L U_r = (m_L)_{\text{diagonal}}$. $M^2_L$ is diagonalized by the unitary matrix

$$V_L = \begin{pmatrix} 0.9891 & 0.14708 & 0.0033 \\ -0.106 & 0.7011 & 0.70496 \\ 0.1013 & -0.6976 & 0.7092 \end{pmatrix}.$$  \hfill (27)

This means that the mixing angles are

$$\sin \theta_{12} = 0.14708, \quad \sin \theta_{13} = 0.0033, \quad \sin \theta_{23} = 0.70496.$$  \hfill (28)

$M^2_E$ is diagonalized by the unitary matrix

$$V_R = \begin{pmatrix} 0.9999 & -0.002605 & 0.00072 \\ 0.00240 & 0.7325 & -0.680 \\ 0.00124 & 0.6806 & 0.7325 \end{pmatrix}.$$  \hfill (29)
This means that the mixing angles are

\[ \sin\theta_{12} = -0.002605, \quad \sin\theta_{13} = 0.00072, \quad \sin\theta_{23} = -0.680. \tag{30} \]

Above one example of \(O(1)\) coefficients really induce \(O(1)\) mixings between the 2nd and the 3rd generations in both left- and right-handed slepton sectors, and also light 3rd generation sfermion masses. There are wide parameter region where the same situation is satisfied. In Ref.[19] it has been shown that this parameter region can induce enough large muon anomalous magnetic moment, \(g_\mu - 2\), within the constraint of \(\tau \rightarrow \mu \gamma\). The severer constraint exists in the process of \(\mu \rightarrow e \gamma\). However it can be satisfied since the mixing angles between the 1st and the 3rd generations can be of \(O(10^{-3})\) in the wide parameter region as in Eqs.(28) and (30).

The improved mass matrices I (large \(\tan\beta\)) in Ref. [10] induce the same sfermion mass matrices in Eqs.(22) and (23). Thus, this case also realizes the suitable decoupling solution. On the other hand, the matter configurations of “anarchy type” with small \(\tan\beta\) and improved I with small \(\tan\beta\) cannot induce the decoupling solutions. It is because the smallest sfermion mass eigenvalues are given by \(M_L^2 \simeq \epsilon^3 + \sqrt{2}(F/M)^2\) and \(M_E^2 \simeq \epsilon^3 - \sqrt{2}(F/M)^2\) in these cases. They are different in more than two order magnitude from each other, which are not suitable for the decoupling solutions.

As for the improved mass matrices II in Ref. [10], it is not suitable in our scenario as the following argument. Note that \(O(100\text{GeV})\) sfermion mass of the 3rd generation is coming from the gravity mediation as mentioned above,

\[ \delta(y - s) \int d^3\theta \frac{X^1X^1}{M_*^3} Q_iQ_j \rightarrow \tilde{m}_{ij}(\text{gravity}) \sim \frac{F^2a}{M_*^3} \exp[-(s - yQ_i)^2 - (s - yQ_j)^2]. \tag{31} \]

The gravity mediated scalar masses are related to the gaugino mass as

\[ \tilde{m}_{ij}(\text{gravity}) \sim \begin{cases} 0.10(M_*L_c)(aL_c)M_\lambda^2 & (\text{Anarchy, ImprovementI}), \\ 0.76(M_*L_c)(aL_c)M_\lambda^2 & (\text{ImprovementII}). \end{cases} \tag{32} \]

For \(\tilde{m}_{ij}(\text{gravity})\) to be of order 100 GeV,

\[ (M_*L_c)^2 \simeq \begin{cases} 10 & (\text{Anarchy, ImprovementI}), \\ 1 & (\text{ImprovementII}) \end{cases} \tag{33} \]

are obtained if \(a \simeq M_*\) for simplicity. For example, \(M_* \simeq \sqrt{10} \times 10^{16}\text{ GeV}, L_c^{-1} \simeq 10^{16}\text{ GeV}\) is viable for Anarchy, Improvement I. However, Improvement II seems to be unnatural since \(M_*L_c \simeq 1\) contradicts the constraint \(L^{-1} < a \leq M_*\) in the fat brane scenario.
4 Summary

We have discussed the fermion mass hierarchy and the flavor mixings in the fat brane scenario of a five dimensional SUSY theory taking into account of $O(1)$ coefficients. We consider the case where SUSY breaking brane is put on the center of the 2nd and the 3rd generations’ right-handed electron fields in the 5 dimensional coordinate. In this case the decoupling solution is realized. The sizable flavor mixings both in the left- and right-handed sleptons can be naturally induced, which can realize the suitable magnitude of the muon anomalous magnetic moment within the experimental bounds of lepton flavor violating processes.

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