Probing the Seesaw and Gauge Mediation Scales with \( \text{BR}(\mu \to e\gamma) \) and \( |U_{e3}| \)

Daniel Grossman\(^\dagger\) and Yosef Nir\(^\dagger\)
Department of Particle Physics and Astrophysics, Weizmann Institute of Science, Rehovot 76100, Israel

The new MEG bound on \( \text{BR}(\mu \to e\gamma) \) provides the strongest upper bound on the scale of gauge mediation of supersymmetry breaking. If, in the future, this decay is observed by MEG, the mediation scale will become known to within one order of magnitude, and the seesaw scale will be constrained. In such a case, contributions from Planck mediated supersymmetry breaking are likely to be non-negligible, and an interpretation in terms of purely seesaw parameters will be impossible. The recent evidence for \( |U_{e3}| \sim 0.15 \) further sharpens the predictions of gauge mediated supersymmetry breaking.

**Introduction.** The MEG experiment has recently announced a new bound \(^1\),
\[
\text{BR}(\mu \to e\gamma) \leq 2.4 \times 10^{-12},
\]
a factor of 5 improvement over the previous MEGA bound \(^2\). Within a few years, MEG is expected to explore the range of
\[
\text{BR}(\mu \to e\gamma) \gtrsim 10^{-13}. \tag{2}
\]

The Standard Model, extended with the seesaw mechanism to account for neutrino masses, predicts a branching ratio that is about forty orders of magnitude lower. Thus, future observation of \( \text{BR}(\mu \to e\gamma) \) within the range of \( 2.4 \times 10^{-12} - 10^{-13} \) will signal new physics. A leading candidate to account for such a rate will be the supersymmetric standard model, where lepton flavor violation arises, in general, in the soft supersymmetry breaking slepton mass-squared terms. The implications will be particularly interesting in the framework of models with universal soft terms, such as gauge mediated supersymmetry breaking (GMSB), where the supersymmetric flavor problem is naturally solved. In such a framework, lepton flavor violation can still arise via the renormalization group evolution (RGE) effect of the singlet neutrinos on the soft breaking terms \(^3\), \(^4\). (For reviews, see for example \(^5\), \(^6\).)

An observation of \( \mu \to e\gamma \), when interpreted within the GMSB framework, will put a lower bound on the seesaw scale. The lower the seesaw scale, the smaller the neutrino Yukawa couplings, leading eventually to negligible RGE effects and to a decay rate below the experimental sensitivity. A GMSB interpretation of such an observation will also put an upper bound on the seesaw scale. It must be lower than the mediation scale in order for the RGE effects to take place. The mediation scale itself is bounded from above, or else Planck scale mediation (PMSB), where the flavor structure is not expected to be universal, becomes significant. Conversely, an upper bound on the rate of \( \mu \to e\gamma \) provides an upper bound on the scale of gauge mediation, to avoid large contributions from PMSB, and an upper bound on the seesaw scale, if it is lower than the mediation scale.

The purpose of this work is to obtain the constraints on the seesaw and mediation scales that follow from the new bound \((1)\). We further study the constraints that will follow from observing \( \mu \to e\gamma \), to understand whether such an observation can be translated into constraints on the seesaw flavor parameters, and to explore further phenomenological implications.

**Low energy neutrino parameters.** The GMSB predictions for lepton flavor violation depend also on the low energy neutrino parameters. Specifically, the dependence is on the three following combinations:
\[
\begin{align*}
\mu_{\mu e} &= s_{12} c_{12} c_{23} (m_2 - m_1) + s_{13} s_{23} (m_3 - m_1), \\
\mu_{\tau e} &= -s_{12} c_{12} s_{23} (m_2 - m_1) + s_{13} c_{23} (m_3 - m_1), \\
\mu_{\tau\mu} &= c_{23} s_{23} [-c_{12} (m_2 - m_1) + (m_3 - m_1)],
\end{align*}
\]
where \( s_{ij} \equiv \sin \theta_{ij}, c_{ij} \equiv \cos \theta_{ij} \), with \( \theta_{ij} \) the angles of the leptonic mixing matrix \( U \), and we use \( c_{13} = 1 \). To evaluate the numerical values of the \( \mu_{ij} \) parameters, we assume normal hierarchy for the neutrino masses, whereby \( m_2 - m_1 \simeq \sqrt{\Delta m^2_{31}} \) and \( m_3 - m_1 \simeq \sqrt{\Delta m^2_{31}} \) \(^7\):
\[
\begin{align*}
m_2 - m_1 &= 0.009 \text{ eV}, \\
m_3 - m_1 &= 0.05 \text{ eV}.
\end{align*}
\]

The recent intriguing measurements of \( |U_{e3}| \) by the T2K \(^8\) and MINOS \(^9\) experiments have been combined into a global analysis of oscillation data \(^10\), yielding
\[
s_{13} \simeq 0.15 \pm 0.03 \ (1\sigma). \tag{5}
\]

For the other two mixing angles, we use the central values from a recent fit \(^11\):
\[
s_{12} = 0.56, \quad s_{23} = 0.68. \tag{6}
\]

We obtain (taking into account only the \( s_{13} \)-related uncertainty):
\[
\begin{align*}
\mu_{\mu e} &= 0.0082 \pm 0.0010 \text{ eV}, \\
\mu_{\tau e} &= 0.0027 \pm 0.0011 \text{ eV}, \\
\mu_{\tau\mu} &= 0.022 \text{ eV}. \tag{7}
\end{align*}
\]
The Model. In minimal GMSB models \cite{12,15} (for earlier work, see \cite{14,18}), slepton-doublet masses are given, at the messenger scale, by

$$m_L^2 \equiv m_L^2(M_m) = \frac{3n_5(5\alpha_2^2 + \alpha_1^2)M^2_2}{16\pi^2},$$

where $M_m$ is the mass scale of the messenger fields (that is, the mediation scale), $M_4$ is a scale of order a (few) hundred TeV, and $n_5$ is the number of messenger fields in the $5\times5$ representation of $SU(5)$. (In this work, we focus on models with weakly coupled messenger fields. We will explore more general models of gauge mediation \cite{19} in future work.)

The leptonic part of the superpotential reads

$$W_L = E_i^c \lambda_i^c L_i H_1 + N_i^c \lambda_i^c L_i H_2 - \frac{1}{2} M_{2i}^c L_i N_j^c,$$

where $L_i$, $E_i^c$, and $N_j^c$ are lepton superfields in the $(2)-1/2$, $(1)_{+1}$, and $(1)_0$ representations of $SU(2) \times U(1)$, respectively, and $H_{1,2}$ are the Higgs superfields in the $(2)_{1/1,2}$ representations. The matrices $\lambda_c$ and $\nu_c$ are complex $3 \times 3$ Yukawa matrices, while $M_N$ is the symmetric Majorana mass matrix. Without loss of generality, we work in a basis where $\lambda_c$ and $M_N$ are diagonal.

The light neutrino Majorana mass matrix is given by

$$M_{\nu} = v_L^{-1} \lambda^T \nu N^{-1} \lambda \nu,$$

where $v_i = (H_i^0)$. Diagonalization of $M_{\nu}$ leads to the mass eigenvalues $m_i$, $i = 1, 2, 3$, and to the leptonic mixing matrix $U$.

If some of the eigenvalues of $M_{\nu}$ are lower than $M_m$, then RGE of $m_{ij}^2$ between the scales $M_m$ and $M_N$ will generate off-diagonal terms in $m_{ij}^2$. In a generic supersymmetric model, the RGE is given by \cite{20,21}

$$\mu \frac{d}{d\mu} (m_{ij}^2) = \frac{d}{d\mu} \left( m_{ij}^2 \right)_{\text{MSSM}}$$

$$+ \frac{1}{16\pi^2} \left[ (m_{ij}^2 \lambda_i \lambda_j + \lambda_i^\dagger \lambda_j m_{ij}^2)_{ij} + 2(\lambda_i^\dagger m_N^2 \lambda_j + m_{ij}^2 \lambda_i \lambda_j + A_{ij}^L A_{ij}^R)_{ij} \right].$$

In GMSB models, the trilinear scalar couplings are negligible, $A_{ij} = 0$, and the soft masses-squared of singlet fields vanish, $m_S = 0$. Moreover, at the messenger scale $m_{ij}^2 = m_{ij}^{\text{grav}}$. For simplicity, we adopt from here on the minimal lepton flavor violation (MLFV) ansatz of Ref. \cite{22} and take the singlet neutrinos to be degenerate,

$$M_N = m_N \mathbf{1}.$$

In the leading log approximation, and with the GMSB boundary conditions \cite{3}, the off-diagonal elements of the doublet slepton mass matrix at low energy are given by

$$m_{ij}^2 = \frac{\tilde{m}_L^2}{4\pi^2} (\lambda_i^\dagger \lambda_j)_{ij} \ln \frac{M_m}{m_N}.$$
The $\mu \to e\gamma$ decay. Within our framework, where the main features of the spectrum are determined by GMSB with a perturbative messenger sector, the $\mu \to e\gamma$ decay is dominated by the chargino/sneutrino loop diagrams. An approximate expression for the branching ratio can be written (based on Ref. [21]) as follows:

$$\text{BR}(\mu \to e\gamma) \simeq \frac{6\pi\alpha_\mu^2\beta^4}{\hat{m}_L^2} \left[ \left( \frac{m_2}{m_L} \right)^2, \left( \frac{\mu}{m_L} \right)^2 \right] \delta^2_{\mu e},$$

(20)

where $t_\beta \equiv \tan \beta$, $M_2$ is the Wino mass, $\mu$ is the Higgsino mass term, $m_{\tilde{L}}$ is the average mass of the quasi-degenerate sneutrinos,

$$\delta_{\mu e} \equiv \frac{(m^2_{\tilde{L}})_{21} - \hat{\delta}_{\mu e}}{m^2_{\tilde{L}}} = \frac{\hat{\delta}_{\mu e}}{m^2_{\tilde{L}}},$$

(21)

and

$$h(x, y) = \frac{xy}{(x - y)^2} \left[ g(x) - g(y) \right]^2,$$

$$g(x) = -5 + 4x + x^2 - 2 \ln x - 4x \ln x.$$

(22)

We note that for very large $n_5$, the approximation of Eq. (20) becomes less accurate, as the contributions from neutralino/charged-slepton loop diagrams become non-negligible. The bounds of Eqs. (28) and (30) below are derived from numerical calculations (based on Ref. [29]), which do take into account these additional contributions.

An upper bound on the radiative decay,

$$\text{BR} (\mu \to e\gamma) < \frac{\text{BR}}{\text{R}},$$

(23)

implies

$$\delta^L_{\mu e} < 6 \cdot 10^{-5} \left( \frac{20}{t_\beta} \right) \sqrt{\frac{0.1}{h}} \left( \frac{m_{\tilde{L}}}{300 \text{ GeV}} \right)^2 \sqrt{\frac{\text{BR}}{2.4 \cdot 10^{-12}}}.$$

(24)

Using Eqs. (17) and (8), we then obtain an upper bound on the scale of gauge mediation (for anarchical PMSB slepton mass-squared terms):

$$\frac{M_m}{M_{Pl}} < 3 \cdot 10^{-5} \left( \frac{\text{BR}}{2.4 \cdot 10^{-12}} \right)^{1/4} \left( \frac{\hat{m}_{\tilde{L}}}{300 \text{ GeV}} \right) \times \left( \frac{m_{\tilde{L}}}{m_{\tilde{L}}} \right)^2 \left( \frac{20n_5}{t_\beta} \right)^{1/2} \left( \frac{0.1}{h} \right)^{1/4}. $$

(25)

The properties of the GMSB spectrum [30], and in particular the fact that the ratio between the wino and doublet-slepton masses increases with $n_5$, lead to

$$h \sim 0.1 \to 0.01 \text{ for } n_5 = 1 \to 20.$$

(26)

The running of $m_{\tilde{L}}$ is dominated by gaugino masses, such that for fixed $\hat{m}_{\tilde{L}}$

$$m_{\tilde{L}} \approx 1.01 + \frac{9n_5}{20\pi^2},$$

(27)

Plugging these results back into Eq. (25) we obtain for the two extreme cases of $n_5 = 1, 20$:

$$\frac{M_m}{M_{Pl}} \lesssim \begin{cases} 3 \times 10^{-5} & n_5 = 1, \\ 1 \times 10^{-3} & n_5 = 20. \end{cases}$$

(28)

These bounds scale like $\sqrt{\frac{20}{t_\beta} \left( \frac{\hat{m}_{\tilde{L}}}{300 \text{ GeV}} \right) \left( \frac{\text{BR}}{2.4 \cdot 10^{-12}} \right)^{1/4}}$.

They are to be compared with Eq. (18). We conclude that the new MEG bound on $\mu \to e\gamma$ provides the strongest upper bound on the scale of gauge mediation.

Inserting Eq. (15) into the expression (20), we obtain, for the purely GMSB contributions and for $M_m > m_N$,

$$\text{BR}(\mu \to e\gamma) \simeq \frac{3\alpha^3 h}{2\pi^3 s^2_{\beta\gamma} s_W^2} \left[ m_{\tilde{L}}^4 m_N^2 \delta_{\mu e} \left( \frac{\ln M_m}{m_N} \right)^2 \right]$$

$$= 5 \times 10^{-13} \left( \frac{m_{\tilde{L}}}{m_L} \right)^8 \left( \frac{0.1}{s_{\beta\gamma}} \right)^2 \left( \frac{h}{0.1} \right) \times \left( \frac{m_N}{10^{12} \text{ GeV}} \right)^2 \left( \frac{\mu_{\mu e}}{0.008 \text{ eV}} \right)^2$$

$$\times \left( \frac{300 \text{ GeV}}{\hat{m}_{\tilde{L}}} \right)^4 \left( \frac{\ln \frac{M_m}{m_N}}{4} \right)^2.$$

Using the expression (20), the upper bound of Eq. (11) implies that

$$m_N \gtrsim M_m \text{ or } m_N \lesssim \begin{cases} 2.7 \times 10^{12} \text{ GeV} & n_5 = 1, \\ 1.1 \times 10^{14} \text{ GeV} & n_5 = 20, \end{cases}$$

(30)

where the bounds scale like $(\sin 2\beta/0.1)(\hat{m}_{\tilde{L}}/300 \text{ GeV})^2$.

In the literature, the less well-motivated case of universal PMSB (known as mSUGRA or CMSSM) is often considered [31,32]. The supersymmetric spectrum is different from the GMSB framework, and the approximation (20) is replaced by

$$\text{BR}(\mu \to e\gamma) \simeq \frac{8\pi\alpha^2 t_\beta^4 \delta_{\mu e}^4}{7\hat{m}_{\tilde{L}}^4}.$$  

(31)

where $\hat{m}$ is the scale of the soft breaking parameters. The upper bound on BR($\mu \to e\gamma$) of Eq. (11) provides an upper bound on $\delta_{\mu e}$:

$$\delta_{\mu e} \lesssim 1.4 \times 10^{-4} \left( \frac{20}{t_\beta} \right) \left( \frac{\hat{m}_{\tilde{L}}/300 \text{ GeV}}{2.4 \cdot 10^{-12}} \right)^{1/4}.$$

(32)

An important feature of this scenario is that the seesaw scale, $m_N$, is necessarily below the mediation scale, $M_{Pl}$. Thus, the generation of off-diagonal entries in the doublet slepton mass-squared matrix by RGE from $M_{Pl}$ to $m_N$ is unavoidable. Consequently, Eq. (11) provides an upper bound on $m_N$. Using Eq. (15) with $M_m$ replaced by $M_{Pl}$, Eq. (32) yields, for $\tan \beta = 10$ and $\hat{m} = 300 \text{ GeV},$

$$m_N \lesssim 1.4 \times 10^{12} \text{ GeV}.$$  

(33)
This bound scales like \((\sin 2\beta/0.2)\). In the parameter region of interest, the bound scales roughly as \((\tilde{n}/300 \text{ GeV})^{2/17}\). (For related work on seesaw parameters in the PMSB framework, see [34, 38].) Consequences of observing \(\mu \to e\gamma\). If \(\mu \to e\gamma\) is actually observed by MEG, this will have far reaching consequences for GMSB. Explicitly, if \(\text{BR}(\mu \to e\gamma)\) is established to be within the range of Eq. (2), we will be able to make the following statements:

- The following lower bounds apply:
  \[
  M_m \gtrsim 2 \times 10^{13} \text{ GeV or } M_m > m_N \gtrsim 3 \times 10^{11} \text{ GeV.}
  \] (34)
  Thus, low scale gauge mediation will be excluded.

- At the same time, a bound similar to (and perhaps somewhat stronger than) Eq. (25) applies. Thus, the range of \(M_m\) will be determined to within about an order of magnitude.

- Combining these considerations leads us to suspect that \(x_{\mu e}\) of Eq. (19) cannot be negligibly small. For example, taking \(M_m \sim 5 \times 10^{13} \text{ GeV and } m_N \sim 10^{12} \text{ GeV, gives } x_{\mu e} \sim 1\).

  Thus, in case that \(\mu \to e\gamma\) is observed and interpreted within the GMSB framework, we should not neglect PMSB contributions to \(\delta_{\mu e}\). In particular, in Eq. (20), we should replace
  \[
  \delta_{\mu e}^2 \to (\delta_{\mu e} + \delta_{\mu e}^{L, \text{grav}})^2 + (1/16)(\alpha_1/\alpha_2)^2(\delta_{\mu e}^{R, \text{grav}})^2,
  \] (35)
  where \(\delta_{\mu e}\) stands for the pure gauge-mediation contribution of Eq. (13), while \(\delta_{\mu e}^{L, \text{grav}}\) and \(\delta_{\mu e}^{R, \text{grav}}\) stand for the contributions from gravity mediation to, respectively, \((m_L^{2, 21})\) and \((m_R^{2, 21})\), both of which we take to be estimated by Eq. (17). Within our framework, the contribution from \(\delta_{\mu e}^{R, \text{grav}}\) is thus of order one percent of that from \(\delta_{\mu e}^{L, \text{grav}}\), and can be safely neglected for most purposes.

  The interplay between gauge-mediated and gravity-mediated contributions can be understood from Fig. 1. (A similar figure, which neglects, however, the PMSB contributions, was presented in Ref. [4].) We can make the following statements regarding an explanation of \(\text{BR}(\mu \to e\gamma) > 10^{-13}\) within our framework:

- For \(M_m \gtrsim 5 \times 10^{13} \text{ GeV, gravity mediated contributions are, in general, too large.}\)

- For \(10^{11} \text{ GeV} \lesssim m_N < M_m \lesssim 10^{14} \text{ GeV, gauge-mediated contributions can be large enough.}\)

  Given this situation, we can now make a more precise statement about \(x_{\mu e}\). For a given value of \(\text{BR}(\mu \to e\gamma)\), the \(x_{ij}\) parameters are minimized when \(\ln(M_m/m_N) = 3/4\) and \(M_m\) takes the minimal value consistent with the decay rate, which we denote by \(\hat{M}_m\). From Eq. (29) we obtain (neglecting the uncertainty in \(m_{\mu e}\))
  \[
  \frac{\hat{M}_m}{10^{13} \text{ GeV}} \simeq \sqrt{\frac{\text{BR}(\mu \to e\gamma) \times 20}{8 \times 10^{-13} \tan \beta (300 \text{ GeV})^2}}.
  \] (36)
  Putting this value in Eq. (19), we find
  \[
  x_{\mu e} \simeq 0.07 \sqrt{\frac{\text{BR}(\mu \to e\gamma) \times 20}{8 \times 10^{-13} n_5 \tan \beta (300 \text{ GeV})^2}}.
  \] (37)

  We remind the reader that this minimal value of \(x_{\mu e}^{\text{min}}\) is obtained under the assumption that the gravity mediated contribution to the slepton masses-squared is of order \((F_s/M_{P1})^2\). It can be further suppressed if this scale is accompanied with small or hierarchical dimensionless coefficients.

  In addition to setting a lower bound on \(x_{\mu e}\), a measurement of \(\mu \to e\gamma\) yields both upper and lower bounds on \(M_m\), as can be seen in Fig. 1. The upper bounds for the cases of \(n_5 = 1, 20\) are given in Eq. (28). To determine the lower bound we require \(\frac{\partial}{\partial m_N} \delta_{\mu e} = 0\). Along a curve of constant \(\text{BR}(\mu \to e\gamma)\) this occurs at \(\ln \frac{\hat{M}_m}{m_N} = 1\).

FIG. 1: Curves of fixed \(\text{BR}(\mu \to e\gamma) = 10^{-11}, 10^{-12}, 10^{-13}\) in the \(M_m - m_N\) plane.
(for which \(x_{\mu e}\) is larger than the bound given by Eq. [37]. The resulting lower bound is:

\[
\frac{M_m}{M_P} \gtrsim \begin{cases} 3 \times 10^{-6} & n_5 = 1, \\ 2 \times 10^{-4} & n_5 = 20, \end{cases}
\]  

(38)

where the bounds scale like \(\sqrt{\frac{\text{Br}(\mu \rightarrow e\gamma)}{2.4 \times 10^{-12} \left(\frac{m_\mu}{300 \text{ GeV}}\right)^2 \frac{20}{13}}\)}

We conclude that a measurement of \(\text{Br}(\mu \rightarrow e\gamma)\) will fix \(M_m\) to within one order of magnitude.

**Related observables.** The GMSB framework, combined with the MLFV seesaw mechanism, relate the \(\mu \rightarrow e\gamma\) decay rate to a number of other lower energy observables. In this section we review and update these relations in view of the new measurements.

1. \((g - 2)\) of the muon:
   The dependence of \(\text{Br}(\mu \rightarrow e\gamma)\) on the supersymmetric flavor-conserving parameters \(\tan \beta\) and \(m_{\tilde{E}_L}\), Eq. (20), can be eliminated by the use of the supersymmetric contribution to \((g - 2)\) of the muon:

\[
\text{Br}(\mu \rightarrow e\gamma) = 3 \times 10^{-5} \left(\frac{\delta_{\mu e}^{\text{SUSY}}}{10^{-9}}\right)^2 \delta_{\mu e}^2.
\]  

(39)

Future experimental and theoretical developments will decide whether it is more useful to use collider measurements of \(\tan \beta\) and \(m_{\tilde{E}_L}\), or the measurement of \((g - 2)\), to extract \(\delta_{\mu e}\) from \(\text{Br}(\mu \rightarrow e\gamma)\) and by that test the GMSB framework.

2. \(\mu \rightarrow e\) conversion:
   An interesting observable is the \(\mu \rightarrow e\) conversion rate \(R(\mu \rightarrow e)\). In the GMSB framework, the photon penguin diagram dominates both the radiative decay and the \(\mu \rightarrow e\) conversion, and thus the two rates are related in a way that is independent of the model parameters:

\[
\frac{R(\mu \rightarrow e \text{ in Ti})}{\text{Br}(\mu \rightarrow e\gamma)} \simeq 0.005.
\]  

(40)

Clearly, all the implications that we describe here for a signal in the \(\mu \rightarrow e\gamma\) decay can be made in a similar way from a signal of \(\mu \rightarrow e\) conversion.

3. Radiative \(\tau\) decays:
   Models with quasi-degenerate sleptons correlate all three radiative charged lepton decay rates in a simple way:

\[
\text{Br}(\tau \rightarrow \ell\gamma) = 0.175 \text{Br}(\mu \rightarrow e\gamma)(\delta_{\tau\ell}/\delta_{\mu e})^2.
\]  

(41)

GMSB in its version employed here translates Eq. (41) into an expression that depends only on light neutrino flavor parameters:

\[
\text{Br}(\tau \rightarrow \ell\gamma) = 0.175 \text{Br}(\mu \rightarrow e\gamma)(\mu_{\tau\ell}/\mu_{\mu e})^2.
\]  

(42)

For our two sets of mixing parameters, we get the following predictions:

\[
\begin{align*}
\frac{\text{Br}(\tau \rightarrow e\gamma)}{0.175 \text{Br}(\mu \rightarrow e\gamma)} & \simeq 0.04 - 0.17, \\
\frac{\text{Br}(\tau \rightarrow \mu\gamma)}{0.175 \text{Br}(\mu \rightarrow e\gamma)} & \simeq 5.4 - 9.0.
\end{align*}
\]  

(43)

Whether the gravity mediated contributions are significant depends on the \(x_{ij}\) parameters. Here, our framework gives

\[
\frac{x_{\tau\ell}}{x_{\mu e}} \sim \frac{\mu_{\mu e}}{\mu_{\tau\ell}}.
\]  

(44)

which leads to

\[
\begin{align*}
x_{\tau\ell}/x_{\mu e} & \simeq 2.2 - 4.0, \\
x_{\tau\mu}/x_{\mu e} & \simeq 0.33 - 0.42.
\end{align*}
\]  

(45)

We learn that our framework predicts \(\text{Br}(\tau \rightarrow \mu\gamma) \gtrsim \text{Br}(\tau \rightarrow e\gamma)\) and, furthermore, among the three rates of radiative decays, \(\tau \rightarrow \mu\gamma\) is the closest to represent the seesaw flavor parameters.

**Conclusions.** Within the framework of gauge mediated supersymmetry breaking, flavor violation is minimal. Yet, observable lepton flavor violation can be generated if the seesaw scale is lower than the scale of gauge mediation. In particular, a measurement of the \(\mu \rightarrow e\gamma\) decay rate is sensitive to the gauge-mediation and seesaw scales, and to the seesaw parameters.

In this work, we focus our attention on “ordinary gauge mediation”, with weakly coupled messenger fields, and on “minimally lepton flavor violating” seesaw sector, with degenerate singlet neutrinos and real Yukawa couplings. Many of our conclusions (such as the significance of Planck scale mediation) hold in more generality, but we postpone a detailed study of these generalizations to future work.

The recent MEG upper bound, \(\text{Br}(\mu \rightarrow e\gamma) \lesssim 2.4 \times 10^{-12}\), gives the strongest upper bound up-to-date on the mediation scale, which for a small number of messenger fields reads \(M_m \lesssim 10^{13}\) GeV [see Eq. (28)]. The seesaw scale is either higher than the gauge mediation scale or lower than about \(10^{13}\) GeV [see Eq. (30)]. The bound on the seesaw scale in the framework of gravity mediated supersymmetry breaking is of order \(10^{12}\) GeV.

If the \(\mu \rightarrow e\gamma\) decay is eventually observed by MEG, it will determine the scale of gauge mediation to within one order of magnitude. It will be impossible, however, to interpret the measurement in terms of purely seesaw parameters, because contributions from gravity mediation will be, in general, significant.

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[1] J. Adam et al. [MEG collaboration], Phys. Rev. Lett. 107, 171801 (2011) [arXiv:1107.5547 [hep-ex]].

[2] M. L. Brooks et al. [MEGA Collaboration], Phys. Rev. Lett. 83, 1521 (1999) [arXiv:hep-ex/9905013].

[3] F. Borzumati and A. Masiero, Phys. Rev. Lett. 57, 961 (1986).

[4] K. Tobe, J. D. Wells and T. Yanagida, Phys. Rev. D 69, 035010 (2004) [arXiv:hep-ph/0310148].

[5] A. Masiero, S. K. Vempati and O. Vives, J. Phys. G 37, 075021 (2010).

[6] K. Nakamura et al. [Particle Data Group], J. Phys. G37, 075021 (2010).

[7] K. Abe et al. [T2K Collaboration], Phys. Rev. Lett. 107, 041801 (2011) [arXiv:1006.5022 [hep-ex]].

[8] P. Adamson et al. [MINOS Collaboration], Phys. Rev. Lett. 107, 181802 (2011) [arXiv:1108.0142 [hep-ex]].

[9] G. L. Fogli, E. Lisi, A. Marrone, A. Palazzo, A. M. Rotunno, Phys. Rev. D84, 053007 (2011) [arXiv:1106.6028 [hep-ph]].

[10] M. Dine and A. E. Nelson, Phys. Rev. D 48, 1277 (1993) [hep-ph/9303230].

[11] M. Dine, A. E. Nelson and Y. Shirman, Phys. Rev. D 51, 1362 (1995) [hep-ph/9408384].

[12] M. Dine, A. E. Nelson, Y. Nir and Y. Shirman, Phys. Rev. D 53, 2658 (1996) [hep-ph/9507378].

[13] G. F. Giudice and R. Rattazzi, Phys. Rept. 322, 419 (1999) [hep-ph/9801271].

[14] M. Dine and W. Fischler, Phys. Lett. B 110, 227 (1982); Nucl. Phys. B 204, 346 (1982).

[15] C. R. Nappi and B. A. Ovrut, Phys. Lett. B 113, 175 (1982).

[16] L. Alvarez-Gaume, M. Claudson and M. B. Wise, Nucl. Phys. B 207, 96 (1982).

[17] P. Meade, N. Seiberg and D. Shih, Prog. Theor. Phys. Suppl. 177, 143 (2009) [arXiv:0801.3278 [hep-ph]].

[18] J. Hisano, T. Moroi, K. Tobe, M. Yamaguchi and T. Yanagida, Phys. Lett. B 357, 579 (1995) [arXiv:hep-ph/9501407].

[19] J. Hisano, T. Moroi, K. Tobe and M. Yamaguchi, Phys. Rev. D 53, 2442 (1996) [arXiv:hep-ph/9510309].

[20] Y. Cavioni, B. Grinstein, G. Isidori and M. B. Wise, Nucl. Phys. B 728, 121 (2005) [arXiv:hep-ph/0507001].

[21] G. Hiller, Y. Hochberg and Y. Nir, JHEP 0903, 115 (2009) [arXiv:0812.0511 [hep-ph]].

[22] G. Hiller, Y. Hochberg and Y. Nir, JHEP 1003, 079 (2010) [arXiv:1001.1513 [hep-ph]].

[23] C. D. Froggatt and H. B. Nielsen, Nucl. Phys. B 147, 277 (1979).

[24] S. Davidson and A. Ibarra, JHEP 0109, 013 (2001) [arXiv:0107005 [hep-ph]].

[25] J. Hisano, T. Moroi, K. Tobe, M. Yamaguchi and T. Yanagida, Phys. Lett. B 357, 579 (1995) [arXiv:hep-ph/9501407].

[26] J. Hisano, T. Moroi, K. Tobe and M. Yamaguchi, Phys. Rev. D 53, 2442 (1996) [arXiv:hep-ph/9510309].

[27] V. Cirigliano, B. Grinstein, G. Isidori and M. B. Wise, Nucl. Phys. B 728, 121 (2005) [arXiv:hep-ph/0507001].

[28] G. Hiller, Y. Hochberg and Y. Nir, JHEP 0903, 115 (2009) [arXiv:0812.0511 [hep-ph]].

[29] G. Hiller, Y. Hochberg and Y. Nir, JHEP 1003, 079 (2010) [arXiv:1001.1513 [hep-ph]].

[30] C. D. Froggatt and H. B. Nielsen, Nucl. Phys. B 147, 277 (1979).

[31] Y. Nir and N. Seiberg, Phys. Lett. B 309, 337 (1993) [arXiv:hep-ph/9304307].

[32] M. Leurer, Y. Nir and N. Seiberg, Nucl. Phys. B 420, 468 (1994) [arXiv:hep-ph/9310320].

[33] Y. Grossman and Y. Nir, Nucl. Phys. B 448, 30 (1995) [arXiv:hep-ph/9502418].

[34] S. C. Petcov, S. Profumo, T. Takamori and C. E. Yaguna, Nucl. Phys. B 708, 193 (2005) [arXiv:hep-ph/0506195].

[35] S. Dimopoulos, S. D. Thomas, J. D. Wells, Nucl. Phys. B 588, 39-91 (1997) [hep-th/9609454].

[36] J. Hisano, M. Nagai, P. Paradisi and Y. Shimizu, JHEP 0912, 030 (2009) [arXiv:0904.2080 [hep-ph]].

[37] P. Paradisi, JHEP 0510, 006 (2005) [arXiv:hep-ph/0505046].

[38] E. Arganda and M. J. Herrero, Phys. Rev. D 73, 055003 (2006) [arXiv:hep-ph/0510405].

[39] S. Davidson and A. Ibarra, JHEP 1009, 013 (2001) [arXiv:hep-ph/0104076].

[40] P. Deppisch, H. Pas, A. Redelbach, R. Ruckl and Y. Shimizu, Eur. Phys. J. C 28, 365 (2003) [hep-ph/0206122].

[41] A. Masiero, S. K. Vempati and O. Vives, Nucl. Phys. B 649, 189 (2003) [hep-ph/0209033].

[42] A. Ibarra and C. Simonetto, JHEP 0908, 113 (2009) [arXiv:0903.1776 [hep-ph]].

[43] S. Davidson and M. Elmer, Eur. Phys. J. C 71, 1804 (2011) [arXiv:1108.0548 [hep-ph]].

[44] J. Hisano and K. Tobe, Phys. Lett. B 510, 197 (2001) [arXiv:hep-ph/0102315].