CP violation in models with TeV-scale SUSY breaking

Motoi Endo * and Yutaka Sakamura †

Department of Physics, Tohoku University
Sendai 980-8578, Japan

Abstract

We estimate contributions of the goldstino supermultiplet to the electric dipole moment (EDM) in the case that the SUSY breaking scale $\Lambda_S$ is of order the TeV scale. We found that such contributions can saturate the experimental bound if $\Lambda_S$ is close to the soft mass scale. We also discuss EDM in the gaugino mediated scenario on the warped geometry as an example of models with the TeV-scale SUSY breaking.

* e-mail address: endo@tuhep.phys.tohoku.ac.jp
† e-mail address: sakamura@tuhep.phys.tohoku.ac.jp
1 Introduction

Supersymmetry (SUSY) is one of the most promising candidate for the physics beyond the standard model. It must be, however, broken at low-energy since no superparticles have been observed yet. Although the SUSY breaking scale $\Lambda_S$ is considered to be much higher than the weak scale $G_F^{-1/2}$ in the conventional scenarios, an experimental lower bound on it is relatively weak, i.e., $\Lambda_S \gtrsim G_F^{-1/2}$, which come from the collider experiments\(^1\) [1]. In this paper, we will consider a possibility that $\Lambda_S$ is close to its lower bound, typically the TeV scale. In such a case, we should take into account not only the fields of the minimal supersymmetric standard model (MSSM), but also the *goldstino* supermultiplet as the physical degrees of freedom in the low-energy effective theory. Couplings between the goldstino supermultiplet and the MSSM fields are generally suppressed by negative powers of $\Lambda_S$ [2], and thus they become relevant to the low-energy phenomena [3, 4, 5] when $\Lambda_S$ is close to the soft SUSY-breaking mass scale $\tilde{m}$.

Among low-energy phenomena, CP violation is very sensitive to the physics beyond the weak scale. Due to the appearance of the goldstino supermultiplet, there are additional sources of CP violation besides the MSSM ones. Such new CP violating sources generally induce sizable contributions to the electric dipole moments (EDMs) and thus in this paper, we will estimate the contributions of the goldstino supermultiplet to the EDMs.

The paper is organized as follows. In the next section, we will provide notations and an effective theory used in our discussion. In Sect. 3, we will calculate the contributions of the goldstino supermultiplet to the electron EDM. In Sect. 4, we will estimate EDM in the gaugino mediation model with the warped geometry which is an example of models with the TeV-scale SUSY breaking. Sect. 5 is devoted to the summary.

2 Effective theory

We will basically follow the effective theory approach in Ref. [4]. In general, the goldstino is absorbed into the gravitino and becomes a longitudinal component of the massive gravitino. In the case considered here, the gravitino mass $m_{3/2}$ is several orders of magnitude below the eV scale. Therefore, a typical energy scale is much higher than $m_{3/2}$ in most of the physical processes, and the gravitino interactions can be well approximated by those of the goldstino in the global SUSY limit, thanks to the supersymmetric equivalence theorem [8]. So we will work in the global SUSY limit in the following.

The relevant fields to our discussion consist of the following supermultiplets\(^2\). The left-handed lepton doublet $L = (\tilde{\ell}, \ell, F^l)$, the left-handed positron $E^c = (\tilde{e}^c, e^c, F^{e^c})$, the goldstino supermultiplet $Z = (z, \psi_z, F^z)$, the Higgs doublets $H_i = (H_i, \tilde{h_i}, F^{H_i})$ ($i = 1, 2$), and the gauge supermultiplets $V_1 = (\lambda_1, A_{1\mu}, D_1)$ and $V_2^a = (\lambda_2^a, A_{2\mu}^a, D_2^a)$ ($a = 1, 2, 3$) for $U(1)_Y$ and $SU(2)_W$, respectively.

\(^1\)Constraints from cosmology and astrophysics are somewhat weaker [2, 3].
\(^2\)We use the notations of Ref. [7] in this paper.
The effective theory is described, up to higher-derivative terms, by a Kähler potential $K$, a superpotential $w$ and gauge kinetic functions $f_i$ ($i = 1, 2$) as

$$L = \int d^4\theta K + \left[ \int d^2\theta w + \int d^2\theta \frac{1}{4} \left\{ f_1 W_1^2 + f_2 \text{tr} \left( W_2^2 \right) \right\} + \text{h.c.} \right],$$

where $W_{1\alpha}$ and $W_{2\alpha}$ are superfield strength of $V_1$ and $V_2 \equiv V_2^a \cdot (\sigma^a/2)$, respectively. Their general forms with the above field content are given by

$$K = |Z|^2 + \bar{L}e^{2(\frac{-i}{2}g_1 V_1 + g_2 V_2 \cdot \sigma)} L + \bar{E}^c e^{2g_1 V_1} E^c - \frac{\alpha_e}{4\Lambda^2} |Z|^4$$

$$- \frac{\alpha_l}{\Lambda^2} |Z|^2 |L|^2 - \frac{\alpha_e}{\Lambda^2} |Z|^2 |E^c|^2 - \left\{ \frac{\gamma_K}{2\Lambda^3} \bar{Z}^2 \bar{E}^c H_1 + \text{h.c.} \right\} + \cdots,$$

$$w = -F^* Z - \frac{\sigma}{6} |Z|^3 + y_e L^c H_1 - \frac{\rho}{\Lambda} Z L^c H_1 + \mu H_1 H_2 + \cdots,$$

$$f_i = 1 + \frac{2\eta_i}{\Lambda} Z + \frac{\gamma_{f_i}}{\Lambda^3} L^c H_1 + \cdots, \quad (i = 1, 2)$$

where $g_1$ and $g_2$ are the gauge coupling constants of $U(1)_Y$ and $SU(2)_L$ respectively, and $y_e$ is the Yukawa coupling constant for the electron. Complex parameters $F$ and $\mu$ have a mass-dimension two and one respectively, and the other parameters are all dimensionless. $\Lambda$ denotes the cut-off scale of the effective theory and set to be real and positive. The ellipses denote terms that either do not play any role in the following discussion or can be eliminated by analytic field redefinitions.

We will assume that the only Higgs fields have non-zero VEVs.

$$\langle H_1^0 \rangle = \frac{v_1}{\sqrt{2}}, \quad \langle H_2^0 \rangle = \frac{v_2}{\sqrt{2}},$$

The other fields do not have any non-zero VEVs $^4$. Therefore, we can check that

$$\langle F^Z \rangle = F.$$

Namely, SUSY is broken spontaneously.

Using the $SU(2)_L \times U(1)_Y$ gauge symmetry, we can set the Higgs VEVs $v_1, v_2$ and the Yukawa coupling $y_e$ to be real. Furthermore, we will set the parameter $F$ to be real and positive by the field redefinition of $Z$. As a result, the complex parameters in the theory are $\gamma_K, \sigma, \rho, \mu, \eta_i$, and $\gamma_{f_i}$.

Next, let us see how the soft SUSY breaking parameters are expressed in terms of the above parameters. Considering the fact that SUSY is broken by the $F$-term of $Z$, the

$^3$The most general form of the gauge kinetic function also contains the $SU(2)_L$-triplet term $f^{(\text{trp})\alpha} = (\gamma_f^{(\text{trp})}/\Lambda^3) E^c \sigma^a H_1 + \cdots$. Contributions of such terms to EDM are similar to those of $\gamma_{f_i}$-terms in $f_i$. Then, in order to simplify the discussion, we will neglect such terms in the following.

$^4$As a result, the Kähler metric and the gauge kinetic functions are canonical at the vacuum, so that the component fields are all canonically normalized.
fermionic component $\psi_z$ corresponds to the goldstino, and thus is massless\textsuperscript{5}. Masses of the scalar partners of the goldstino, the sgoldstino, are given by

$$m^2_S = \frac{\alpha_z}{\Lambda^2} F^2 + |\sigma| F, \quad m^2_P = \frac{\alpha_z}{\Lambda^2} F^2 - |\sigma| F,$$

where

$$z \equiv \frac{e^{-i\varphi/2}}{\sqrt{2}} (S + iP). \quad (\varphi \equiv \arg(\sigma)).$$

The gaugino masses are

$$M_i \equiv \frac{\eta_i}{\Lambda} F,$$

and the $A$-parameter $A_e$ is given by

$$y_e A_e \equiv \frac{\rho}{\Lambda} F.$$

The selectron mass matrix in the basis ($\tilde{e}, \tilde{e}^\ast$) is

$$\mathcal{M}^2_e = \begin{pmatrix} \tilde{m}_{LL}^2 + m_e^2 & \tilde{m}_{RL}^2 \tilde{m}_{RR}^2 + m_e^2 \\ \tilde{m}_{RL}^2 \tilde{m}_{RR}^2 + m_e^2 & \tilde{m}_{RL}^2 \tilde{m}_{RR}^2 + m_e^2 \end{pmatrix},$$

where $m_e \equiv y_e v_1/\sqrt{2}$ is the electron mass, and

$$\tilde{m}_{LL}^2 \equiv \frac{\alpha_l}{\Lambda^2} F^2, \quad \tilde{m}_{RR}^2 \equiv \frac{\alpha_e}{\Lambda^2} F^2,$n

$$\tilde{m}_{RL}^2 \equiv m_e (A_e + \mu^* \tan \beta). \quad (\tan \beta \equiv \frac{v_2}{v_1})$$

Using these expressions, we can rewrite all the parameters except $\gamma_K$ and $\gamma_{fi}$ in Eq.(2) in terms of the soft SUSY breaking parameters and the order parameter of SUSY breaking $\sqrt{F}$.

Then, the independent CP violating phases that are relevant to the following discussion are $\arg(M_i^* A_e)$, $\arg(M_i \mu)$, $\arg(M^2 e^{-i\varphi})$, $\arg(M \gamma_K)$ and $\arg(M^* \gamma_{fi})$. The first two types of the phases are the usual MSSM ones, but the remaining ones are the new phases associated with the TeV-scale SUSY breaking.

### 3 Electron EDM

Now we will estimate EDM of the electron. For simplicity, we will suppose that the gauge kinetic functions are common, and denote the common gaugino mass and the coupling constant as $M$ and $\gamma_f$, respectively.

\textsuperscript{5}Strictly speaking, $\psi_z$ is not the goldstino itself. Since $\langle F^{H_1^0} \rangle = -\mu^* v_2/\sqrt{2}$, $\langle F^{H_2^0} \rangle = -\mu^* v_1/\sqrt{2}$, $\langle D_1 \rangle = g_1 (v_1^2 - v_2^2)/4$ and $\langle D_2 \rangle = -g_2 (v_1^2 - v_2^2)/4$ after the electroweak symmetry breaking, the genuine goldstino $G$ has its components also in the higgsinos and gauginos, and thus $\psi_z$ has a small mass term suppressed by $v_1 v_2/\Lambda^2$ [9]. However, we will refer to $\psi_z$ as the ‘goldstino’ in the following, because $\langle F^Z \rangle$ is the dominant source of SUSY breaking, i.e., $\Lambda_S \simeq \sqrt{F}$. 


Figure 1: Additional diagrams to MSSM ones which contribute to the electron EDM. $\tilde{\chi}^0$ stands for the neutralinos, and $\lambda$ is the photino.

For large $\tan \beta$, the conventional MSSM contribution to $d_e$ is given by [9]

$$
\frac{d_e^{(\text{MSSM})}}{e} \approx \frac{5\alpha \text{Im}(M^*\tilde{m}^2_{RL})}{96\pi \sin^2 \theta_W |M|^4},
$$

(11)

where $\alpha$ is the fine structure constant and $\theta_W$ is the Weinberg angle. In this case, $\tilde{m}^2_{RL} \simeq m_e \mu^* \tan \beta$. Here, we have assumed that $|M| \simeq \tilde{m}_{LL} \simeq \tilde{m}_{RR} \simeq |\mu|$, for simplicity. Since we are interested in the goldstino contributions besides the MSSM contribution, we will assume that the above $d_e^{(\text{MSSM})}$ is negligibly small, i.e., the relative phase between $M$ and $\tilde{m}^2_{RL}$ is tuned to be small, in the following.

Due to the interactions involving the goldstino supermultiplet $Z$, there are additional contributions besides the usual MSSM ones. The relevant diagrams to $d_e$ are essentially the same as those to the anomalous magnetic moment $a_\mu$, which are shown in Fig. 1 of Ref. [4]. Among them, four types of the diagrams shown in Fig. 1 can have a non-vanishing CP phases, and thus can contribute to EDM.

Contributions of (a) and (b) to EDM are suppressed by $m_e^2$ compared to the other contributions, and thus can be neglected$^6$. The remaining diagrams are logarithmically divergent.

$$
\frac{d_e^{(c)}}{e} = -\frac{m_e \text{Im}(MA_e e^{-i\phi})}{16\pi^2 F^2} \ln \frac{m_S^2}{m_P^2} - \frac{\text{Im}(\gamma K M)v_1}{16\sqrt{2}\pi^2 \Lambda^3} \left( \ln \frac{\Lambda_{UV}^2}{m_S^2} + \ln \frac{\Lambda_{UV}^2}{m_P^2} \right),
$$

(12)

$^6$Here, we neglect regularization-dependent terms which emerge from the diagrams (a) and (c) because they are cancelled with each other [4].
Figure 2: Constant contours of the value of $d_e$ on the $(\sqrt{F}, m_S/m_P)$-plane. The parameters are set as $|M| = A_e = 1$ TeV, $\arg(M_A e^{-i\varphi}) = \pi/2$, and $\gamma_K = \gamma_f = 0$. The numbers in the plot shows the values of $d_e$ in units of $\text{ecm}$.

and

$$d_e^{(d)} = -\frac{\text{Im}(\gamma_f M^*) v_1}{32\sqrt{2\pi^2}\Lambda^3} \left\{ \frac{|M|^2 \ln(\Lambda_{UV}^2/|M|^2) - \tilde{m}_{LL}^2 \ln(\Lambda_{UV}^2/\tilde{m}_{LL}^2)}{|M|^2 - \tilde{m}_{LL}^2} + (\tilde{m}_{LL}^2 \leftrightarrow \tilde{m}_{RR}^2) \right\},$$

where $\Lambda_{UV}$ is a cut-off scale. Here, we have neglected terms suppressed by $m_e$, and assumed that the both sgoldstinos $S$ and $P$ are much heavier than the electron.

First, we will consider the case that $\gamma_K$ and $\gamma_f$ are negligibly small. In this case, the result becomes finite and can be expressed by the SUSY breaking masses and the order parameter $\sqrt{F}$. Fig. 2 shows constant contours of the value of $d_e$ on the $(\sqrt{F}, m_S/m_P)$-plane. From this plot, we can see that the sgoldstino contribution can saturate the current experimental bound: $|d_e| < 4.3 \times 10^{-27}\text{ecm}$ \cite{10}, when $\sqrt{F}$ is close to the soft mass scale.

Notice that EDM vanishes when $m_S = m_P$. The reason for this is as follows. We have considered the case of $\gamma_K, \gamma_f \ll 1$, and thus the only CP phase besides the MSSM ones is $\varphi$, the phase of the parameter $\sigma$ in the superpotential. The degeneracy of the two sgoldstinos means that $\sigma = 0$. (See Eq. (5).) Therefore, when two sgoldstinos are degenerate, there exists no CP violating sources contributing to the electron EDM at one-loop level.

So far, we have assumed that $\gamma_K$ and $\gamma_f$ are suppressed by some mechanism. If they are of order one, the resultant EDM value far exceeds its experimental upper bound. Then, to make the discussion complete, we will investigate the constraints on $\gamma_K$ and $\gamma_f$ from the experimental bound on $d_e$. Here, we will assume that all soft SUSY breaking masses are equal, for simplicity. Then, the expression of $d_e$ can be simplified as

$$\frac{d_e}{e} \approx -\frac{\text{Im}(\gamma_K M) v_1}{4\sqrt{2\pi^2}\Lambda^3} \ln \frac{\Lambda}{m} - \frac{\text{Im}(\gamma_f^* M) v_1}{8\sqrt{2\pi^2}\Lambda^3} \left( \ln \frac{\Lambda}{m} - \frac{1}{2} \right),$$

(14)
where $\tilde{m} \equiv m_S = m_P = |M| = \tilde{m}_{LL} = \tilde{m}_{RR}$. Since the momentum cut-off scale $\Lambda_{UV}$ is thought to be the same order as the scale parameter $\Lambda$ that controls the higher-order terms in the effective Lagrangian, we have identified these two scales, i.e., $\Lambda_{UV} = \Lambda$.

Fig. 3 shows the value of $|d_e|$ on the $(\Lambda, |\gamma_K|)$-plane calculated by Eq. (14) with $\gamma_f = 0$. We have set the parameters as $\tilde{m} = 1$ TeV and $\tan \beta = 10$, and the CP phase is assumed to be unsuppressed, i.e., $\arg(\gamma_K M) = \pi/4$. From the current experimental bound on $|d_e|$, the dimensionless parameter $\gamma_K$ must be suppressed at least by four to five orders of the magnitude when the cut-off scale $\Lambda$ is in the range of 4-10 TeV, for example. For the smaller values of $\tan \beta$, the constraint on $\gamma_K$ becomes severer. This constraint on $\gamma_K$ can be realized if we introduce an approximate chiral symmetry that is broken explicitly by the order of the lepton masses $[4]$. In such a case, the flavor structure of $\gamma_K$ is inferred as $m_\mu (\gamma_K)_e \simeq m_\tau (\gamma_K)_\mu$. Then, we can relate $d_e$ to the anomalous magnetic moment of the muon $a_\mu$. That is,

$$a_\mu^{\text{nonSM}} \equiv a_\mu - a_\mu^{\text{SM}} \simeq \frac{2 m_\mu^2}{m_e \tan \phi_K} \cdot \frac{d_e}{e},$$

(15)

where $a_\mu^{\text{SM}}$ is the standard model contribution to $a_\mu$, $\phi_K \equiv \arg(\gamma_K M)$ is the CP phase, which is set to be $\pi/4$ now. Using this relation, the contour of $10^{-27} e\text{cm}$ in Fig. 3 corresponds to that of $2 \times 10^{-12}$ for $a_\mu^{\text{nonSM}}$. Therefore, we can see that the constraint on $\gamma_K$ from the experimental bound for $a_\mu^{\text{nonSM}}$ is weaker than that for $d_e$ by about three orders of magnitude$^7$.

Similarly, the constraint on $\gamma_f$ can be obtained by setting $\gamma_K = 0$ in Eq. (14). The

$^7$We have used $\delta a_\mu \equiv a_\mu^{\exp} - a_\mu^{\text{SM}} < \mathcal{O}(10^{-9})$ as a bound on $a_\mu^{\text{nonSM}}$ $[12]$. 

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Figure 3: Constant contours of the value of $|d_e|$ on the $(\Lambda, |\gamma_K|)$-plane. The parameters are set as $\tilde{m} = 1$ TeV, $\tan \beta = 10$ and $\arg(\gamma_K M) = \pi/4$. The numbers in the plot shows the values of $d_e$ in units of $e\text{cm}$. 

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required suppression of \( \gamma_f \) is roughly the same order as that of \( \gamma_K \). \(^8\)

Although the naturalness considerations suggest that the sgoldstino masses are favored to be the same order as the slepton mass scale \(^1\), the possibility of the light sgoldstinos is not excluded if some dynamical mechanism that protects their small masses works. For instance, in the case that the sgoldstinos are interpreted as the (Pseudo-) Nambu-Goldstone bosons for some (approximate) global symmetries, their small masses are protected against the quantum corrections. In such a case, i.e., \( m_{S}^2, m_{P}^2 \ll m_{e}^2 \), the sgoldstino contribution to EDM becomes

\[
\frac{d_{e}^{(c)}}{e} = -\frac{m_{e} \text{Im}(M^* A_{e})}{16\pi^2 F^2}.
\]

Note that this is proportional to the factor Im\((M^* A_{e})\) that is common to the conventional MSSM contribution. So if there is some mechanism that suppresses the MSSM contributions, the above sgoldstino contribution also receives the same suppression.

It is straightforward to extend the discussion to the neutron and the mercury EDMs. We can see that the situations become similar to that of the electron EDM.

### 4 EDM in the warped Gaugino mediation scenario

In this section, we will consider the gaugino mediation model with the warped geometry \(^1\) as an example of models with the TeV-scale SUSY breaking, and provide a rough estimation of the electron EDM.

The background metric is

\[
ds^2 = e^{-2k|y|} \eta_{\mu\nu} dx^\mu dx^\nu + dy^2,
\]

where \(1/k\) is the AdS curvature radius, and \(x^\mu (\mu = 0, 1, 2, 3)\) and \(y\) are the coordinates of our four dimensions and the fifth dimension, respectively.

The effective mass scale on the boundary at \(y = 0\) is the Planck mass \(M_p\), while that on the boundary at \(y = \pi R\) is \(M_p e^{-\pi k R}\), which will be associated with the TeV scale provided \(k R \simeq 12\). So we will refer to the boundaries at \(y = 0\) and \(y = \pi R\) as the Planck brane and the TeV brane, respectively. SUSY is supposed to be broken on the TeV brane, so that the order parameter of SUSY breaking \(A_S\) is \(\mathcal{O}(\text{TeV})\). We have assumed that the quark, the lepton and the Higgs supermultiplets are localized on the Planck brane and the gauge supermultiplets live in the bulk. Therefore, the SUSY breaking effects are mediated by the gauge interactions. Namely the gaugino masses are generated at tree-level, while the sfermion masses and the \(A\)-parameters are induced at one-loop level. As a result, the CP phases of the \(A\)-parameters are automatically aligned with those of the gauginos. So, the MSSM contributions to EDM are automatically suppressed in this model. On the other hand, as discussed in the previous section, the goldstino supermultiplet \(Z\) might

\(^8\)Of course, these constraints become weaker if an accidental cancellation between two terms in Eq. \(^1\) occurs.
provide a significant contributions to EDM because $\Lambda_S = O(\text{TeV})$ in this model. However, due to the separation along the fifth dimension, there are no direct couplings between $Z$ and the matter fields. So, possible one-loop diagrams are only those of type (d) in Fig. 4. Contribution from them is, however, negligible due to the large suppression by $M_P$ because the fundamental scale at $y = 0$ is $M_P$. Therefore, there are no sizable contributions to EDM at one-loop level.

The leading contributions to EDM come from the two-loop diagrams. For example, in the large $\tan \beta$ case, the diagram shown in Fig. 4 provides the dominant contribution to $d_e$.

Since the gauge supermultiplets propagate in the bulk, there are an infinite number of the Kaluza-Klein (K.K.) modes in the four-dimensional perspective. However, such K.K. modes do not provide sizable contributions because their couplings to the matter fields, which are localized on the Planck brane are highly suppressed [13]. So, only the zero-modes contribute to the estimation of $d_e$.

In this model, $\tilde{m}_{RL}^2 \simeq m_e \mu^* \tan \beta$ since the $A$-parameter $A_e$ is suppressed by the loop factor. Therefore, the electron EDM is roughly estimated as

$$\frac{|d_e|}{e} \simeq \frac{\alpha}{(4\pi)^3 \cos^2 \theta_W F^2} \left| \text{Im}(M \tilde{m}_{RL}^2 e^{-i\phi}) \right| \simeq \frac{\alpha m_e |\mu| \tan \beta}{(4\pi)^3 \cos^2 \theta_W |M| F^2} \left| \text{Im}(M^2 e^{-i\phi}) \right|. \quad (18)$$

Here, we have assumed that $|M|^2 \sim |\sigma F| = |m_S^2 - m_P^2|/2$, and that the conventional MSSM CP phase $\text{arg}(M^* \tilde{m}_{RL}^2)$ is suppressed by some mechanism. For example, in the case that $|M| \sim |\mu| \sim \sqrt{F} \sim 1 \text{ TeV}$ and the CP phase is maximal, the estimated value of $d_e$ is

$$|d_e| \sim 10^{-28} \tan \beta \text{ (ecm)}. \quad (19)$$

Hence, the predicted value of $d_e$ can be close to its experimental bound in this model if the SUSY breaking scale $\Lambda_S \simeq \sqrt{F}$ is close to the gaugino masses and $\tan \beta$ takes a value of $O(10)$. 

Figure 4: The diagram that provides the dominant contribution to $d_e$ in the large $\tan \beta$ case in the warped gaugino mediation scenario.
5 Summary and discussion

In a class of models where SUSY is spontaneously broken at the TeV scale, interactions with the goldstino supermultiplet becomes relevant to low-energy observables. In this paper, we have investigated contributions of the goldstino supermultiplet to the electron EDM $d_e$, which is one of the most sensitive observables for high-energy physics beyond the standard model. We found that the goldstino interactions induce sizable contributions to $d_e$, and they can saturate the experimental upper bound on $d_e$ even if the conventional MSSM contributions are suppressed when the SUSY breaking scale $\Lambda_S$ is close to the soft mass scale.

Concerning the relevant terms to $d_e$ at one-loop level, most of the parameters in the effective theory can be expressed in terms of the soft SUSY breaking parameters and $\Lambda_S$. Only two parameters $\gamma_K$ and $\gamma_f$ are left as free parameters. However, these two parameters must be suppressed by the factor $10^{-5} - 10^{-4}$ due to the requirement that the predicted value of $d_e$ does not exceed the current experimental bound. Such suppression can be realized, for example, by introducing a chiral symmetry that is explicitly broken by the order of the lepton masses.

One of the specific example of models with the TeV-scale SUSY breaking is the gaugino mediation scenario on the warped geometry. We have estimated the value of $d_e$ in such a model. Since there are no contributions to $d_e$ at one-loop level in this model, the dominant contributions are induced at two-loop level. Although they are suppressed by the loop factor, the value of $d_e$ can be close to the experimental bound especially in the large $\tan \beta$ case.

Another possibility of the TeV-scale SUSY breaking is the existence of the strong coupling dynamics just above the TeV scale [14, 15]. Among such a class of models, a certain type of models can be interpreted as a dual theory of the above mentioned sequestered SUSY-breaking model with the warped geometry [15], from the viewpoint of the AdS/CFT correspondence [16]. Furthermore, there is an interesting model that solves the strong CP problem in the context of the strong coupling dynamics [17]. Applying our discussion to such a model is an intriguing subject.

The search of flavor changing processes also provide valuable information on high-energy physics beyond the weak scale. In fact, the goldstino interactions contribute such processes, and their contributions can be sizable when $\Lambda_S$ is close to the soft mass scale [5]. In Ref. [5], they focused the contributions to the flavor changing processes which depend on the interactions related to the mass spectra. In this case, the sources of flavor violation of the goldstino interactions are common to those of MSSM sector. Then, if there is some mechanism that suppresses the MSSM contributions to the flavor violation, the goldstino contributions also receive the same suppression. On the other hand, there is an additional CP source to the MSSM ones in the goldstino sector even in the absence of $\gamma$-terms. Thus, the goldstino contributions can saturate the experimental bound regardless of suppression mechanisms of the MSSM ones.

The $\gamma_K$ or $\gamma_f$ term may also become the source of flavor violation. Then the goldstino contributions to the flavor violating processes become large as in the case of CP violation.
However if the chiral suppression is assumed to suppress the contribution to the EDM, the constraints from the flavor changing processes are relaxed simultaneously. Moreover though the effective Lagrangian Eq. (2) is the most general one in the context of the particle content given in this paper, there might be possible to introduce additional particles and interactions which contribute to the flavor changing processes. Those terms generally also become new sources of CP violation like the EDMs. Thus CP violation is one of the most useful phenomena to detect a signal of the TeV-scale SUSY breaking.

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