Effective supersymmetric theory and muon anomalous magnetic moment with $R$ parity violation

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Abstract

The effective supersymmetric theory (ESUSY) with $R$ parity conservation cannot give a large anomalous magnetic moment of $\mu$. It is pointed out that the flavor conservation and a large $(g - 2)_{\mu}$ within the experimental limits are achievable in the ESUSY with $R$ parity violating couplings involving the third generation superparticles.

[Key words: anomalous magnetic moment of muon, effective SUSY, $R$ parity violation]

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The recent hint that the anomalous magnetic moment of muon \( a_\mu = \frac{1}{2}(g - 2)_\mu \) might not fall within the predicted range in the standard model attracted a great deal of attention in the hope of probing new physics possibility. Probably, in the future this excitement can be compared to that on the discussion of the weak neutral current thirty years ago.

In general, the loop correction to the magnetic moment of \( \mu \) arises if there exists a left to right (or right to left) chirality transition of the external muon lines. This chirality transition needs an insertion of a fermion mass or a Yukawa coupling vertex. If heavy fermions are introduced, it has been known that a large magnetic moment is possible, with both left-handed and right-handed currents. Without a heavy fermion, the mass insertion or the Yukawa coupling leads to a magnetic moment proportional to the external fermion mass. But in the standard model there does not exist a heavy fermion and we expect small lepton magnetic moments of order \( a_\nu \sim (m_\nu m_e/M_W^2) \sim 10^{-18} \) times the electron Bohr magneton for the eV range neutrino mass and \( a_\mu \sim (m_\mu^2/M_W^2) \) times the muon Bohr magneton. There is a new heavy (but relatively light to resolve the hierarchy problem) fermion in supersymmetric models, i.e. the chargino. Then the above argument may imply an anomalously large anomalous magnetic moment of \( \mu \), which however is not realized due to the chiral nature of the supersymmetric models.

Indeed, the recent observation of the anomalous magnetic moment of muon are around the electroweak scale order, i.e. of order predicted in the standard model

\[
a_\mu^{\exp} - a_\mu^{\text{th}} = 426(165) \times 10^{-11},
\]

but off from the standard model prediction by 2.6\( \sigma \). The theoretical prediction quoted above is from Ref. which gives the most stringent error bar compared to the other published results, the main difference coming from the treatment of the hadronic contributions. Therefore, at present it is premature to conclude the existence of a new physics beyond 100 GeV.

Nevertheless, it is tempting to search for possibilities of generating \( (g - 2)_\mu \) of order the electroweak scale or even a larger value. Then, the mass scale of new physics must be close
to the electroweak scale since the possible deviation is of order the predicted range in the standard model. In this spirit, already there appeared several explanations \[2,3\]. Among these, the supergravity scenario is particularly interesting because a relatively light (mass < 500 GeV) superpartners are around the corner as shown in \[2\], which gives a hope of probing whole spectrum of superparticles below TeV at TevatronII and LHC.

However, the supergravity models have the notorious problem regarding the flavor changing processes \[9\]. One of the models resolving this flavor changing problem is the gauge mediated supersymmetry breaking at low energy \[9\]. Another model is the effective supersymmetric theory(ESUSY) in which the superpartners of the first two generations are heavy (masses greater than tens of TeV) while the superpartners of the third generation is relatively light (masses a few hundred GeV) \[10\]. Thus, the ESUSY with R parity conservation cannot account for the possible extra contribution to the \(g - 2\) of muon, reported recently \[1\], since smuonneutrino(\(\tilde{\nu}_\mu\)) and smuon (\(\tilde{\mu}\)) masses are very heavy > 20 TeV in ESUSY.

In this paper, we consider both the flavor changing problem and the BNL experiment \[1\] seriously, and study the anomalous magnetic moment of \(\mu\) in the ESUSY. Here, we do not digress into how the ESUSY phenomenon results at low energy, but mainly discuss its effects on the \((g - 2)_\mu\). In the effective SUSY scenario, sparticle masses of the first two generations are required to be heavier than 20 TeV \[10\]. These sparticles of the first two generations are not responsible for a large anomalous magnetic moment of \(\mu\). Therefore, the SUSY loops considered in Ref. \[2\] for \(a_\mu\) are negligible in ESUSY. For the gauge hierarchy solution, however, there is an effective SUSY for the third generation particles. The phenomenological constraints on the third generation sparticle masses are expected to be lighter than 1 TeV \[10\]. Within the experimentally allowed regions of these mass parameters, we look for a possibility of generating a large anomalous magnetic moment of \(\mu\) by introducing R parity violating interactions.

The ESUSY requires that top squarks must be lighter than 1 TeV for a hierarchy solution and \(SU(2) \times U(1)\) breaking by the top quark Yukawa coupling \[1\]: the left-handed stop
mass $M_{\tilde{t},L} < 1$ TeV and the right-handed stop mass $M_{\tilde{t},R} < 1$ TeV. Because of the $SU(2)$ symmetry left-handed sbottom mass $M_{\tilde{b},L}$ must be smaller than 1 TeV, while right-handed sbottom mass $M_{\tilde{b},R}$ need not to be constrained to $< 1$ TeV. Similarly, $M_{\tilde{\nu},L}, M_{\tilde{\tau},L}$ and $M_{\tilde{\tau},R}$ are not restricted. However, the Fayet-Iliopoulos anomaly should lead to a light Higgs boson below 20 TeV for the $SU(2) \times U(1)$ breaking, i.e. there remains an effective SUSY for the third generation below 20 TeV, which gives a condition $^{10}$

$$\text{Tr } Y M_{\text{heavy}}^2 = 0,$$  \hspace{1cm} (2)

where $Y$ is the electroweak hypercharge. If the full third generation sparticles are light, the above trace is can be satisfied under the assumption that the first two generations satisfy the trace rule. In addition, note that if the third generation scalar leptons are heavy with the same mass and the third generation scalar quarks are light then Rule (2) can also be satisfied, which however is not considered below. Therefore, for simplicity and in search of a possibility, we assume that the full third generation masses are light with mass less than 1 TeV.

In the estimation of the anomalous magnetic moment, there can be a mixing between the left scalar and the right scalar fermions, which however gives a contribution negligible compared to the case without the mixing. The main reason for this is that the standard model is a chiral theory and scalars introduced in the supersymmetrization remembers this chiral nature. So if there is a mixing term between left- and right-sfermion masses, then there must be a chirality flip by a mass insertion in the fermion line in the loop. However, the $SU(3) \times SU(2) \times U(1)$ symmetry forbids the mass insertion for charged leptons and quarks as an internal fermion line. The only allowed diagram for the internal neutrino line, which is possible only by introducing neutrino masses beyond the standard model, is as shown in Fig. 1 (b). We can anticipate Fig. 1 (a), but it is absent due to the illegal $M_{\tilde{\nu}}$ insertion in the supersymmetric limit. (Note however that there will be a small correction if we consider the $A$ term arising from $\chi_{333}$ coupling which must be very small in view of a small $\nu_\tau$ mass.)
Since neutrino masses are below 10 eV, we can neglect Fig. 1 (b) also and hence in the remainder of this paper we will neglect the mass mixing between the left- and right-handed sfermions. Also, for the mixing between squarks, there does not appear an important contribution.

The relevant $R$ parity violating superpotential($\lambda$ and $\lambda'$ couplings) is given by

$$W = \frac{1}{2}\lambda_{ijk}L_iL_jE^c_k + \lambda'_{ijk}L_iQ_jD^c_k,$$

where $i, j,$ and $k$ are generation indices, $L_i$ and $Q_i$ are the left-handed lepton and quark doublet superfields, and $E^c_i$ and $D^c_i$ are the left-handed anti-lepton and anti-quark singlet superfields. The bosonic symmetry implies, $\lambda_{ijk} = -\lambda_{jik}$.

The needed vertices for the $(g-2)_\mu$ calculation can be read from the above superpotential,

$$-\mathcal{L}_Y = \lambda_{ijk}(\bar{\nu}_iL\bar{e}_jL\bar{e}_kL + \bar{e}_jL\nu_iL\bar{e}_kL + \bar{e}^*_kR\nu_iL\bar{e}_jL)$$

$$+ \lambda'_{ijk}(\bar{\nu}_iL\bar{d}_jL\bar{d}_kL + \bar{d}_jL\nu_iL\bar{d}_kL + \bar{d}^*_kR\nu_iL\bar{d}_jL - \bar{e}_iL\bar{u}_jL\bar{d}_kL - \bar{\nu}_jL\bar{e}_iL\bar{d}_kL - \bar{\nu}_jL\bar{d}_iL\bar{d}_kL - \bar{d}^*_kR\bar{e}_iL\bar{u}_jL)$$

$$+ h.c.$$

(4)

In terms of the four component spinors, we obtain

$$-\mathcal{L}_Y = \lambda_{ij2}(\bar{\nu}_{iL}\bar{\mu}_LP_L\bar{e}_j + \bar{e}_{jL}\bar{\mu}_P\nu_i) + \lambda'_{i2k}(\bar{\nu}_{iL}\bar{\mu}_P\nu_k + \bar{\nu}_{kR}\bar{\mu}_P\bar{\nu}^c_i)$$
where the left-handed and right-handed chirality projecton operator are defined as $P_{L,R} = \frac{1}{2}(1 \mp \gamma^5)$.

As commented above, the slepton mixing diagrams, Fig. 1, are not important for $(g-2)_\mu$.

The $R$ violating $\lambda$ couplings contribute to the muon anomalous magnetic moment as shown in Fig. 2, (a) – (d). In Figs. 2, one of the external lines shown as left-handed muon is understood as shifted to a right-handed one by $m_\mu$ insertion in that external line. Then, we obtain,

\[ a_{\mu}^{\nu(1)} = \frac{m_\mu^2}{16\pi^2} |\lambda_{32k}|^2 \int_0^1 dx \frac{x^2 - x^3}{m_\mu x^2 + (m_{\tilde{\tau}_k}^2 - m_\mu^2)x + m_{\tilde{\nu}_L}^2(1 - x)} \]

\[ \approx \frac{|\lambda_{32k}|^2}{48\pi^2} \frac{m_\mu^2}{m_{\tilde{\nu}_L}^2}, \]

\[ a_{\mu}^{\tau R} = \frac{m_\mu^2}{16\pi^2} |\lambda_{i2k}|^2 \int_0^1 dx \frac{x^3 - x^2}{m_\mu x^2 + (m_{\tilde{\tau}_k}^2 - m_\mu^2)x + m_{\tilde{\nu}_L}^2(1 - x)} \]

\[ \approx -\frac{|\lambda_{i2k}|^2}{96\pi^2} \frac{m_\mu^2}{m_{\tilde{\tau}_R}^2}, \]

\[ a_{\mu}^{\nu(2)} = \frac{m_\mu^2}{16\pi^2} |\lambda_{ij2}|^2 \int_0^1 dx \frac{x^2 - x^3}{m_\mu x^2 + (m_{\tilde{\tau}_j}^2 - m_\mu^2)x + m_{\tilde{\nu}_L}^2(1 - x)} \]

\[ \approx \frac{|\lambda_{ij2}|^2}{48\pi^2} \frac{m_\mu^2}{m_{\tilde{\nu}_L}^2}, \]

\[ a_{\mu}^{\tau L} = \frac{m_\mu^2}{16\pi^2} |\lambda_{ij2}|^2 \int_0^1 dx \frac{x^3 - x^2}{m_\mu x^2 + (m_{\tilde{\tau}_j}^2 - m_\mu^2)x + m_{\tilde{\nu}_L}^2(1 - x)} \]

\[ \approx -\frac{|\lambda_{ij2}|^2}{96\pi^2} \frac{m_\mu^2}{m_{\tilde{\tau}_L}^2}. \]

Therefore, the $\lambda$ contribution to the muon anomalous magnetic moment is

\[ a_\mu^\lambda = \frac{m_\mu^2}{96\pi^2} \left( |\lambda_{32k}|^2 \frac{2}{m_{\tilde{\nu}_\tau}} + |\lambda_{3k2}|^2 \left[ \frac{2}{m_{\tilde{\tau}_\nu}} - \frac{1}{m_{\tilde{\tau}_L}} \right] - \frac{|\lambda_{k23}|^2}{m_{\tilde{\tau}_R}^2} \right) \]

where the sum over the repeated indices is implied.
The contribution from the $\lambda$ couplings. One $\mu_L$ (or $\mu_R$) is changed to $\mu_R$ (or $\mu_L$) by $m_\mu$ insertion.

The $\lambda$ couplings also contribute to the muon anomalous magnetic moment. The effect of soft supersymmetry breaking term can be considered, but we find that they are not important for our purpose. In addition, the constraint from the flavor changing effects restrict possible models, and here we concentrate on one allowed model as discussed in the introduction, the ESUSY model; thus the allowed couplings are rather restricted. Then, we obtain \[12\]

$$a_{\tilde{t}_L}^{(1)} = 3 \left( \frac{1}{3} \right) \frac{m_\mu^2}{16\pi^2} |\lambda'_{2jk}|^2 \int_0^1 dx \frac{x^2 - x^3}{m^2_\mu x^2 + (m^2_{d_k} - m^2_\mu) x + m^2_{u_jL} (1 - x)}$$

$$\approx \frac{|\lambda'_{2jk}|^2 m_\mu^2}{48\pi^2} \frac{m^2_{\tilde{t}_L}}{m^2_{u_jL}}$$

(11)
\[
\tilde{a}_{\mu}^{(1)} = 3 \left( \frac{1}{3} \right) \frac{m_{\mu}^2}{16\pi^2} |\lambda'_{jk}|^2 \int_0^1 dx \frac{x^3 - x^2}{m_{\mu}^2 x^2 + (m_{d_{kR}}^2 - m_{\mu}^2)x + m_{u_j}(1-x)} \\
\simeq -\frac{|\lambda'_{2j3}|^2}{96\pi^2} \frac{m_{\mu}^2}{m_{u_j}^2 - m_{u_j}^2} (1 + F_1(x_{u_j d_{kR}})) \tag{12}
\]
\[
\tilde{a}_{\mu}^{(2)} = 3 \left( \frac{2}{3} \right) \frac{m_{\mu}^2}{16\pi^2} |\lambda'_{2jk}|^2 \int_0^1 dx \frac{x^3 - x^2}{m_{\mu}^2 x^2 + (m_{u_j}^2 - m_{\mu}^2)x + m_{d_{kR}}^2(1-x)} \\
\simeq -\frac{|\lambda'_{2j3}|^2}{24\pi^2} \frac{m_{\mu}^2}{m_{u_j}^2 - m_{d_{kR}}^2} (1 + F_2(x_{u_j d_{kR}})) \tag{13}
\]
\[
a_{\mu}^{i_L} = 3 \left( \frac{2}{3} \right) \frac{m_{\mu}^2}{16\pi^2} |\lambda'_{2jk}|^2 \int_0^1 dx \frac{x^3 - x^2}{m_{\mu}^2 x^2 + (m_{u_j}^2 - m_{\mu}^2)x + m_{d_{kR}}^2(1-x)} \\
\simeq -\frac{|\lambda'_{2j3}|^2}{48\pi^2} \frac{m_{\mu}^2}{m_{u_j}^2} \tag{14}
\]

where

\[
F_1(x) = \frac{6x}{1-x} \left[ \frac{1}{2} - \frac{1}{1-x} - \frac{x}{(1-x)^2} \ln x \right], \tag{15}
\]
\[
F_2(x) = \frac{3x}{1-x} \left[ \frac{1}{2} + \frac{1}{1-x} + \frac{1}{(1-x)^2} \ln x \right], \tag{16}
\]

and \(x_{u_j d_{kR}} = m_{u_j}^2 / m_{d_{kR}}^2\). These \(F\) functions are important only for the top quark, i.e. for \(x_t \equiv x_{u_3 d_{kR}}\). Note that \(a_{\mu}^{i_L(1)}\) and \(a_{\mu}^{i_L(2)}\) add up to zero. Thus, the remaining Eqs. (12) and (13) give the \(\lambda'\) contribution,

\[
a_{\mu}^{\lambda'} = \frac{m_{\mu}^2}{32\pi^2} \left\{ \frac{1}{m_{b_{kR}}^2 - m_{\mu}^2} |\lambda'_{2j3}|^2 \left( 1 + \frac{2x_t}{1-x_t} \left[ \frac{1}{2} + \frac{3}{1-x_t} + \frac{2 + x_t}{(1-x_t)^2} \ln x_t \right] \right) \right\} \\
+ \frac{m_{\mu}^2}{32\pi^2} \left\{ \frac{1}{m_{b_{kR}}^2} \left( |\lambda'_{213}|^2 + |\lambda'_{223}|^2 \right) \right\} > 0. \tag{17}
\]

Note that this contribution is positive definite.
Fig. 3. The contribution from $\lambda'$ couplings. One $\mu_L$ is changed to $\mu_R$ by $m_\mu$ insertion.

We can show from the above expression that in the supersymmetric limit, i.e. $M_t = m_t$, $M_b = m_b$, etc, the contribution to $(g - 2)_\mu$ vanishes, which is a necessary consistency check. In the expression (17), the unknown mass is $m_{b_R}$. Without the left- and right-sbottom mixing and the flavor mixing, it is the physical mass.

For completeness, however, we present the formula with the left- and right-sbottom mixing. We need to consider a $2 \times 2$ sbottom mass matrix, in the $(\tilde{b}_L, \tilde{b}_R^*)^T$ basis,

$$M^2 = \begin{pmatrix} M_{11}^2 & M_{12}^2 \\ M_{21}^2 & M_{22}^2 \end{pmatrix}. \tag{18}$$

After diagonalization of $M^2$, $(M^2)^{-1}_{22}$ should replace $1/m_{b_R}^2$,

$$\left(V^\dagger (M_D^2)^{-1} V\right)_{22} = \sum_k \frac{|V_{k2}|^2}{M_k^2} = \frac{1}{m_{b_R}^2} \left[ 1 - \frac{m_b^2 (A_b + \mu \tan \beta)^2}{m_{b_R}^2 m_{b_L}^2} \right]^{-1} \tag{19}$$

where $V$ is the diagonalizing unitary matrix, $M_D^2$ is the diagonalized mass matrix, and the last relation is given for the MSSM case.

Therefore, we obtain the total contribution to $\Delta a_\mu$ in the ESUSY with $R$ parity violating couplings as

$$\Delta a_\mu^R = a_\mu^\lambda + a_\mu^{\lambda'} \tag{20}$$

where $a_\mu^\lambda$ and $a_\mu^{\lambda'}$ are given in Eqs. (10) and (17), respectively. In principle, the anomalous magnetic moment can take both signs. But for the most range of the parameter space,
in view of Eq. (17), $\Delta a_\mu$ is positive. The possibility of a negative sign arises when the contributions of Fig. 2 (b) and (d) dominate.

![Graph](image)

Fig. 4. The $1\sigma$ and $2\sigma$ band of $\Delta a_\mu$ in the $\lambda_{322}$ and $\tilde{m} \equiv m_{\tilde{\nu}} = m_\tau$ plane.

The low mass cut is the LEP bound of 81.0 GeV on the stau mass [13].

In view of the Brookhaven data of the anomalous magnetic moment of muon with the $2\sigma$ deviation [1], we obtain the necessary lower bound on the $R$ parity violating couplings ($\lambda$, $\lambda'$) from Eqs. (14) and (17) as

$$|\lambda| (\text{or } |\lambda'|) > 0.524 \times \left(\frac{\tilde{m}}{100\text{GeV}}\right)$$

where $\tilde{m}$ is the third generation sfermion mass (or the sbottom mass). In that case, the $\lambda'$ couplings cannot give sizable effects for the muon anomaly since there are unavoidable single bounds. $\lambda'_{233}$ is constrained from the radiative muon neutrino mass, $m_{\nu_\mu} \sim [3|\lambda'_{233}|^2 m^2_b(A_b + \mu \tan \beta)]/(8\pi^2 m^2_{b_R})$ as $|\lambda'_{233}| < 0.15 \times \sqrt{(m_{b_R}/100\text{GeV})}$ [14,15]. And we also have single bounds such as $|\lambda'_{213}| < 0.059 \times (m_{b_R}/100\text{GeV})$ from $R_\pi = \Gamma(\pi \rightarrow e\nu)/\Gamma(\pi \rightarrow \mu\nu)$ and
\[ |\lambda_{23}| < 0.21 \times (m_{\tilde{\nu}_{R}}/100\text{GeV}) \] from the decay process \( D \to K\ell\nu \) [14].

For the \( \lambda \) couplings, on the other hand, we can get a naturally large contribution to the muon anomalous magnetic moment in the effective SUSY model we are considering. Let us enumerate the \( \lambda \) contribution from Eq. (10) as

\[
a_\mu^\lambda \approx \frac{m_\mu^2}{32\pi^2 m^2} \left( \frac{2}{3} |\lambda_{321}|^2 + |\lambda_{322}|^2 + \frac{1}{3} |\lambda_{323}|^2 + \frac{1}{3} |\lambda_{312}|^2 - \frac{1}{3} |\lambda_{123}|^2 \right)
\]

(22)

where we set \( m_{\tilde{\nu}_r} = m_{\tilde{\nu}_\ell} = m_{\tilde{\tau}_R} \equiv \tilde{m} \). There also exist relevant single bounds on the \( \lambda \) couplings. Firstly, \( \lambda_{123} \) giving a negative contribution becomes negligible since \( |\lambda_{123}| < 0.05 \times (m_{\tilde{\tau}_R}/100\text{GeV}) \) from the charged current universality [14] and thus the \( \lambda \) contributions can be regarded as positive definite values consistent with the Brookhaven data if the other \( \lambda \) couplings are allowed to be larger than \( |\lambda_{123}| \). And we can also ignore \( \lambda_{312} \) and \( \lambda_{321} \), which just give contributions less than required by factor 10 since they are constrained as \( \sqrt{|\lambda_{312}|^2 + |\lambda_{321}|^2} < 0.25 \times (m_{\tilde{\nu}_r}/100\text{GeV}) \) from asymmetries in \( e^+e^- \) collisions at the Z peak [14]. Lastly, since

\[
\sqrt{|\lambda_{322}|^2 + \left( \frac{\tilde{M}}{\tilde{m}} \right)^2 |\lambda_{323}|^2} < 0.070 \times \left( \frac{\tilde{M}}{100\text{GeV}} \right)
\]

(23)

from \( R_\tau = \Gamma(\tau \to e\nu\bar{\nu})/\Gamma(\tau \to \mu\nu\bar{\nu}) \) or \( R_{\tau\mu} = \Gamma(\tau \to \mu\nu\bar{\nu})/\Gamma(\mu \to e\nu\bar{\nu}) \) [14], \( |\lambda_{323}| < 0.070 \times (\tilde{m}/100\text{GeV}) \) becomes negligible while \( |\lambda_{322}| \) can be as large as order one for \( \tilde{M} \equiv m_{\tilde{\mu}} = 20 \text{ TeV} \), which is a natural assumption in the effective SUSY.

Consequently, it is experimentally viable to take the \( \lambda_{322} \) dominant case among \( R \) parity violating couplings such that

\[
\Delta a_\mu^R \approx 34.9 \times 10^{-10} \left( \frac{100\text{GeV}}{\tilde{m}} \right)^2 |\lambda_{322}|^2.
\]

(24)

In Fig. 4, the above equation is plotted in the \( |\lambda_{322}| \) and \( \tilde{m} \) plane. Here we show the 1\( \sigma \)(dark grey) and 2\( \sigma \)(light grey) bands beyond the standard model prediction [7]. It is shown that most of the parameter space is consistent with the experimental bound (23) for the decoupling assumption \( \tilde{M} = 20 \text{ TeV} \). Note that there exists a enough room for perturbative \( \lambda_{322} \) to explain the BNL experiment even for a relatively large slepton mass in the third
generation, for example $\tilde{m} \sim 500$ GeV can give a perturbative $\lambda_{322}$ coupling within the shaded region. If we consider the other terms in Eqs. (10) and (17) in the phenomenologically allowed region with positive contributions, the $R$ parity violating couplings can be even lowered. [In this regard, we note that the SUSY calculation of $a_\mu$ with $\tan \beta \sim 40$ amounts to the Yukawa couplings of order 1.]

In conclusion, we have shown that the effective supersymmetric model toward a flavor conserving phenomena for the first two light generations and the introduction of $R$ parity violating couplings open up the possibility of a large anomalous magnetic moment of $\mu$ within the accessible range of the BNL experiment and a discovery potential of the third generation sfermions in the TevatronII and LHC accelerators.

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