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Probing supersymmetric model with heavy sfermions using leptonic flavor and CP violations

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ABSTRACT

We study leptonic flavor and CP violating observables in supersymmetric (SUSY) models with heavy sfermions, which is motivated by the recent results of the LHC experiments (i.e., the discovery of the Higgs-like boson with the mass of about 126 GeV and the negative searches for the superparticles). Even if the sfermion masses are of O(10–100 TeV), signals may be within the reach of future leptonic flavor and CP violation experiments assuming that the off-diagonal elements of the sfermion mass matrices are unsuppressed compared to the diagonal ones. We also consider the SUSY contribution to the K^0-$\bar{K}^0$ mixing parameters; we show that the leptonic observables can become as powerful as those in K^0-$\bar{K}^0$ mixing to constrain SUSY models.

Recent progresses of the searches for new particles at the LHC have provided important information about the physics at the electroweak scale and beyond. In particular, supersymmetry (SUSY), which is one of the important candidates of the physics beyond the standard model, has been seriously constrained by the results of the LHC, i.e., the discovery of the Higgs-like boson with the mass of about 126 GeV [1,2], and the negative searches for superparticles [3,4]. In particular, it is notable that the Higgs mass is preferred to close to the Z-boson mass in the minimal SUSY standard model (MSSM) and that the lightest Higgs mass of about 126 GeV is hardly realized in the MSSM unless stops are heavier than 10 TeV or the tri-linear scalar coupling constant of stops is enhanced.

These facts suggest a class of supersymmetric models, i.e., models with heavy sfermions. If the SUSY is broken with the SUSY breaking scale corresponding to $m_{3/2} \sim 10$–100 TeV (with $m_{3/2}$ being the gravitino mass), all the scalars in the MSSM sector (except for the lightest Higgs boson) naturally have a mass of O(10–100 TeV). Such a model is phenomenologically viable; the lightest Higgs mass can be pushed up to 126 GeV, while all the superparticles (in particular, sfermions) can be out of the reach of the LHC experiments. If the gravitino mass is as heavy as $m_{3/2} \sim 10$–100 TeV, it also helps to avoid the serious cosmological gravitino problem [5]. In addition, SUSY model with sfermion masses of O(10–100 TeV) is compatible with the grand unified theory (GUT). Thus, the scenario with heavy sfermions has been recently attracted attentions [6–14].

Heavy sfermions are also advantageous to avoid (or to relax) constraints from flavor and CP violations. In SUSY models in which the masses of superparticles are around 1 TeV, it is often the case that too large flavor and CP violations are induced by loop diagrams with superparticles inside the loop. With heavy enough sfermions, such constraints are supposed to be avoided. However, even if the sfermion masses are around 10–100 TeV, it is non-trivial whether the model evades all the flavor and CP constraints. It is well known that the constraint from the CP violation in the kaon decay (i.e., the constraint from the so-called ϵ_K parameter) often gives the most stringent constraint and that the SUSY contribution to ϵ_K may become larger than the standard-model prediction even with the sfermion masses of O(10–100 TeV) [15].

The purpose of this Letter is to reconsider the flavor and CP constraints on the SUSY model, paying particular attention to the heavy sfermion scenario. We will see that, in some class of well-motivated model, the constraint from ϵ_K is relaxed and that lepton flavor violation (LFV) and CP violation may be also powerful tools to study heavy sfermion scenario. Future experiments measuring $Br(\mu \to e\gamma)$, $Br(\mu \to 3e)$, μ-e conversion rate, and the electron electric dipole moment (EDM) d_e will cover the parameter region with the sfermion masses of O(10–100 TeV).

Let us start our discussion with introducing the framework of the model of our interest. In this Letter, we consider the case where SUSY is dynamically broken by the condensation of a chiral superfield Z. (There may exist more than one chiral superfields responsible for the SUSY breaking, but the following discussion is unaffected.) Allowing higher-dimensional operators suppressed by

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the Planck scale $M_{\rm Pl} \simeq 2.4 \times 10^{18}$ GeV, the Kähler potential may contain the following term

$$K \ni \frac{\kappa_{\Phi,IJ}}{M_{\rm Pl}^2} Z^\dagger Z \Phi_I^\dagger \Phi_J, \tag{1}$$

where Φ denotes chiral superfields in the MSSM sector, corresponding to $q_L(\mathbf{3},\mathbf{2},1/6)$, $u_R^c(\bar{\mathbf{3}},\mathbf{1},-2/3)$, $d_R^c(\bar{\mathbf{3}},\mathbf{1},1/3)$, $l_L(\mathbf{1},\mathbf{2},-1/2)$, and $e_R^c(\mathbf{1},\mathbf{1},1)$ (with the gauge quantum numbers for $SU(3)_C$, $SU(2)_L$ and $U(1)_Y$ being shown in the parenthesis), $\kappa_{\Phi,IJ}$ is a constant, and I and J are generation indices which run from 1 to 3. After the SUSY breaking, the F-component of Z acquires a vacuum expectation value, and the soft SUSY breaking mass squareds of sfermions in the MSSM sector show up. In this Letter, we assume that the SUSY is broken with relatively large value of the gravitino mass of $m_{3/2} \sim O(10\text{–}100$ TeV). Assuming that $\kappa_{\Phi,IJ} \sim O(0.1\text{–}1)$, sfermion masses (as well as the SUSY breaking Higgs mass parameters) are of the same order. In this framework, the Kähler potential may contain the terms like

$$K \ni c_1 H_u H_d + \frac{c_2}{M_{\rm Pl}^2} Z^\dagger Z H_u H_d + \text{h.c.}, \tag{2}$$

where H_u and H_d are up- and down-type Higgses, respectively. With c_1 and c_2 being $O(0.1\text{–}1)$, we expect that the SUSY invariant Higgs mass (so-called μ parameter) and the bi-linear SUSY breaking parameter (so-called B parameter) are both expected to be of the same order of the gravitino mass [16,17,12].

Even if the sfermion masses are of $O(10\text{–}100$ TeV), the gaugino masses are model-dependent. If the SUSY breaking sector contains a singlet field, then the gaugino masses may arise from the direct F-term interaction between the gauge multiplet and the SUSY breaking field. In such a case, we expect that the gaugino masses are of the order of the gravitino mass. We call such a case as heavy gaugino case, in which we assume that the gaugino masses M_A obey the simple GUT relation:

$$\frac{3}{5}\frac{M_1}{g_1^2} = \frac{M_2}{g_2^2} = \frac{M_3}{g_3^2}, \tag{3}$$

where g_A denote standard model gauge coupling constants. (Here, $A = 1$, 2 and 3 correspond to $U(1)_Y$, $SU(2)_L$ and $SU(3)_C$, respectively.) On the contrary, if there is no singlet field, effect of the anomaly-mediated SUSY breaking (AMSB) may dominate the gaugino masses. In such a case, the gaugino masses are obtained as [18,19][1]

$$M_A^{(\rm AMSB)} = -\frac{b_A g_A^2}{16\pi^2} m_{3/2}, \tag{4}$$

where b_A denote coefficients of the renormalization-group (RG) equations for g_A, i.e., $b_A = (-11, -1, 3)$.[2] In the following discussion, we consider both cases. We also comment here on the tri-linear scalar couplings (i.e., so-called A parameters); in the following discussion, their effects are not important, so are neglected.

We parameterize the soft SUSY breaking terms of sfermions in the gauge basis as

$$\mathcal{L}_{\rm soft} = \tilde{q}_{L,I}^\dagger \mathcal{M}_{\tilde{q}_L,IJ}^2 \tilde{q}_{L,J} + \tilde{u}_{R,I}^c \mathcal{M}_{\tilde{u}_R,IJ}^2 \tilde{u}_{R,J}^{c\dagger} + \tilde{d}_{R,I}^c \mathcal{M}_{\tilde{d}_R,IJ}^2 \tilde{d}_{R,J}^{c\dagger} + \tilde{l}_{L,I}^\dagger \mathcal{M}_{\tilde{l}_L,IJ}^2 \tilde{l}_{L,J} + \tilde{e}_{R,I}^c \mathcal{M}_{\tilde{e}_R,IJ}^2 \tilde{e}_{R,J}^{c\dagger}, \tag{5}$$

or in the flavor basis,

$$\mathcal{L}_{\rm soft} = \tilde{u}_{L,i}^\dagger \mathcal{M}_{\tilde{u}_L,ij}^2 \tilde{u}_{L,j} + \tilde{d}_{L,i}^\dagger \mathcal{M}_{\tilde{d}_L,ij}^2 \tilde{d}_{L,j} + \tilde{u}_{R,i}^c \mathcal{M}_{\tilde{u}_R,ij}^2 \tilde{u}_{R,j}^{c\dagger} + \tilde{d}_{R,i}^c \mathcal{M}_{\tilde{d}_R,ij}^2 \tilde{d}_{R,j}^{c\dagger} + \tilde{\nu}_{L,i}^\dagger \mathcal{M}_{\tilde{\nu}_L,ij}^2 \tilde{\nu}_{L,j} + \tilde{e}_{L,i}^\dagger \mathcal{M}_{\tilde{e}_L,ij}^2 \tilde{e}_{L,j} + \tilde{e}_{R,i}^c \mathcal{M}_{\tilde{e}_R,ij}^2 \tilde{e}_{R,j}^{c\dagger}. \tag{6}$$

(Here and hereafter, the tilde is for superpartners of the standard-model particles.) In Eq. (6), all the mass matrices are in the basis in which the fermion mass matrices are diagonalized; thus, $\mathcal{M}_{\tilde{u}_L}^2$ and $\mathcal{M}_{\tilde{d}_L}^2$ are related by the CKM matrix. It is obvious that the non-vanishing values of $\mathcal{M}_{\tilde{f},ij}^2$ (with $i \neq j$) become new sources of flavor violation. In addition, the off-diagonal elements of the mass matrices can have non-vanishing phases, and are new sources of CP violation.

For the following discussion, it is convenient to define

$$\Delta_{\tilde{f},ij} \equiv \frac{\mathcal{M}_{\tilde{f},ij}^2 - m_{\tilde{f}}^2 \delta_{ij}}{m_{\tilde{f}}^2}, \tag{7}$$

where

$$m_{\tilde{f}}^2 \equiv \mathcal{M}_{\tilde{f},11}^2. \tag{8}$$

The value of $\Delta_{\tilde{f},ij}$ is model-dependent. If all the higher-dimensional operators suppressed by the Planck scale are allowed, $\Delta_{\tilde{f},ij} \sim O(0.1)$ is expected.[3]

Now, we discuss the rates of LFV. As we have mentioned, non-vanishing values of $\Delta_{\tilde{e}_L,ij}$, $\Delta_{\tilde{e}_R,ij}$, and $\Delta_{\tilde{\nu}_L,ij}$ induce various LFV processes. With $\Delta_{\tilde{e}_L,ij}$, $\Delta_{\tilde{e}_R,ij}$, and $\Delta_{\tilde{\nu}_L,ij}$ being fixed, the LFV rates become smaller as the sleptons become heavier. To see the mass scale of sleptons accessible with the experiments, we simply assume the following structure of the slepton mass matrix:

$$\mathcal{M}_{\tilde{e}_R,ij}^2 = m_{\tilde{l}}^2 (\delta_{ij} + \Delta_{\tilde{e}_R,ij}), \tag{9}$$

$$\mathcal{M}_{\tilde{e}_L,ij}^2 = m_{\tilde{l}}^2 (\delta_{ij} + \Delta_{\tilde{l}_L,ij}), \tag{10}$$

$$\mathcal{M}_{\tilde{\nu}_L,ij}^2 = m_{\tilde{l}}^2 (\delta_{ij} + \Delta_{\tilde{l}_L,ij}), \tag{11}$$

and calculate the LFV rate; the detailed formulae of the LFV rates can be found in [22].

We first consider the $\mu \to e\gamma$ process. The simplest possibility to induce the $\mu \to e\gamma$ process is to introduce non-vanishing values of $\Delta_{\tilde{l}_L,12}$ and/or $\Delta_{\tilde{e}_R,12}$. In Fig. 1, we plot $Br(\mu \to e\gamma)$ in such a case as a function of the slepton mass $m_{\tilde{l}}$. Here, we take $\Delta_{\tilde{l}_L,12} = \Delta_{\tilde{l}_L,21} = \Delta_{\tilde{e}_R,12} = \Delta_{\tilde{e}_R,21} = 0.1$. (Other components of $\Delta_{\tilde{f},ij}$ are taken to be zero; notice that, in such a case, $Br(\mu \to e\gamma)$ is approximately proportional to $\Delta_{l,12}^2$.) In order to see how large $Br(\mu \to e\gamma)$ can be, we adopt relatively large value of $\tan\beta$ (which is the ratio of up- and down-type Higgs bosons); in our numerical calculation, we take $\tan\beta = 50$. (For the case of large $\tan\beta$, $Br(\mu \to e\gamma)$ is approximately proportional to $\tan^2\beta$.) For the gaugino mass, we consider two cases: the heavy gaugino case with GUT

[1] If the μ parameter is as large as the gravitino mass, the gaugino masses may not obey the simple anomaly-mediation relation [19,12]. Here, we assume that such an effect is negligible.

[2] The AMSB scenario without singlet field is advantageous for cosmology because the lightest neutralino (which may be the neutral Wino) can be a good candidate of dark matter [19,20], and also because the cosmological problem due to the late-time decay of the SUSY breaking field (i.e., the so-called Polonyi problem) may be avoided [21].

[3] One might think $\Delta_{\tilde{f},ij}$ is naturally of $\sim O(1)$. However, if the off-diagonal elements are larger than the diagonal ones, there may exist a negative eigenvalue of $\mathcal{M}_{\tilde{f},ij}^2$, which results in color and/or charge breaking vacuum. Thus, in order to maintain the $SU(3)_C \times U(1)_{\rm em}$ symmetry, we assume $\Delta_{\tilde{f},ij} \sim O(0.1)$.

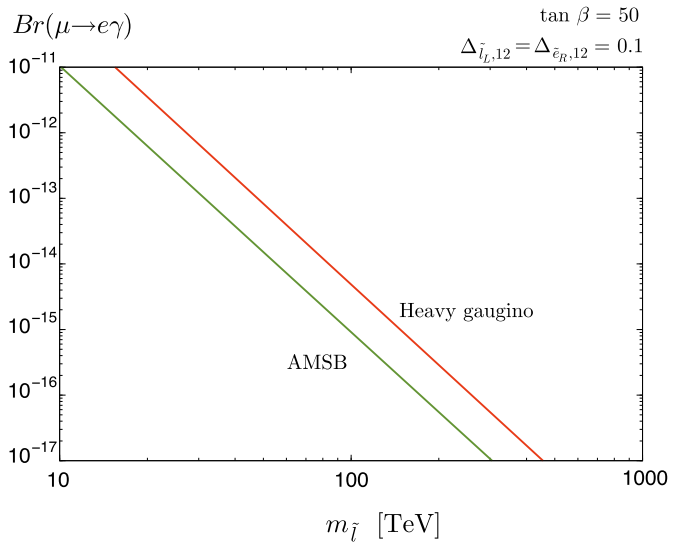

Fig. 1. $Br(\mu \to e\gamma)$ as a function of the slepton mass $m_{\tilde{l}}$ for $\tan\beta = 50$ and $\mu = m_{\tilde{l}}$. In addition, $\Delta_{\tilde{l}_L,12} = \Delta_{\tilde{l}_L,21} = \Delta_{\tilde{e}_R,12} = \Delta_{\tilde{e}_R,21} = 0.1$, while other components of $\Delta_{\tilde{f},ij}$ are taken to be zero. Upper (red) and lower (green) lines are for the heavy gaugino case with $M_3 = m_{\tilde{l}}$ and the AMSB case with $m_{3/2} = 5m_{\tilde{l}}$, respectively. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this Letter.)

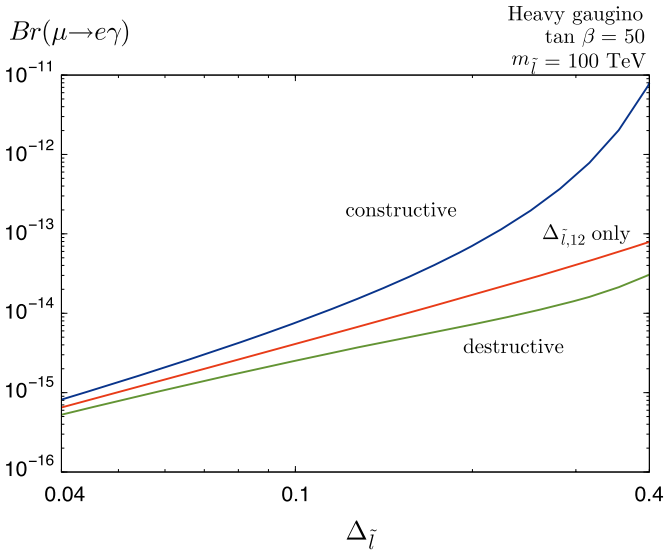

Fig. 2. $Br(\mu \to e\gamma)$ as a function of the slepton mixing parameter $\Delta_{\tilde{l}}$ (which is the absolute value of the off-diagonal elements of $\Delta_{\tilde{l}_L,ij}$ and $\Delta_{\tilde{l}_R,ij}$) for the heavy gaugino scenario. We take $\tan\beta = 50$ and $\mu = m_{\tilde{l}} = 100$ TeV. We show the constructive case ($\Delta_{\tilde{e}_R,12} = \Delta_{\tilde{e}_R,13} = \Delta_{\tilde{e}_R,23} = \Delta_{\tilde{l}_L,12} = \Delta_{\tilde{l}_L,13} = \Delta_{\tilde{l}_L,23}$) and destructive case ($-\Delta_{\tilde{e}_R,12} = \Delta_{\tilde{e}_R,13} = \Delta_{\tilde{e}_R,23} = -\Delta_{\tilde{l}_L,12} = \Delta_{\tilde{l}_L,13} = \Delta_{\tilde{l}_L,23}$). For comparison, we also show the results only with $\Delta_{\tilde{e}_R,12}$ and $\Delta_{\tilde{l}_L,12}$.

relation (with $M_3 = m_{\tilde{l}}$) and the AMSB case (with $m_{3/2} = 5m_{\tilde{l}}$).[4] In addition, we take $\mu = m_{\tilde{l}}$. We can see that $Br(\mu \to e\gamma)$ becomes smaller in the AMSB case. This is because, in the case of large $\tan\beta$, the $\mu \to e\gamma$ process is induced by diagrams with chirality-flip due to the gaugino mass. Thus, the amplitude for the AMSB case is suppressed by the factor of about $M_{1,2}/m_{\tilde{l}}$.

So far, we have considered the case where $\mu \to e\gamma$ is dominantly induced by the 12 elements of $\Delta_{\tilde{l}_L,ij}$ and $\Delta_{\tilde{l}_R,ij}$. Other components may, however, also affect $Br(\mu \to e\gamma)$. In particular, $Br(\mu \to e\gamma)$ can be enhanced if the product $\Delta_{\tilde{e}_R,13}\Delta_{\tilde{l}_L,32}$ or $\Delta_{\tilde{l}_L,13}\Delta_{\tilde{e}_R,32}$ are non-vanishing [23]. This is because, in such a case, left–right mixing occurs due to the Yukawa interaction of tau-lepton instead of that of muon. To see how large the branching ratio can be, we also calculated $Br(\mu \to e\gamma)$ for the case where the absolute values of all the off-diagonal elements of $\Delta_{\tilde{l}_L,ij}$ and $\Delta_{\tilde{l}_R,ij}$ are equal. (Here, we take $\Delta_{\tilde{l}_L,ii} = \Delta_{\tilde{l}_R,ii} = 0$.) In such a case, the magnitude of the amplitude proportional to the tau Yukawa coupling constant and that to the muon Yukawa coupling constant become comparable when the off-diagonal elements are about 0.1. The relative phase between those two amplitudes depends on the phases in the off-diagonal elements. In Fig. 2, for the fixed value of the slepton mass of $m_{\tilde{l}} = 100$ TeV, we plot $Br(\mu \to e\gamma)$ as a function of $\Delta_{\tilde{l}}$ (which is the absolute value of the off-diagonal elements of $\Delta_{\tilde{e}_R,ij}$ and $\Delta_{\tilde{l}_L,ij}$) for the cases where the two amplitudes are constructive and destructive.

Our results should be compared with the experimental bounds on the leptonic flavor and CP violations as well as with the prospects of future experiments (see Table 1). As one can see, the heavy sfermion scenario is already constrained by the present bounds on $Br(\mu \to e\gamma)$ if slepton masses are below $O(10)$ TeV. In the future, experimental bound on $Br(\mu \to e\gamma)$ may be significantly improved by the MEG upgrade, with which $\mu \to e\gamma$ may be found if the branching ratio is larger than $\sim 6 \times 10^{-14}$. With the choice of parameters used in Fig. 1, the MEG upgrade will cover the slepton mass up to ~ 55 TeV, 43 TeV, 25 TeV (36 TeV, 28 TeV, 16 TeV) for $\tan\beta = 50$, 30, and 10, respectively, for the heavy gaugino case (the AMSB case).

Table 1
Current bounds on the leptonic flavor and CP violations (which are with "♡"), as well as the future prospects. (D is the number of the days of operation.)

$\mu \to e\gamma$	MEG (current) [26]♡	$Br(\mu \to e\gamma) < 5.7 \times 10^{-13}$
	MEG Upgrade [27]	$Br(\mu \to e\gamma) \lesssim 6 \times 10^{-14}$
$\mu \to 3e$	SINDRUM I [28]♡	$Br(\mu \to 3e) < 1 \times 10^{-12}$
	Mu3e Phase I [29]	$Br(\mu \to 3e) \lesssim 10^{-15}$
	Mu3e Phase II [29]	$Br(\mu \to 3e) \lesssim 10^{-16}$
μ-e conversion	SINDRUM II [30]♡	$R_{\mu e} < 7 \times 10^{-13}$ (with Au)
	DeeMe [31]	$R_{\mu e} \lesssim 10^{-14}$ (with SiC)
	Mu2e [32]	$R_{\mu e} \lesssim 2.4 \times 10^{-17}$ (with Al)
	COMET [33]	$R_{\mu e} \lesssim 10^{-17}$ (with Al)
	PRISM/PRIME [33]	$R_{\mu e} \lesssim 2 \times 10^{-19}$
Electron EDM	YbF molecule [34]♡	$\lvert d_e\rvert < 10.5 \times 10^{-28} e$ cm
	ThO molecule [35]	$\lvert d_e\rvert \lesssim 3.7 \times 10^{-29}/\sqrt{De}$ cm
	Fr [36]	$\lvert d_e\rvert \lesssim 1 \times 10^{-29} e$ cm
	YbF molecule [37]	$\lvert d_e\rvert \lesssim 1 \times 10^{-30} e$ cm
	WN ion [38]	$\lvert d_e\rvert \lesssim 1 \times 10^{-30} e$ cm

We also study $\mu \to 3e$ and μ-e conversion processes. Here, we are paying particular attention to the case where $\tan\beta$ is large, with which the LFV rates are enhanced. Then, dipole-type operators give dominant contributions to the LFV processes because their coefficients are proportional to $\tan\beta$. In such a case, we can use the following approximated formula to evaluate $Br(\mu \to 3e)$ [24]:

$$\frac{Br(\mu \to 3e)}{Br(\mu \to e\gamma)} \simeq \frac{\alpha}{3\pi}\left(\log\frac{m_\mu^2}{m_e^2} - \frac{11}{4}\right) \simeq 6.6 \times 10^{-3}, \tag{12}$$

where α is the fine structure constant. In addition, in the case of the dipole dominance, the μ-e conversion rate, which is defined as

[4] For the AMSB case, if we naively take $m_{3/2} = m_{\tilde{l}}$, the gluino mass may conflict with the LHC bounds in some parameter region of our study below. Thus, we assume a slight suppression of the slepton mass relative to the gravitino mass in the AMSB case.

$$R_{\mu e} \equiv \frac{\Gamma(\mu N \to eN)}{\Gamma(\mu N \to \text{capture})}, \tag{13}$$

is also approximately proportional to $Br(\mu \to e\gamma)$ as

$$\frac{R_{\mu e}}{Br(\mu \to e\gamma)} \simeq \frac{\pi D_N^2}{m_\mu^5 \tau_\mu \Gamma(\mu N \to \text{capture})}, \tag{14}$$

where τ_μ is the lifetime of muon, and D_N is the overlap integral for the conversion process with nucleus N. Using the values of D_N (with method 2) and capture rates given in [25], the ratio $R_{\mu e}/Br(\mu \to e\gamma)$ is given by 2.5×10^{-3} and 3.0×10^{-3}, for N being $^{27}_{13}$Al and $^{197}_{79}$Au, respectively. We can see that $\mu \to 3e$ and μ-e conversion processes may be also within the reaches of future experiments. (See Table 1.)

Before discussing the electron EDM, we comment on LFV decay processes of τ-lepton. If $\Delta_{\tilde{l},23}$ and $\Delta_{\tilde{l},12}$ are of the same order, $Br(\tau \to \mu\gamma)$ becomes comparable to $Br(\mu \to e\gamma)$. Given the fact that even the BELLE II experiment will reach $Br(\tau \to \mu\gamma) \sim 2.4 \times 10^{-9}$ [39], it is harder to find the LFV processes of τ-lepton in models with heavy sfermions.

Next, let us consider the SUSY contribution to the electron EDM $d_e^{(\text{SUSY})}$. There are many sources of CP violation in the present model; the phase in the μ parameter, those in the off-diagonal elements of the sfermion mass matrices, and so on. With those, the SUSY contribution to the electron EDM may become sizable as we see below.

The SUSY contribution to the electron EDM can be generated even without flavor violation. However, with non-vanishing values of $\Delta_{\tilde{e}_R,ij}$ and $\Delta_{\tilde{l}_L,ij}$, $d_e^{(\text{SUSY})}$ may be enhanced [23,40]. This is because the left–right mixing of the smuon and/or stau, which are larger than that of selectron, may contribute to $d_e^{(\text{SUSY})}$. In the following, we show results for the case with and without sizable flavor violations. We calculate one-loop SUSY diagrams contributing to the electron EDM, with slepton and gaugino inside the loop.[5]

In Fig. 3, we plot $d_e^{(\text{SUSY})}$ as a function of the slepton mass $m_{\tilde{l}}$. We consider the cases of heavy gauginos and AMSB. We adopt two cases of off-diagonal elements of the sfermion mass matrices. The first one is the case without flavor violation; we take $\Delta_{\tilde{e}_R,ij} = \Delta_{\tilde{l}_L,ij} = 0$. In this case, using the fact that the phase of μ can be arbitrary, we choose $\text{Arg}(\mu)$ which maximizes $d_e^{(\text{SUSY})}$.[6] In addition, to see the effects of muon and tau Yukawa coupling constants, we consider the second case in which 13 and 31 components of $|\Delta_{\tilde{e}_R,ij}|$ and $|\Delta_{\tilde{l}_L,ij}|$ are 0.1. (Other components of $\Delta_{\tilde{e}_R,ij}$ and $\Delta_{\tilde{l}_L,ij}$ are taken to be zero.) In this case, we assumed that the phase in μ is negligible, and the phases of the off-diagonal elements are chosen such that $d_e^{(\text{SUSY})}$ is maximized. Notice that, in the second case, $d_e^{(\text{SUSY})}$ is proportional to $|\Delta_{\tilde{e}_R,13}\Delta_{\tilde{l}_L,31}|$ (with the phases of off-diagonal elements being fixed). As in the case of the LFV, the electron EDM in the heavy gaugino case is larger than that in the AMSB case. In addition, d_e is more enhanced as $\tan\beta$ becomes larger; d_e is approximately proportional to $\tan\beta$.

So far, we have concentrated on the leptonic flavor and CP violations. However, it is well known that the SUSY models may also affect flavor and CP violations of baryons. It is often the case that K^0-$\bar{K}^0$ mixing parameters, in particular the ϵ_K parameter, give very stringent constraints on the scale of the superparticle masses

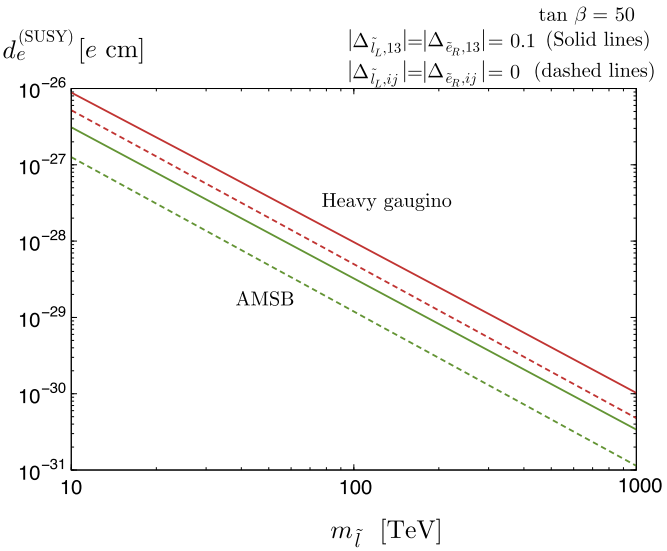

Fig. 3. The SUSY contribution to the electron EDM as a function of the slepton mass $m_{\tilde{l}}$ for $\tan\beta = 50$ and $|\mu| = m_{\tilde{l}}$. Upper (red) and lower (green) lines are for the heavy gaugino case with $M_3 = m_{\tilde{l}}$ and the AMSB case with $m_{3/2} = 5m_{\tilde{l}}$, respectively. For the dashed lines, all the elements in $\Delta_{\tilde{e}_R,ij}$ and $\Delta_{\tilde{e}_l,ij}$ are taken to be zero, and the phase of μ is chosen to maximize $d_e^{(\text{SUSY})}$. For the solid lines, $\Delta_{\tilde{e}_R,ij} = \Delta_{\tilde{l}_L,ij} = 0$ except for $|\Delta_{\tilde{e}_R,13}| = |\Delta_{\tilde{e}_R,31}| = |\Delta_{\tilde{l}_L,13}| = |\Delta_{\tilde{l}_L,31}| = 0.1$ and $\text{Arg}(\mu) = 0$; in this case, the phases of the off-diagonal elements are chosen to maximize $d_e^{(\text{SUSY})}$. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this Letter.)

[15,42–44]. To see the importance of the constraints from the SUSY contribution to K^0-$\bar{K}^0$ mixing parameters, we parameterize the mass matrices as[7]

$$\mathcal{M}^2_{\tilde{d}_R,ij} = m^2_{\tilde{q}}(\delta_{ij} + \Delta_{\tilde{d}_R,ij}), \tag{15}$$

$$\mathcal{M}^2_{\tilde{d}_L,ij} = m^2_{\tilde{q}}(\delta_{ij} + \Delta_{\tilde{d}_L,ij}), \tag{16}$$

$$\mathcal{M}^2_{\tilde{u}_R,ij} = m^2_{\tilde{q}}\delta_{ij}, \tag{17}$$

and calculate the SUSY contribution.

In Fig. 4, we plot the SUSY contribution $\epsilon_K^{(\text{SUSY})}$ for the heavy gaugino case. (We checked that the result for the AMSB case does not change so much, and hence the following arguments are also applicable to the AMSB case.) In the calculation of $\epsilon_K^{(\text{SUSY})}$, we use the formulae given in [45]. The dashed line shows the result for the case where there is no cancellation; here, we take $|\Delta_{\tilde{d}_R,ij}| = |\Delta_{\tilde{d}_L,ij}| = 0.1$ $(i \neq j)$ and $\text{Arg}(\Delta_{\tilde{d}_R,12}\Delta_{\tilde{d}_L,12}) = \pi/2$.

Comparing the result with the bound on the possible extra contribution to ϵ_K, which we conservatively take $|\epsilon_K^{(\text{SUSY})}| = 9.8 \times 10^{-4}$,[8] the squarks are required to be heavier than 450 TeV. (Notice that $\epsilon_K^{(\text{SUSY})}$ is approximately proportional to $|\Delta_{\tilde{d}_R,12}\Delta_{\tilde{d}_L,12}|$ unless there is cancellation.) Even if the squark and slepton masses are of the same order, the future electron EDM experiment may have a sensitivity to such a parameter region if the bound on d_e is improved by three orders of magnitude (see Fig. 3).

If the constraint from ϵ_K is relaxed, the future LFV experiments also play important role to probe the heavy sfermion scenario. One observation is that $\epsilon_K^{(\text{SUSY})}$ is suppressed if the following relation

[5] If the μ parameter is as small as gaugino masses, two-loop diagram may dominate the SUSY contribution to d_e [41]. In the present case, the μ parameter is as large as sfermion masses.

[6] We use the convention such that the gaugino masses as well as the vacuum expectation values of the Higgs bosons are real.

[7] The right-handed up-type squarks are not important for the discussion of K^0-$\bar{K}^0$ mixing, so we simply take $\Delta_{\tilde{u}_R,ij} = 0$.

[8] We take into account the large uncertainty of the standard model contribution $\epsilon_K^{(\text{SM})} = (1.81 \pm 0.28) \times 10^{-3}$ [46], while we use the experimental value $\epsilon_K^{(\text{exp})} = (2.228 \pm 0.011) \times 10^{-3}$ [47].

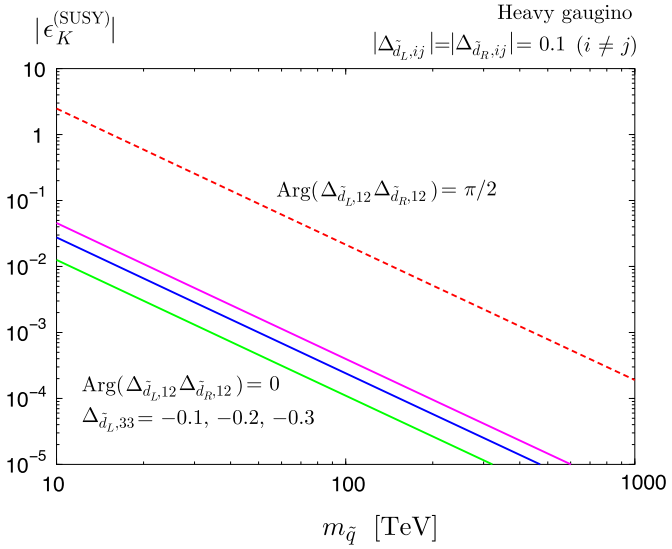

Fig. 4. The SUSY contribution to the ϵ_K parameter as a function of $m_{\tilde{q}}$. Here, we consider heavy gaugino scenario taking $M_3 = m_{\tilde{q}}$. For the dashed line, we take $|\Delta_{\tilde{d}_R,ij}| = |\Delta_{\tilde{d}_L,ij}| = 0.1$ $(i \neq j)$ and $\mathrm{Arg}(\Delta_{\tilde{d}_R,12}\Delta_{\tilde{d}_L,12}) = \pi/2$. For solid lines, we take $|\Delta_{\tilde{d}_R,ij}| = 0.1$ $(i \neq j)$ with the $SO(10)$ relation given in Eq. (18), except for $\Delta_{\tilde{d}_R,33} = 0$ and $\Delta_{\tilde{d}_L,33} = -0.1,\ -0.2,$ and -0.3 from below; the phases of the off-diagonal elements are chosen to maximize $\epsilon_K^{(\mathrm{SUSY})}$.

(which we call the $SO(10)$ relation because of the reason explained below) holds:

$$SO(10):\ \mathcal{M}^2_{\tilde{d}_R,ij} = \mathcal{M}^{2*}_{\tilde{d}_L,ij}. \tag{18}$$

This is from the fact that, if we limit ourselves to the sector consisting of down-type (s)quarks and gluino, which gives the dominant contribution to $\epsilon_K^{(\mathrm{SUSY})}$, the Lagrangian becomes invariant under the exchange of d_L and d_R^c (i.e., C invariance). Thus, if the relation (18) is exact, the SUSY contribution to ϵ_K comes from $SU(2)_L$ and $U(1)_Y$ interactions, resulting in a significant suppression of $\epsilon_K^{(\mathrm{SUSY})}$.

It is notable that the relation (18) naturally arises from $SO(10)$ unification of the gauge groups (at least at the GUT scale). In $SO(10)$ model, $q_{L,I}$, $u^c_{R,I}$, $d^c_{R,I}$, $l_{L,I}$ and $e^c_{R,I}$ (as well as right-handed neutrino) are embedded into a single **16**-representation multiplet of $SO(10)$, which we denote $\mathbf{16}_I$. In such a model, the MSSM Yukawa interactions responsible for the quark and lepton masses arise from terms which are quadratic in **16**-multiplets, and the Yukawa matrices in the MSSM originate to symmetric 3×3 matrices.[9] The down-type Yukawa matrix in this basis can be diagonalized by a single unitary matrix, which we denote U_d. Then, the down-type squarks in the flavor-eigenstate basis are related to $\tilde{d}_{L,I}$ and $\tilde{d}^c_{R,I}$ (which are in the same $\mathbf{16}_I$-multiplet) as $\tilde{d}_{L,i} = U_{d,iJ}\tilde{d}_{L,J}$ and $\tilde{d}^{c*}_{R,i} = U^*_{d,iJ}\tilde{d}^{c*}_{R,J}$, from which we obtain the $SO(10)$ relation (18).

Even if the relation (18) is satisfied at the GUT scale, however, it may not hold at the lower energy scale. In particular, the RG effects change the relation. The most important RG effect is from the Yukawa coupling constants of third generation quarks (i.e., top and bottom quarks), with which the 33 components of the mass matrices of $\tilde{q}_L$ and $\tilde{d}^c_R$ are reduced. Detailed values of $\Delta_{\tilde{d}_L,33}$ and $\Delta_{\tilde{d}_R,33}$ depend on various parameters. To see how large the difference may become, we simply adopted the assumption of the universal scalar masses at the GUT scale and estimated $\Delta_{\tilde{d}_L,33}$ and $\Delta_{\tilde{d}_R,33}$ at the scale of $m_{\tilde{q}}$. Then, we found that $(\Delta_{\tilde{d}_L,33}, \Delta_{\tilde{d}_R,33}) \simeq (0.7, 1)$, $(0.7, 0.8)$, and $(0.5, 0.5)$, for $\tan\beta = 10$, 30, and 50, respectively.[10] We also calculate the SUSY contribution to ϵ_K for the case where the relation (18) holds except for 33 components; we show the results in Fig. 4 for the cases with $\Delta_{\tilde{d}_R,33} \neq \Delta_{\tilde{d}_L,33}$. Off-diagonal elements are taken as $\Delta_{\tilde{d}} = |\Delta_{\tilde{d}_R,ij}| = 0.1$. (In this case, $\epsilon_K^{(\mathrm{SUSY})}$ is approximately proportional to $\Delta_{\tilde{d}}^3$ with $\Delta_{\tilde{d}_R,33} - \Delta_{\tilde{d}_L,33}$ being fixed.) As one can see, the bound on $m_{\tilde{q}}$ is significantly relaxed in this case in particular when $\Delta_{\tilde{d}_R,33}$ and $\Delta_{\tilde{d}_L,33}$ are close. In fact, in the case where $\epsilon_K^{(\mathrm{SUSY})}$ is suppressed, one should also consider the constraint from the K_L-K_S mass difference Δm_K. We have also calculated the SUSY contribution $\Delta m_K^{(\mathrm{SUSY})}$, and found that $m_{\tilde{q}}$ smaller than about 25 TeV is excluded with the present choice of parameters by requiring that $\Delta m_K^{(\mathrm{SUSY})}$ should be smaller than the present experimental value.

In summary, we have studied the leptonic flavor and CP violations in supersymmetric models with heavy sfermions. We have shown that the SUSY contribution to the leptonic flavor and CP violations can be so large that the future experiments may observe the signal even if $m_{\tilde{l}} \sim O(10\text{–}100$ TeV). The ϵ_K parameter often gives a very stringent constraint on the mass scale of superparticles if the off-diagonal elements of the sfermion mass matrices are sizable. However, it should be noted that the SUSY contributions to the leptonic flavor and CP violations and those to K^0-$\bar{K}^0$ mixing parameters depend on different parameters. Thus, it is important to look for signals of new physics using leptonic flavor and CP violation experiments. In particular, we have shown that, in some class of model like the $SO(10)$ unification model, suppression of $\epsilon_K^{(\mathrm{SUSY})}$ may occur because of the automatic cancellation due to the approximate C invariance. In addition, $\epsilon_K^{(\mathrm{SUSY})}$ may be suppressed due to an accidental cancellation. In this Letter, we have concentrated on the leptonic sector (as well as the constraints from K^0-$\bar{K}^0$ mixing). Other possible signals of the heavy sfermion scenario may be hidden in the B physics. Such a possibility, as well as more detailed studies of the leptonic flavor and CP violations, will be given elsewhere [48].

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[9] Here, we assume that the effects of higher-dimensional operators are subdominant.

[10] One of the reasons why $\Delta_{\tilde{d}_L,33}$ and $\Delta_{\tilde{d}_R,33}$ are close to each other in particular for the large $\tan\beta$ case is that the SUSY breaking masses for up- and down-type Higgses are taken to be equal at the GUT scale.

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