Messenger-matter mixing and lepton flavor violation

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Abstract

We consider the Minimal Gauge Mediated Model (MGMM) with messenger-matter mixing and find that existing experimental limits on $\mu \to e\gamma$ decay and $\mu - e$ - conversion place significant constraints on relevant coupling constants. On the other hand, the rates of $\tau \to e\gamma$ and $\tau \to \mu\gamma$ in MGMM are well below the existing limits. We also point out the possibility of sizeable slepton oscillations in this model.

1. Presently, much attention is paid to lepton flavor physics in supersymmetric theories. Lepton flavor violation naturally emerges in those SUSY models where supersymmetry is broken by supergravity interactions. The corresponding soft-breaking terms are often assumed to be universal at the Planck scale. This universality breaks down due to the renormalization group evolution below the scale of Grand Unification. In this way one gets sizeable mixing in slepton sector at low energies \cite{1}. This mixing leads to lepton flavor violation for ordinary leptons ($\mu \to e\gamma$, etc.) at rates close to existing experimental limits \cite{1}, and also to slepton oscillations \cite{2} possibly observable at the Next Linear Collider.

Another class of SUSY models assumes gauge mediated supersymmetry breaking \cite{3}. In these models, the gauge interactions do not lead to lepton flavor violation because the messenger-matter interactions are flavor blind. Nevertheless there is a way to introduce flavor changing lepton interactions in these theories. This possibility is based on the observation that some of the messenger fields have the same quantum numbers as some of the usual fields. So it becomes natural to introduce direct mixing between messenger and matter fields \cite{4}. In such variation of the gauge mediated models, messengers not only transfer SUSY breaking to usual matter, but also generate lepton flavor violation.

The purpose of this paper is to show that in a reasonable range of parameters, messenger-matter mixing in minimal gauge-mediated models (MGMM) \cite{5} gives rise to observable rates of rare lepton flavor violating processes like $\mu \to e\gamma$ and $\mu - e$ conversion.

We will see that the tree level mixing between leptons is small due to seesaw type suppression. The tree level mixing between sleptons is also small in the part of the parameter space that is natural for MGMM. However, we observe that radiative corrections involving interactions with ordinary Higgs sector induce much stronger mixing of sleptons. The point is that messengers obtain masses not through interactions with ordinary Higgs fields, but through interactions with hidden sector. So, the overall matrix of trilinear couplings is
not proportional to the mass matrix (unlike in the Standard Model). In the basis of eigenstates of the tree level mass matrix, the largest trilinear terms involve sleptons as well as messenger and Higgs superfields. It is these terms that produce slepton mixing at the one loop level.

As a result, at messenger masses of order $10^5$ GeV and Yukawa couplings of order $10^{-2}$, the rates of $\mu \rightarrow e\gamma$ and $\mu - e$ conversion are comparable to their experimental limits.

2. The MGMM contains, in addition to MSSM particles, two messenger multiplets $Q = (q, \psi_q)$ and $\bar{Q} = (\bar{q}, \psi_{\bar{q}})$ which belong to 5 and $\bar{5}$ representations of $SU(5)$. In what follows we are interested in lepton sector so we will consider only colorless components of messenger multiplets. These multiplets couple to a MSSM singlet $Z$ through the superpotential term

$$\mathcal{L}_{ms} = \lambda Z Q \bar{Q}$$

$Z$ obtains non-vanishing vacuum expectation values $F$ and $S$ via hidden-sector interactions,

$$Z = S + F\theta\theta$$

Gauginos and scalar particles of MSSM get masses in one and two loops respectively. Their values are

$$M_{\lambda_i} = c_1 \frac{\alpha_i}{4\pi} \Lambda f_1(x)$$

for gauginos and

$$\tilde{m}^2 = 2\Lambda^2 \left[ C_3 \left( \frac{\alpha_3}{4\pi} \right)^2 + C_2 \left( \frac{\alpha_2}{4\pi} \right)^2 + \frac{5}{3} \left( \frac{Y}{2} \right)^2 \left( \frac{\alpha_1}{4\pi} \right)^2 \right] f_2(x)$$

for scalars. Here $\alpha_1 = \alpha / \cos^2 \theta_W$; $c_1 = 5/3$, $c_2 = c_3 = 1$; $C_3 = 4/3$ for color triplets (zero for singlets), $C_2 = 3/4$ for weak doublets (zero for singlets), $Y$ is the weak hypercharge. The two parameters entering eqs. (2) and (3) are:

$$\Lambda = \frac{F}{S}$$

and

$$x = \frac{\lambda F}{\lambda^2 S^2}$$

The dependence of masses on $x$ is very mild, as the functions $f_1(x)$ and $f_2(x)$ are close to 1. In the absence of mixing between messengers and leptons (and/or quarks), the predictions of this theory at realistic energies are determined predominantly by the value of $\Lambda$, while the value of $x$ is almost unimportant. It has been argued in ref. that $\Lambda$ must be larger than 70 TeV, otherwise the theory would generically be inconsistent with mass limits from LEP. The characteristic features of the model are large $\tan \beta$ (an estimate of ref. is $\tan \beta > 48$) and large squark masses. The next lightest supersymmetric particle (NLSP) is argued to be a combination of $\tilde{\tau}_R$ and $\tilde{\tau}_L$. Bino is slightly heavier, but lighter than $\tilde{\mu}_{L,R}$ and $\tilde{e}_{L,R}$. 

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Unlike the masses of MSSM particles, the masses of messenger fields strongly depend on $x$. Namely, the vacuum expectation value of $Z$ mixes scalar components of messenger fields and gives them masses

$$M^2_\pm = \frac{\Lambda^2}{x^2}(1 \pm x)$$

It is clear that $x$ must be smaller than 1. Masses of fermionic components of messenger superfields are all equal to $\Lambda x$.

Components of $Q$ have the same quantum numbers as left leptons (and $\bar{d}$ quarks), so one can introduce messenger-matter mixing

$$\mathcal{L}_{mm} = H_D L \tilde{Y}_{ij} E_j$$

where

$$H_D = (h_D, \chi_D)$$

is one of the Higgs doublet superfields,

$$L_i = (\tilde{l}_i, l_i) = \left\{ \begin{array}{c} (\tilde{e}_{iL}, e_{iL}) \\
(\tilde{q}, \psi_q) \end{array} \right\} ; \ i = 1, \ldots, 3$$

are left doublet superfields and

$$E_j = (\bar{e}_{Rj}, e_{Rj})$$

are right lepton singlet superfields. Hereafter $\tilde{i}, \tilde{j} = 1, \ldots, 4$ label the three left lepton generations and the messenger field, $i, j = 1, \ldots, 3$ correspond to the three leptons and $Y_{\tilde{i}j}$ is the $4 \times 3$ matrix of Yukawa couplings,

$$Y_{\tilde{i}j} = \left( \begin{array}{ccc}
Y_e & 0 & 0 \\
0 & Y_\mu & 0 \\
0 & 0 & Y_\tau \\
Y_{\tilde{4}1} & Y_{\tilde{4}2} & Y_{\tilde{4}3}
\end{array} \right)$$

In terms of component fields, the tree level scalar potential and Yukawa terms are

$$V = \lambda S^2 q^* q + \mu^2 h_D h_D^* + |\lambda S \bar{q} + h_D \tilde{e}_{Rj} Y_{4j}|^2 + |\mu h_U + Y_{\tilde{i}j} \tilde{e}_{Rj}|^2 + |Y_{ij} h_D \tilde{e}_{Rj}|^2 + |Y_{\tilde{i}j} \tilde{l}_i h_D|^2 + \left( \lambda S \psi_q \bar{\psi}_q - \lambda F \bar{q} \bar{q} \right) + Y_{\tilde{i}j} (h_D \tilde{e}_{Rj} + \chi_D \tilde{l}_i \tilde{e}_{Rj} + \chi_D \tilde{t}_i \tilde{e}_{Rj}) + \mu \chi_D \chi_U + h.c. \right)$$

where $\mu$ is the usual parameter of MSSM and

$$H_U = (h_U, \chi_U)$$

is the second Higgs superfield. Besides these terms, there are soft-breaking terms coming from loops involving messenger fields. In the absence of messenger-matter mixing they have the form (at the SUSY breaking scale, which is of order of $\Lambda$)

$$\mathcal{L}_{sb} = \tilde{m}^2_{Li} \tilde{l}_i \tilde{l}_i^* + \tilde{m}^2_{Ri} \tilde{e}_{Ri} \tilde{e}_{Ri}^* + B \mu h_U h_D$$

where
where \( \tilde{m}_{Lj}^2 \) and \( \tilde{m}_{Rj}^2 \) are given by eq.(3). Messenger-matter mixing modifies eq.(3); we will consider this modification later on. We will not discuss sneutrinos in what follows, so we will use the notation \( e_{Lj} \) for charged components of lepton doublets.

The main emphasis of this paper is lepton flavor violation induced by the messenger-matter mixing, eq.(3). However, let us note in passing that another effect of this mixing is the absence of heavy stable charged (and colored) particles (messengers) in the theory [1].

3. To see that lepton mixing is small at the tree level, let us write the fermion mass matrix (that includes left and right fermionic messengers) in the following form

\[
\mathcal{M}_f = \mathcal{U}_f \mathcal{D}_f \mathcal{U}_f^\dagger
\]

where

\[
\mathcal{U}_f = \begin{pmatrix}
1 & 0 & 0 & -\frac{y_e y_L^*}{\lambda S} \\
0 & 1 & 0 & -\frac{y_\mu y_L^*}{\lambda S} \\
\frac{y_\mu y_e}{\lambda S} & \frac{y_e y_\mu}{\lambda S} & \frac{y_\tau y_\mu}{\lambda S} & 1
\end{pmatrix}
\]

and

\[
\mathcal{D}_f = \begin{pmatrix}
1 - \frac{|y_1|^2}{\lambda S} & -\frac{y_1 y_2}{\lambda S} & -\frac{y_1 y_3}{\lambda S} & -\frac{y_1}{\lambda S} \\
-\frac{y_1 y_2}{\lambda S} & 1 - \frac{|y_2|^2}{\lambda S} & -\frac{y_2 y_3}{\lambda S} & -\frac{y_2}{\lambda S} \\
-\frac{y_1 y_3}{\lambda S} & -\frac{y_2 y_3}{\lambda S} & 1 - \frac{|y_3|^2}{\lambda S} & -\frac{y_3}{\lambda S} \\
\frac{|y_1|^2}{\lambda S} & \frac{|y_2|^2}{\lambda S} & \frac{|y_3|^2}{\lambda S} & 1 - \frac{|y_1|^2 + |y_2|^2 + |y_3|^2}{\lambda S}
\end{pmatrix}
\]

are mixing matrices to the leading order in \( \frac{\lambda}{\lambda S} \). Here \( v_U \) and \( v_D \) are Higgs expectation values,

\[
y_i = Y_i v_D, \quad y_{e,\mu,\tau} = Y_{e,\mu,\tau} v_D
\]

and

\[
\mathcal{D}_f = \text{diag}\left(y_e (1 - \frac{|y_e|^2}{\lambda S}), y_\mu (1 - \frac{|y_\mu|^2}{\lambda S}), y_\tau (1 - \frac{|y_\tau|^2}{\lambda S}), \lambda S (1 + \frac{|y_e|^2}{\lambda S}), \lambda S (1 + \frac{|y_\mu|^2}{\lambda S}), \lambda S (1 + \frac{|y_\tau|^2}{\lambda S})\right)
\]

is the matrix of mass eigenvalues. In principle, the off-diagonal terms in eqs.(8) and (9) may lead to lepton flavor violation due to one loop diagrams involving sleptons and gauginos [1]. However, these mixing terms are negligible due to see-saw type mechanism: in MGMM one definitely has \( \lambda S > 10^4 \) GeV, \( \tan \beta > 1 \) and even at \( Y_i \sim 1 \) the mixing terms are smaller than \( 10^{-4} \). The corresponding contributions to lepton flavor violating rates are too small to be observable.

Mixing in slepton sector is also small at the tree level. Indeed, the mass term of scalars has the following form

\[
\mathcal{V}_{sc} = s \mathcal{M}_{sc} s^\dagger
\]

where

\[
s = (\tilde{e}_L, \tilde{\mu}_L, \tilde{e}_R, \tilde{\mu}_R, \tilde{\tau}_L, \tilde{\tau}_R, q, \bar{q}^*)
\]
are the scalar fields, and

\[
\mathcal{M}_{sc} = \begin{pmatrix}
\tilde{m}_{e_L}^2 & 0 & 0 & \mu Y_e v_U & 0 & 0 & \bar{y}_e y_1^* & 0 \\
0 & \tilde{m}_{\mu_L}^2 & 0 & 0 & \mu Y_\mu v_U & 0 & \bar{y}_\tau y_2^* & 0 \\
0 & 0 & \tilde{m}_{\tau_L}^2 & 0 & 0 & \mu Y_\tau v_U & \bar{y}_\tau y_3^* & 0 \\
\mu Y_e v_U & 0 & 0 & \tilde{m}_{e_R}^2 & y_1^* y_2 & y_1^* y_3 & \mu Y^*_e v_U & \lambda S y_1^* \\
0 & \mu Y_\mu v_U & 0 & y_1 y_2^* & \tilde{m}_{\mu_R}^2 & y_2^* y_3 & \mu Y^*_\mu v_U & \lambda S y_2^* \\
0 & 0 & \mu Y_\tau v_U & y_1 y_3^* & y_2 y_3^* & \tilde{m}_{\tau_R}^2 & \mu Y^*_\tau v_U & \lambda S y_3^* \\
y_1 y_2 & y_\mu y_2 & y_\tau y_3 & \mu Y_1 v_U & \mu Y_2 v_U & \mu Y_3 v_U & \lambda^2 S^2 - \lambda F & -\lambda F \\
0 & 0 & 0 & \lambda S y_1 & \lambda S y_2 & \lambda S y_3 & -\lambda F & \lambda^2 S^2
\end{pmatrix}
\]

After the scalar messengers are integrated out at the tree level, the lepton flavor violating terms in the mass matrix of right sleptons are of order \(Y_i Y_j (v^2 D^2 \Lambda^2 + \mu^2 v^2 U^2 \Lambda^2)\). These terms are smaller than the one loop contributions (see below) at \(\Lambda \gtrsim 10\) TeV, which is certainly the case in MGMM. Flavor violating mixing between left sleptons is even smaller. The only substantial non-diagonal terms in this matrix are \((-\lambda F)\) in the messenger sector and \(\tilde{\tau}_R - \tilde{\tau}_L\) mixing. Due to the latter, the NLSP is likely to be \(\tilde{\tau}\).

4. The most substantial mixing in slepton sector is due to one loop diagrams coming from trilinear terms in the superpotential, that involve \(H_D\). The fact that the one loop mixing terms of sleptons are proportional to the large parameter \(\Lambda^2\) is obvious from eq.(5): say, one of the cubic terms in the scalar potential, \([\lambda SY_{4i}q^* h_D \tilde{e}_{Rj} + \text{h.c.}]\), contains \(\lambda S = \frac{\Lambda}{x}\) explicitly.

After diagonalizing the messenger mass matrix we obtain the diagrams contributing to slepton mixing to the order \((\lambda S)^2\), which are shown in fig.1.

![Figure 1: The diagrams giving main contribution to the scalar mixing matrix.](image)

If supersymmetry were unbroken, their sum would be equal to zero. In our case of broken supersymmetry the resulting contribution to the mass matrix of sleptons, to the order \(\Lambda^2\) is

\[
\delta m_{ij}^2 = -\frac{1}{8\pi^2} \frac{\Lambda^2}{x^2} \left\{ -\ln(1-x^2) - \frac{x}{2} \ln \left( \frac{1+x}{1-x} \right) \right\} Y^*_i Y_{4j}
\]

Since sleptons get negative shifts in \(\tilde{m}_{ij}^2\), this equation immediately implies theoretical bounds on Yukawa couplings \(Y_{4i}\) which come from the requirement \(\tilde{\tau}\) that none of the slepton masses become negative (see below).
5. Let us now consider the effect of slepton mixing on the usual leptons. We begin with $\mu \rightarrow e\gamma$. The dominant contribution to the amplitude of this decay comes from diagrams shown in fig. 2, where $N_n$ are neutralino mass eigenstates with masses $M_n$.

![Figure 2: Diagrams contributing to $\mu \rightarrow e\gamma$ decay](image)

In general, their contribution gives

$$\Gamma(\mu \rightarrow e\gamma) = \frac{\alpha}{4} m_\mu^3 |F|^2 \quad (13)$$

$$F = \frac{\alpha}{4 \pi \cos^2 \theta_W} m_\mu \delta m_{12}^2 G(\tilde{m}_{eR}^2) \quad (14)$$

where

$$G(m^2) = \sum_{k=1}^{4} \frac{H_{bk}^2}{M_k^2} g \left( \frac{m^2}{M_k^2} \right)$$

$$g(r) = \frac{1}{6(r-1)^5} \left( 17 - 9r - 9r^2 + r^3 + 6(3r + 1) \ln r \right)$$

Here $H_{bk}$ are coefficients in the decomposition of bino in terms of mass eigenstates. To obtain numerical estimates we point out that neutralino mixing is small, at least at large mass of bino and large $\mu$. Neglecting this mixing, recalling eqs. (2) and (3) and collecting all factors, we find

$$\Gamma(\mu \rightarrow e\gamma) = 2.6 \cdot 10^{10} \frac{m_\mu^5}{\Lambda^4} |Y_{41}Y_{42}|^2 \phi(x) \quad (15)$$

where

$$\phi(x) = \frac{1}{f_1^8} \left[ g \left( \frac{6 f_2}{5 f_1^2} \right) \right]^2 \left[ \ln (1 - x^2) + \frac{3}{2} \ln \frac{1 + x}{1 - x} \right]^2 x^4 \quad (16)$$

and the functions $f_1(x)$ and $f_2(x)$ can be found in refs. [7, 8]. This rate strongly depends on $x$, or, in physical terms, on the messenger masses. In particular, at small $x$

$$\phi(x) = \frac{1}{36} \left[ g \left( \frac{6}{5} \right) \right]^2 x^4 = 4.2 \cdot 10^{-5} x^4$$

To see what eq. (17) means, we plot the limit on the product of Yukawa couplings as function of $x$ at $\Lambda = 100$ TeV in fig. 3, where we make use of the existing experimental limit on the rate of $\mu \rightarrow e\gamma$ decay [9]. The case of arbitrary $\Lambda$ is

\[ \text{We neglect here the running of slepton masses from the SUSY breaking scale $\Lambda$ down to weak scale. In fact, this running is small.} \]
Figure 3: Upper limit on $Y = |Y_{41}Y_{42}|^{1/2}$ as function of $x$ at $\Lambda = 100$ TeV (solid line). Above the dashed line slepton oscillations are unsuppressed at maximal mixing.

straightforward: as follows from eq. (15), the limit on $|Y_{41}Y_{42}|^{1/2}$ scales like $\Lambda$ at fixed $x$ and varying $\Lambda$. We conclude that $\mu \to e\gamma$ decay is observable in this model in a reasonable part of the parameter space.

The slepton mixing in this model gives rise also to $\mu - e$ - conversion. The dominant contribution to $\Gamma(\mu \to e\gamma)$ is given by penguin-type diagrams, while box diagrams are suppressed by squark masses. So, there is a simple relation between $\mu - e$ - conversion and $\mu \to e\gamma$ rates [1]:

$$\Gamma(\mu \to e) = 16\alpha^4 Z_{eff}^4 Z|F(q)|^2 \Gamma(\mu \to e\gamma)$$

(17)

For $\text{Ti}_{22}^{48}$ with $Z = 22$, $Z_{eff} = 17.6$, $|F(q)| = 0.54$ [10] one expects

$$\frac{\Gamma(\mu \to e)}{\Gamma(\mu \to e\gamma)} = 2.8 \cdot 10^{-2}$$

while the ratio of experimental limits [1] is

$$\frac{\Gamma(\mu \to e)_{\text{exp}}^{\text{lim}}}{\Gamma(\mu \to e\gamma)_{\text{exp}}^{\text{lim}}} = 1.1 \cdot 10^{-1}$$

Hence, the existing limit on $\mu - e$ - conversion gives weaker (by a factor 1.4) bounds on the product of Yukawa couplings $|Y_{41}Y_{42}|^{1/2}$.

Similar mechanism leads also to flavor changing $\tau$-decays, $\tau \to e\gamma$ and $\tau \to \mu\gamma$. Crude estimates of the rates are straightforward to obtain by neglecting $\tilde{\tau}_R - \tilde{\tau}_L$ mixing. In this approximation, $\tau$-decay rates are given by eq. (15) with obvious substitutions $m_\mu \to m_\tau$, $Y_{42} \to Y_{43}$ (and $Y_{41} \to Y_{42}$ in the case of
Figure 4: The upper limits on matrix $Y_{ij}$ from flavor changing $\tau$ - decays. The dashed line is the limit on $|Y_{41}Y_{43}|^{1/2}$ from $\tau \to e\gamma$ decay, the same line represents also the limit on $|Y_{42}Y_{43}|^{1/2}$ from $\tau \to \mu\gamma$ decay. The solid line corresponds to theoretical constraint, which is the same for $|Y_{41}Y_{43}|^{1/2}$ and $|Y_{42}Y_{43}|^{1/2}$.

$\tau \to \mu\gamma$). The corresponding upper bounds derived from experimental limits on $\Gamma(\tau \to \mu\gamma)$ and $\Gamma(\tau \to e\gamma)$ [3] are shown in figure [4]. In fact, these bounds are weaker than theoretical constraints inherent in this model. Indeed, with loop corrections [12] to slepton mass matrix included, its eigenvalues $\tilde{m}_i^2$ are all positive only if

$$|Y_{41}|^2 + |Y_{42}|^2 + |Y_{43}|^2 < \frac{5}{3} \alpha f_2(x) \frac{x^2}{\ln \frac{1-x}{1+x} - \ln(1-x^2)} \quad (18)$$

Making use of the inequality

$$|Y_{41}Y_{43}| < \frac{1}{2} (|Y_{41}|^2 + |Y_{42}|^2 + |Y_{43}|^2)$$

one obtains from eq. (18) theoretical constraint on $|Y_{41}Y_{43}|^{1/2}$. Precisely the same constraint applies to $|Y_{12}Y_{43}|^{1/2}$. The result is shown in fig. [4]. We see that self-consistence of the model requires that the rates $\Gamma(\tau \to e\gamma)$ and $\Gamma(\tau \to \mu\gamma)$ are lower than the present experimental limits at least by factor $10^{-2}$.

6. Pair production of right sleptons (which decay into leptons and bino) in $e^+ - e^-$ annihilation at the Next Linear Collider will result in acoplanar $\mu^+ - \mu^-$ an $e^+ - e^-$ events with missing energy. The NLSP in the model under discussion is $\tilde{\tau}$, so there will be also four $\tau$-leptons produced in each event (two $\tau$ will come from bino decays into $\tau$ and $\tilde{\tau}$ and two more from subsequent $\tilde{\tau}$
decays). In the presence of slepton mixing, the slepton oscillations leading to lepton flavor violating $\mu^\pm - e^\mp$ events, are possible \[2\].

Oscillations of sleptons are characterized by the mixing angle, which in this model can be found from eq.(12)

$$\tan 2\phi = 2\frac{|Y_{41}Y_{42}|}{|Y_{41}|^2 - |Y_{42}|^2}$$

If $Y_{41} \sim Y_{42}$ then mixing is close to maximal. The cross section of $e^+e^- \rightarrow e^\pm\mu^\mp + 4\tau$ may, however, be suppressed even at large mixing if the life-time of $\mu_R$ and $e_R$ is small compared to the period of oscillations. The condition for the absence of such suppression is \[2\]:

$$2\Gamma M_{e_R} < |M^2_{e_R} - M^2_{\mu_R}|$$

where $M_{e_R}$ and $M_{\mu_R}$ denote the true slepton masses. For slepton decay width $\Gamma$ we have

$$\Gamma = \frac{\alpha_1}{2} \frac{m_{\tilde{R}}}{\mu} \left(1 - \frac{M^2_{\text{bino}}}{m^2_{\tilde{R}}}\right)^2$$

By making use of eqs.(2),(3) and (12), it is straightforward to translate the condition (19) into a condition imposed on Yukawa couplings $Y_{41}$ and $Y_{42}$. For example, at small $x$ one has

$$(|Y_{41}|^2 + |Y_{42}|^2)x^2 > 1.5 \cdot 10^{-6}$$

In the case of maximal mixing, $Y_{41} = Y_{42}$, the region of validity of eq.(19) is shown in fig.3. Therefore, there is a fairly wide range of parameters in which $\mu \rightarrow e\gamma$ - decay and slepton oscillations are both allowed. Note, that unlike $\mu \rightarrow e\gamma$, slepton oscillation parameters $\sin 2\phi$ and $\frac{2\Gamma M_{e_R}}{|M^2_{e_R} - M^2_{\mu_R}|}$ are independent of $\Lambda$.

7. Messenger-matter mixing is possible also in strongly interacting sector, as some messengers carry quantum numbers of right down-quarks. It will lead to flavor changing processes involving ordinary quarks \[4\]. Yet another possibility emerges in modifications of the model that are obtained by changing the messenger sector. For example, one can consider messengers belonging to 10 and $\bar{10}$ of $SU(5)$. An interesting feature of this model is the presence of messengers with quantum numbers of up-quarks. So one can introduce mixing of the type $\lambdaQLD$ and suggest an interpretation of possible HERA leptoquark \[12\] as a scalar messenger. However, to have its mass of the order of 200 GeV one has to set $x$ very close to 1, which implies fine tuning

$$\frac{(\lambda S)^2 - \lambda F}{(\lambda S)^2} \sim 10^{-6}$$

We are not aware of any mechanism that would make such fine tuning natural, but the very possibility of messenger interpretation of HERA events seems to be interesting.

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References

[1] R.Barbieri and L.Hall  Phys. Lett. B338 (1994) 212; R.Barbieri, L.Hall, A.Strumia  Nucl. Phys. B445 (1995) 219.
[2] N.V.Krasnikov  Phys. Lett. B388 (1996) 783.
[3] M.Dine and A.Nelson,  Phys. Rev. D47 (1993) 1277; M.Dine, A.Nelson and Y.Shirman,  Phys. Rev. D51 (1995) 1362; M.Dine, A.Nelson,Y.Nir and Y.Shirman,  Phys. Rev. D53 (1996) 2658.
[4] M.Dine, Y.Nir and Y.Shirman,  Phys. Rev. D55 (1997) 1501.
[5] K.S.Babu, C.Kolda and F.Wilczek,  Phys. Rev. Lett. 77 (1996) 3070.
[6] F.Borzumati, hep-ph/9702307.
[7] S.Dimopoulos, G.F.Giudice and A.Pomarol,  Phys. Lett. B389 (1996) 37.
[8] S.Martin,  Phys. Rev. D55 (1997) 3177.
[9] R.D.Bolton et al.,  Phys. Rev. D38 (1988) 2077.
[10] J.Bernabeu, E.Nardi and D.Tommasini,  Nucl. Phys. B409 (1993) 69.
[11] K.W.Edwards et al.,  Phys. Rev. D55 (1997) 3919.
[12] C.Adolf et al., DESY 97-024, hep-ex/9702012; J.Breitweg et al., DESY 97-025, hep-ex/9702013.