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QCD Corrections to Quark (Chromo)Electric Dipole Moments in High-scale Supersymmetry

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Abstract

Recent results from the LHC experiments, both for the Higgs mass measurement and the direct search for supersymmetric (SUSY) particles, might indicate that the SUSY breaking scale is much higher than the electroweak scale. Although it is difficult to investigate such a scenario at collider experiments, the measurement of the hadronic electric dipole moments is one of promising ways to detect the effects of the SUSY particles. These effects are expressed in terms of the CP-violating effective operators defined at the SUSY breaking scale, which involve quarks, gluons, photons, and gluinos. In this paper, we discuss the QCD corrections to the effective operators in the high-scale SUSY scenario. To appropriately evaluate the radiative corrections in the presence of large mass hierarchy among the SUSY particles, we exploit an effective theoretical approach based on the renormalization-group equations. As a result, it is found that the low-energy quark electric and chromoelectric dipole moments may differ from those evaluated in previous works by $\mathcal{O}(100)$ % and $\mathcal{O}(10)$ %, respectively.

1 Introduction

The supersymmetric (SUSY) extension of the Standard Model (SM) is a leading candidate for physics beyond the SM. So far, however, the weak-scale SUSY models have been severely restricted since no evidence for new physics has been found yet; for instance, the latest results from the LHC experiments have imposed stringent limits on the masses of the SUSY particles, especially those of colored particles [1, 2]. In addition, the Higgs boson with a mass of $\sim$ 126 GeV [3], which was recently discovered at the LHC [4, 5], might also indicate that the SUSY particles are well above $\mathcal{O}(1)$ TeV, since in the minimal supersymmetric Standard Model (MSSM) sufficient radiative corrections are required in order to realize the mass of the Higgs boson [6–10]. Unless the Higgs sector is modified nor the left- and right-handed stops adequately mix with each other, such a large quantum effect is only provided with heavy stops having masses of much higher than the electroweak scale.

The current situation motivates us to study models with a high SUSY breaking scale. Such models assume that SUSY is broken at a scale of $\mathcal{O}(10^{2-3})$ TeV to yield the 126 GeV Higgs boson [11–14]. In this case, scalar particles except the lightest Higgs boson acquire masses of the order of the SUSY breaking scale. Fermionic superpartners, on the other hand, may be much lighter than the other sparticles since their masses are protected by chiral symmetries. Indeed, such a mass spectrum is realized with a simple SUSY breaking mechanism in which SUSY is broken by a non-singlet field and the breaking effects are transmitted to the visible sector via a generic Kähler potential. In this framework, the gaugino masses are induced by the anomaly mediation [15, 16] and suppressed by one-loop factors compared with the scalar mass. With the SUSY breaking scale being $\mathcal{O}(10^{2-3})$ TeV, gauginos may lie around the TeV scale. The neutral wino turns out to be the lightest SUSY particle in this model, and may make up a main component of the dark matter in the Universe [17–19]. Further, it is found that the gauge coupling unification is not only preserved but improved in the scenario [20]. Thus, the high-scale SUSY models have interesting features from a phenomenological point of view [21–26], and recently attract a lot of attention especially after the early LHC runnings [27–34].

Although it is difficult to investigate the high-scale SUSY scenario at high-energy collider experiments, the low-energy precision experiments might catch up the SUSY signature. Without any flavor symmetries, soft SUSY breaking parameters in general give rise to extra sources of flavor and/or CP violation [35]. These effects, such as the flavor-changing neutral currents and the hadronic and leptonic electric dipole moments (EDMs), are suppressed by sfermion masses, and thus the high-scale SUSY models do not conflict with the current experimental results for these quantities. Among such experiments, the measurement of the EDMs offers a promising way to look for the signature of the SUSY particles [36, 37]. Since in the SM the EDMs induced by the CP phase in the Cabibbo-Kobayashi-Maskawa (CKM) matrix are considerably below the sensitivities of the present and near future experiments [38, 39], the EDM measurement is free from the SM background, thus provides a clean environment to detect a sign of high-energy physics beyond the SM.

The effective interactions which give rise to the EDMs are expressed in terms of the flavor-conserving CP-violating effective operators. In SUSY models, such operators are induced by diagrams in which SUSY particles run in the loop. In the case of the high-scale SUSY scenario, however, one needs to pay particular attention to the calculation of the diagrams; as mentioned above, there exists a large difference between the mass scales of scalar and fermionic SUSY particles and this hierarchy causes large logarithmic factors which may spoil the perturbation theory. To evade the difficulties, we need to evaluate the effective operators by means of the renormalization-group equations (RGEs). The renormalization corrections are particularly important for the operators including colored particles because of the large value of the strong coupling constant.

In this paper, we study the QCD effects on the flavor-preserving CP-odd quark and gluon operators generated by the squark-gluino interactions. Among the operators, the EDMs and the chromoelectric dipole moments (CEDMs) of quarks have the lowest mass-dimensions, and thus sensitive to the SUSY contribution. We focus on these two operators and study the contribution of the quark-gluino four-Fermi operators to the quantities. The calculation is divided into two steps; first, by integrating out squarks, we construct an effective theory with quarks, gluons, photons, and gluinos. Then, the effective operators are evolved down to the gluino threshold according to the RGEs. During the RGE flow, the CEDMs are radiatively generated from the dimension-six quark-gluino operators. The resultant EDMs and CEDMs evaluated in this way are compared with the results based on the computation of the one-loop diagrams. Possible ways of improvement of the calculation are also discussed.

This paper is organized as follows. In the next section, we write down the CP-violating effective operators involving quarks, gluons, photons, and gluinos which we consider in the following discussion, and present the anomalous dimension matrix for the operators. The Wilson coefficients of the effective operators are evaluated in the MSSM with and without the assumption of the minimal flavor violation in the sections 3 and 4, respectively. Evolving them down according to the RGEs, we obtain the EDMs and CEDMs of light quarks at the hadron scale, and compare them with the explicit one-loop calculation. Section 5 is devoted to conclusion and discussion.

2 Effective Lagrangian

To begin with, we write down the CP-violating effective operators at the hadron scale (~ 1 GeV) which consist of the flavor-diagonal operators of light quarks, photons, and gluons up to dimension-five:

$$\mathcal{L}_{\not{CP}} = -\sum_{q=u,d,s} m_q \bar{q} i\theta_q \gamma_5 q + \theta_G \frac{\alpha_s}{8\pi} G^A_{\mu\nu}\widetilde{G}^{A\mu\nu} - \frac{i}{2}\sum_{q=u,d,s} d_q \bar{q}\sigma^{\mu\nu}\gamma_5 q F_{\mu\nu} - \frac{i}{2}\sum_{q=u,d,s} \tilde{d}_q \bar{q} g_s \sigma^{\mu\nu}\gamma_5 T^A q G^A_{\mu\nu} \ . \tag{1}$$

Here, m_q are the quark masses, g_s is the strong coupling constant ($\alpha_s = g_s^2/4\pi$), and T^A are the generators of the $\mathrm{SU}(3)_\mathrm{C}$. $F_{\mu\nu}$ and $G^A_{\mu\nu}$ are the field strength tensors of photon and gluon, and their dual fields are defined by, *e.g.*, $\widetilde{G}^A_{\mu\nu} \equiv \frac{1}{2}\epsilon_{\mu\nu\rho\sigma}G^{A\rho\sigma}$ with $\epsilon^{0123} = +1$. The second term of the above expression is the effective QCD θ term, which is connected with the first term through the chiral rotation. These two terms are suppressed in the presence of the Peccei-Quinn symmetry [40]. The third and fourth terms represent the EDMs and the CEDMs for light quarks, respectively. They are dimension-five operators, and thus quite sensitive to the high-scale physics beyond the SM.

Apart from the SM contribution, these operators are induced by diagrams where SUSY particles run in the loop. In this paper, we focus on the SUSY contribution and discuss the QCD effects on it at the leading order in α_s. In particular, we consider the case where a large mass difference between the scalar particles and gauginos exists. As mentioned in the Introduction, such a hierarchical mass spectrum often shows up in the high-scale SUSY models. To appropriately include the QCD corrections in the presence of the large mass hierarchy, we evaluate them based on the method of the effective field theory as well as the RGEs. First, by integrating out squark fields, we obtain the effective Lagrangian below the SUSY breaking scale, which involves only gluinos and the SM fields. The short-distance effects, which reflect the CP-violation due to the SUSY particles, are included into the Wilson coefficients of the effective operators matched at the SUSY breaking scale. Next, the effective operators are evolved down to the gluino threshold according to the RGEs. Then, at the threshold, the gluino fields are integrated out to give the effective theory which contains only the SM fields. After this step, the ordinary procedure is applied to estimate the effects of the CP-violating operators on the low-energy physics such as the neutron EDM. The purpose of this paper is to formulate the first two steps in terms of the operator product expansions and the RGEs.

The effective Lagrangian below the SUSY breaking scale is given as follows:

$$\mathcal{L}_{\mathrm{eff}} = \sum_{q=u,d,s} C_1^q(\mu)\mathcal{O}_1^q(\mu) + \sum_{q=u,d,s} C_2^q(\mu)\mathcal{O}_2^q(\mu) + \sum_{q=u,d,s}\sum_{i=1}^{5} \widetilde{C}_i^q(\mu)\widetilde{\mathcal{G}}_i^q(\mu) , \tag{2}$$

where

$$\mathcal{O}_1^q \equiv -\frac{i}{2} e Q_q m_q \overline{q} \sigma^{\mu\nu} \gamma_5 q F_{\mu\nu} \ ,$$
$$\mathcal{O}_2^q \equiv -\frac{i}{2} g_s m_q \overline{q} \sigma^{\mu\nu} \gamma_5 T^A q G^A_{\mu\nu} \ ,$$
$$\widetilde{\mathcal{G}}_1^q \equiv \frac{1}{2} \bar{q} q \overline{\tilde{g}}^A i\gamma_5 \tilde{g}^A \ ,$$
$$\widetilde{\mathcal{G}}_2^q \equiv \frac{1}{2} \bar{q} i\gamma_5 q \overline{\tilde{g}}^A \tilde{g}^A \ ,$$
$$\widetilde{\mathcal{G}}_3^q \equiv \frac{1}{2} d_{ABC} \bar{q} T^A q \overline{\tilde{g}}^B i\gamma_5 \tilde{g}^C \ ,$$
$$\widetilde{\mathcal{G}}_4^q \equiv \frac{1}{2} d_{ABC} \bar{q} i\gamma_5 T^A q \overline{\tilde{g}}^B \tilde{g}^C \ ,$$
$$\widetilde{\mathcal{G}}_5^q \equiv \frac{i}{2} f_{ABC} \bar{q} \sigma^{\mu\nu} i\gamma_5 T^A q \overline{\tilde{g}}^B \sigma_{\mu\nu} \tilde{g}^C \ . \tag{3}$$

Here, Q_q are the electric charges for light quarks with $(Q_u, Q_d, Q_s) = (2/3, -1/3, -1/3)$. The covariant derivative for quarks is defined as $D_\mu \equiv \partial_\mu - ieQ_q A_\mu - ig_s G^A_\mu T^A$ $(e < 0)$ with A_μ and G^A_μ the $\mathrm{U}(1)_{\mathrm{EM}}$ and $\mathrm{SU}(3)_{\mathrm{C}}$ gauge fields, respectively. $\tilde{g}^A$ denotes gluinos, which are Majorana fermions and form an adjoint representation under the $\mathrm{SU}(3)_{\mathrm{C}}$ transformations. The totally symmetric factor d_{ABC} is defined by $d_{ABC} \equiv 2\mathrm{Tr}(\{T_A, T_B\} T_C)$, while f_{ABC} is the structure constant of the SU(3) group with $[T_A, T_B] = i f_{ABC} T_C$. The Wilson coefficients of the operators, C_1^q, C_2^q, and $\widetilde{C}_i^q$, are obtained by integrating out squark fields at the SUSY breaking scale.

In Eq. (2) we only keep the operators which give significant corrections to the quark EDMs and CEDMs. Let us comment on the operators which we ignore in the following analysis. First, we do not consider the dimension-four operators in Eq. (1) since they do not contribute to the RGEs for the dimension-five operators. Especially, when the Peccei-Quinn symmetry is imposed, these dimension-four operators are suppressed and the EDMs and CEDMs give dominant contributions to the hadronic and atomic EDMs. Also, we ignore the dimension-five gluino CEDM, $f_{ABC} g_s \overline{\tilde{g}}^A \sigma^{\mu\nu} \gamma_5 \tilde{g}^B G^C_{\mu\nu}$, since it does not affect the running of the operators in Eq. (2) at the leading order in α_s. As for the dimension-six operators, the Weinberg operator $f_{ABC} G^A_{\mu\nu} \widetilde{G}^{B\nu\lambda} G^{C\mu}_\lambda$ [41], four-quark operators, and four-gluino operators might also yield sizable effects on the radiative corrections to the operators above. These dimension-six operators are, however, generated at $\mathcal{O}(\alpha_s^2)$ in the case of the MSSM, and thus safely neglected in the leading order calculation. [1]

Next, we evaluate the anomalous dimensions of the operators in Eq. (2) at the leading order. The RGE for the Wilson coefficients in Eq. (2) is written as

$$\mu \frac{\partial}{\partial \mu} \vec{C}(\mu) = \vec{C}(\mu) \Gamma \ , \tag{4}$$

[1] Note that once these dimension-six operators are induced, they actually give rise to significant contributions to the EDMs and the CEDMs. For instance, at one-loop level, the Weinberg operator mixes with the quark CEDMs [42], while four-quark operators including heavy quarks radiatively induce both the EDMs and the CEDMs [43].

where $\vec{C}$ is a column vector defined by

$$\vec{C} \equiv (C_1^q, C_2^q, \widetilde{C}_1^q, \widetilde{C}_2^q, \widetilde{C}_3^q, \widetilde{C}_4^q, \widetilde{C}_5^q) \ . \tag{5}$$

Then, we obtain the following anomalous dimension matrix at one-loop level:

$$\Gamma = \begin{pmatrix} \frac{\alpha_s}{4\pi}\gamma_q & 0 \\ \frac{1}{(4\pi)^2}\gamma_{q\tilde{g}} & \frac{\alpha_s}{4\pi}\gamma_{\tilde{g}} \end{pmatrix} \ , \tag{6}$$

with

$$\gamma_q = \begin{pmatrix} 8C_F & 0 \\ 8C_F & 16C_F - 4N \end{pmatrix} \ , \tag{7}$$

$$\gamma_{\tilde{g}} = \begin{pmatrix} -6C_F - 6N & 0 & 0 & 0 & 2 \\ 0 & -6C_F - 6N & 0 & 0 & 2 \\ 0 & 0 & -6C_F & 0 & (N^2-4)/2N \\ 0 & 0 & 0 & -6C_F & (N^2-4)/2N \\ 24 & 24 & 12N & 12N & 2C_F - 4N \end{pmatrix} \ , \tag{8}$$

and

$$\gamma_{q\tilde{g}} = \begin{pmatrix} 0 & 0 \\ 0 & 0 \\ 0 & 0 \\ 0 & 0 \\ 0 & 8N\frac{M_{\tilde{g}}}{m_q} \end{pmatrix} \ . \tag{9}$$

Here, $N(=3)$ is the number of colors, $C_F = (N^2-1)/2N$ is the quadratic Casimir invariant for the fundamental representation, and $M_{\tilde{g}}$ is the mass of gluino. The anomalous dimension matrix for the dimension-five operators γ_q is readily obtained from that for the dipole-type operators relevant to the $b \to s\gamma$ process [44, 45]. Note that the coefficient of $\gamma_{q\tilde{g}}$ is not suppressed by the strong coupling constant α_s. In this case, the scale-dependence arises from a mismatch in the dimension between the dimension-five and -six operators [46]. A similar feature is found in the case of four-quark operators mixing into the quark EDMs and CEDMs, as discussed in Ref. [43].

3 MSSM with minimal flavor violation

Now all we have to do is to compute the Wilson coefficients in Eq. (2) by integrating out squarks in a certain model. Then, by evolving them down according to the RGE (4), we obtain the quark EDMs and CEDMs in the low-energy region. In the following discussion, we take up the MSSM as an example. Also, in this section, we focus on the case with the so-called minimal flavor violation [47, 48], which assumes that the CKM matrix is the only source for all of the flavor-violating terms in the MSSM.

In the present case, the tree-level squark exchanging diagrams give rise to the CP-odd quark-gluino four-Fermi operators, which are to induce the quark CEDMs radiatively. The squark mass matrix has the following form:

$$\mathcal{L}_{\rm mass} = -\begin{pmatrix}\tilde{q}_L^* & \tilde{q}_R^*\end{pmatrix}\begin{pmatrix} m_{\tilde{q}_L}^2 & m_q X_q \\ m_q X_q^* & m_{\tilde{q}_R}^2 \end{pmatrix}\begin{pmatrix}\tilde{q}_L \\ \tilde{q}_R\end{pmatrix} , \tag{10}$$

where $\tilde{q}_L$ and $\tilde{q}_R$ represent the left- and right-handed squarks, respectively, and $X_u \equiv A_u^* - \mu\cot\beta$ ($X_d \equiv A_d^* - \mu\tan\beta$) for up-type (down-type) quarks. Here we assume that the trilinear soft scalar couplings (the so-called A-terms) A_q are proportional to the corresponding Yukawa couplings. In X_q, μ is the higgsino-mass parameter and $\tan\beta$ is the ratio of the vacuum expectation values of the MSSM Higgs fields. Throughout this article we take a convention where the gaugino masses are set to be real parameters, without loss of generality. On the assumption of the minimal flavor violation, the flavor-mixings in the squark mass matrix are considerably suppressed, so we neglect them in the present discussion. We also take $m_{\tilde{q}_L}^2 = m_{\tilde{q}_R}^2 = M_S^2$ and $m_q X_q \ll M_S^2$, for simplicity. Then, by evaluating the squark exchanging diagrams, we readily obtain the Wilson coefficients at the scalar mass scale M_S:

$$C_1^q(M_S) = C_2^q(M_S) = 0 , \tag{11}$$

and

$$\begin{aligned}
\widetilde{C}_1^q(M_S) = \widetilde{C}_2^q(M_S) &= -\frac{1}{2N}\frac{g_s^2 m_q}{M_S^4}{\rm Im}(X_q) , \\
\widetilde{C}_3^q(M_S) = \widetilde{C}_4^q(M_S) &= -\frac{1}{2}\frac{g_s^2 m_q}{M_S^4}{\rm Im}(X_q) , \\
\widetilde{C}_5^q(M_S) &= +\frac{1}{4}\frac{g_s^2 m_q}{M_S^4}{\rm Im}(X_q) .
\end{aligned} \tag{12}$$

Notice that the quark EDMs and CEDMs vanish at tree-level. They are induced radiatively through the mixing terms in RGEs and also from the short-distance contribution, as will be shown below.

By using Eqs. (11) and (12) as initial conditions, we solve the RGE (4) to evaluate the Wilson coefficients at the gluino threshold. Especially, in the leading-logarithmic approximation, the quark CEDMs are generated as

$$C_2^q(M_{\tilde{g}}) \simeq -\frac{1}{(4\pi)^2}8N\frac{M_{\tilde{g}}}{m_q}\ln\left(\frac{M_S}{M_{\tilde{g}}}\right)\widetilde{C}_5^q(M_S) , \tag{13}$$

while the EDMs vanish at the leading order. This result is to be compared with the explicit calculation of the one-loop gluino-squark diagrams. In the limit of $M_{\tilde{g}} \ll M_S$, we have [49]

$$C_1^q|_{\rm 1loop} = +\frac{1}{(4\pi)^2}\frac{16}{3}\frac{M_{\tilde{g}}}{m_q}\widetilde{C}_5^q(M_S) , \tag{14}$$

$$C_2^q|_{\rm 1loop} = -\frac{1}{(4\pi)^2}\left[8N\ln\left(\frac{M_S}{M_{\tilde{g}}}\right) - \frac{88}{3}\right]\frac{M_{\tilde{g}}}{m_q}\widetilde{C}_5^q(M_S) . \tag{15}$$

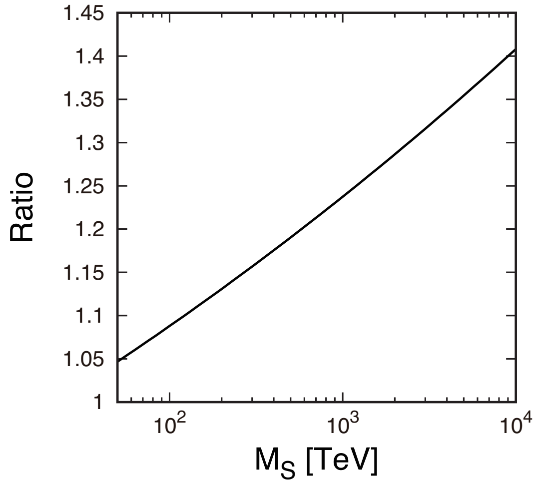

Figure 1: Ratio $C_2^q(M_{\tilde{g}})/C_2^q|_{\rm 1loop}^{(L)}$ against the squark mass M_S. Gluino mass is fixed to $M_{\tilde{g}} = 3$ TeV.

The first term in Eq. (15) is consistent with Eq. (13). The non-logarithmic terms in Eqs. (14) and (15) result from the short-distance contribution; it is induced by the processes in which the loop integrals are dominated by momenta around M_S. In that sense, the first term in Eq. (15) is to be regarded as the long-distance contribution, with the factorization scale around the squark mass scale.

To see the significance of the running effects, we evaluate C_2^q at the gluino threshold numerically and compare it with the long-distance part of Eq. (15), *i.e.*,

$$C_2^q|_{\rm 1loop}^{(L)} = -\frac{1}{(4\pi)^2} 8N \frac{M_{\tilde{g}}}{m_q} \ln\left(\frac{M_S}{M_{\tilde{g}}}\right) \widetilde{C}_5^q(M_S) \ . \tag{16}$$

The difference is caused by the running of the parameters and the mixing among the effective operators. In Fig. 1, we plot the ratio $C_2^q(M_{\tilde{g}})/C_2^q|_{\rm 1loop}^{(L)}$ against the squark mass M_S. Here, the gluino mass is fixed to $M_{\tilde{g}} = 3$ TeV. In $C_2^q|_{\rm 1loop}^{(L)}$ and $\widetilde{C}_5^q(M_S)$, we use $M_{\tilde{g}}$ and m_q evaluated at the squark mass scale. Moreover, in order to obtain $C_2^q(M_{\tilde{g}})$, the RGEs are solved using the beta function of the strong coupling constant which contains the contribution of both gluino and SM particles. Figure 1 shows that as the squark mass scale becomes large, the running-effects yield the $\mathcal{O}(10)\%$ difference between $C_2^q(M_{\tilde{g}})$ and $C_2^q|_{\rm 1loop}^{(L)}$.

Next, we take the threshold short-distance contributions into account, and evaluate both $C_1^q(M_{\tilde{g}})$ and $C_2^q(M_{\tilde{g}})$ in terms of the RGEs. Then, they are compared with the explicit one-loop results in Eqs. (14) and (15). The initial conditions for C_1^q and C_2^q are

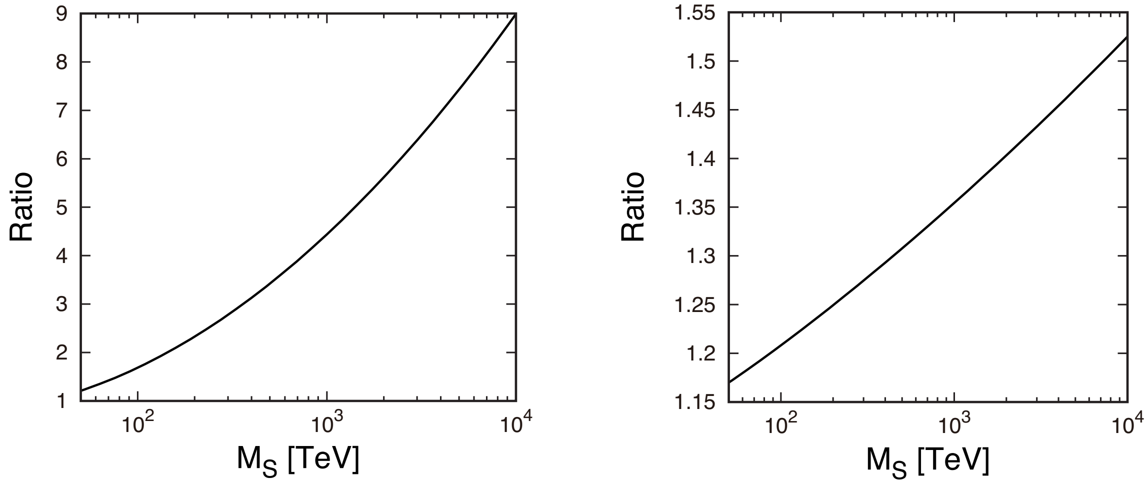

Figure 2: Ratios $C_1^q(M_{\tilde{g}})/C_1^q|_{1\mathrm{loop}}$ and $C_2^q(M_{\tilde{g}})/C_2^q|_{1\mathrm{loop}}$ as functions of M_S in left and right graphs, respectively. In both graphs, gluino mass is fixed to $M_{\tilde{g}} = 3$ TeV.

given by the short-distance contribution in Eqs. (14) and (15), that is,

$$
\begin{aligned}
C_1^q(M_S) &= +\frac{1}{(4\pi)^2}\frac{16}{3}\frac{M_{\tilde{g}}}{m_q}\,\widetilde{C}_5^q(M_S)\ , \\
C_2^q(M_S) &= +\frac{1}{(4\pi)^2}\frac{88}{3}\frac{M_{\tilde{g}}}{m_q}\,\widetilde{C}_5^q(M_S)\ , \qquad (17)
\end{aligned}
$$

while those for $\widetilde{C}_i^q$ $(i = 1, \ldots, 5)$ are given by Eq. (12). In Fig. 2, the results are plotted as functions of M_S. Here again, the gluino mass is fixed to $M_{\tilde{g}} = 3$ TeV. The left (right) panel in Fig. 2 represents the ratio $C_1^q(M_{\tilde{g}})/C_1^q|_{1\mathrm{loop}}$ $(C_2^q(M_{\tilde{g}})/C_2^q|_{1\mathrm{loop}})$. As for C_2^q, it is again found that the variation of the squark mass scale may change the ratio by $\mathcal{O}(10)\%$. In the case of C_1^q, on the other hand, the RGE result is several times larger than the explicit one-loop result, which is quite drastic compared to the case of C_2^q. It is found that this enhancement is caused by the mixing of the CEDM operators, whose contribution becomes dominant as the squark mass scale taken to be higher.

4 MSSM with a generic flavor structure

In the high-scale SUSY scenario, flavor-violation in the soft mass terms of squarks is allowed to be sizable, which motivates us to consider the case where squark mass matrices have a generic flavor structure. In such a case the dominant contributions to the EDMs and CEDMs of light quarks come from the flavor-violating processes [50, 51]. These contributions are also evaluated with the prescription described in the previous section. The Wilson coefficients of the effective operators at the SUSY breaking scale in the present case are given as

$$
\begin{aligned}
\widetilde{C}_1^q(M_S) = \widetilde{C}_2^q(M_S) &= -\frac{1}{2N}\frac{g_s^2 m_{q_3}}{M_S^4}\mathrm{Im}\big[(\delta_{LL})_{qq_3}X_{q_3}(\delta_{RR})_{q_3q}\big] \ , \\
\widetilde{C}_3^q(M_S) = \widetilde{C}_4^q(M_S) &= -\frac{1}{2}\frac{g_s^2 m_{q_3}}{M_S^4}\mathrm{Im}\big[(\delta_{LL})_{qq_3}X_{q_3}(\delta_{RR})_{q_3q}\big] \ , \\
\widetilde{C}_5^q(M_S) &= +\frac{1}{4}\frac{g_s^2 m_{q_3}}{M_S^4}\mathrm{Im}\big[(\delta_{LL})_{qq_3}X_{q_3}(\delta_{RR})_{q_3q}\big] \ ,
\end{aligned}
\tag{18}
$$

where q_3 denotes t-quark (b-quark) for the up-type (down-type) quarks, and the mass insertion parameters [35, 52] $(\delta_{LL})_{ij}$ and $(\delta_{RR})_{ij}$ are defined by

$$
(\delta_{LL})_{ij} \equiv \frac{(m^2_{\tilde{q}_L})_{ij}}{M_S^2} \ , \qquad (\delta_{RR})_{ij} \equiv \frac{(m^2_{\tilde{q}_R})_{ij}}{M_S^2} \ . \tag{19}
$$

It is possible for them to be $\mathcal{O}(1)$ in the high-scale SUSY scenario [37]. Thus, the above coefficients are enhanced by the Yukawa coupling constants of the third generation quarks without suffering from the suppression. In addition, we take into account the short-distance threshold corrections at one-loop level for C_1^q and C_2^q:

$$
\begin{aligned}
C_1^q(M_S) &= +\frac{1}{(4\pi)^2}\frac{16}{3}\frac{M_{\tilde{g}}}{m_q}\ \widetilde{C}_5^q(M_S) \ , \\
C_2^q(M_S) &= +\frac{1}{(4\pi)^2}\frac{118}{3}\frac{M_{\tilde{g}}}{m_q}\ \widetilde{C}_5^q(M_S) \ .
\end{aligned}
\tag{20}
$$

These initial conditions as well as the RGE (4) are again consistent with the one-loop results given in Ref. [51].

By using a similar procedure to that described in the previous section, we readily evaluate the EDMs and CEDMs at the gluino mass scale with initial conditions (18) and (20). Let us now evolve them down to the hadron scale. Below the gluino threshold, the gluino fields are integrated out and the effective theory includes only the SM fields. The tree-level matching condition is applied to C_1^q and C_2^q, and then they are evolved down to the hadronic scale in terms of the SM RGEs. The quark EDMs and CEDMs are then given as

$$
\begin{aligned}
d_q &= m_q(\mu_H) e Q_q C_1^q(\mu_H) \ , \\
\tilde{d}_q &= m_q(\mu_H) C_2^q(\mu_H) \ ,
\end{aligned}
\tag{21}
$$

with $\mu_H \sim 1$ GeV the hadron scale.

In Fig. 3, the absolute values of the quark EDMs $|d_q|$ and CEDMs $e|\tilde{d}_q|$ at the hadron scale $\mu_H = 1$ GeV are plotted as functions of the squark mass scale M_S. The solid and dashed lines represent the CEDMs and EDMs, respectively. The upper two red lines correspond to the EDM and CEDM of up quark, while the lower two green lines

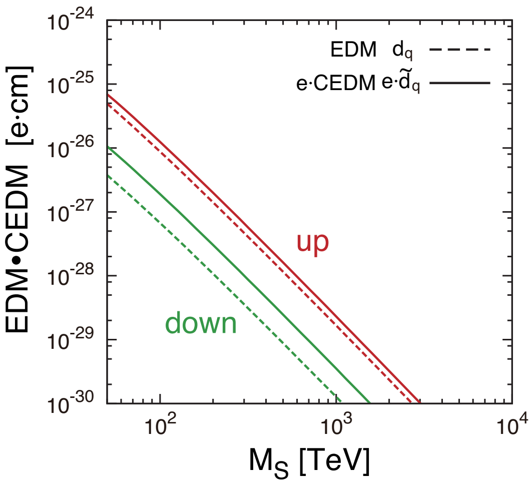

Figure 3: Quark EDMs $|d_q|$ and CEDMs $e|\tilde{d}_q|$ at the hadron scale $\mu_H = 1$ GeV as functions of M_S. Solid and dashed lines represent the CEDMs and EDMs, respectively. Upper two red lines are for up quark, while lower two green lines for down quark. We take $M_{\tilde{g}} = 3$ TeV, $\tan\beta = 3$, $|\mu| = M_S$, and $A_q = 0$. Mass insertion parameters and phase factor are assumed to be $|(\delta_{LL})_{qq_3}| = |(\delta_{RR})_{q_3 q}| = 1/3$ and $\sin\theta_q = 1/\sqrt{2}$, respectively.

to those of down quark. Here, we take $M_{\tilde{g}} = 3$ TeV, $\tan\beta = 3$, $|\mu| = M_S$, and $A_q = 0$.[2] In addition, the mass insertion parameters and the phase factor are assumed to be $|(\delta_{LL})_{qq_3}| = |(\delta_{RR})_{q_3 q}| = 1/3$ and $\sin\theta_q = 1/\sqrt{2}$ with $\theta_q \equiv \mathrm{Arg}[\mu(\delta_{LL})_{qq_3}(\delta_{RR})_{q_3 q}]$, respectively. From this figure, we find that the CEDMs dominate the EDMs, though the latter are not negligible at all. Further, the contribution of up quark is larger than that of down quark in the case of low $\tan\beta$, which is favored from the viewpoint of the 126 GeV Higgs mass in the high-scale SUSY scenario [11–14]. We would like to remark that the EDMs and CEDMs are proportional to the gluino mass except for the renormalization factors, and thus their values corresponding to other gluino masses are readily obtained by means of the scaling law, as long as $M_{\tilde{g}} \ll M_S$.

By using the EDMs and CEDMs computed above, we finally calculate the neutron EDM d_n. To that end, we need to express the neutron EDM in terms of d_q and $\tilde{d}_q$. At present, only the calculations based on the QCD sum-rules [53, 54] include both of these contributions on an equal footing. Their theoretical error is, however, still significant, though partial use of lattice results for the low-energy QCD constants may reduce the uncertainty [54]. Moreover, this approach seems to lack the strange quark contributions. For instance, when one imposes the Peccei-Quinn symmetry, the strange CEDM contribution to the neutron EDM completely vanishes in the case of the sum-rule calculations, while it is expected to be sizable from the estimation based on the chiral perturbation theory [55]. At this moment, both methods have large uncertainty and no consensus has been reached

[2] In the anomaly mediation, the A-terms are suppressed by one-loop factors, thus negligible in our calculation.

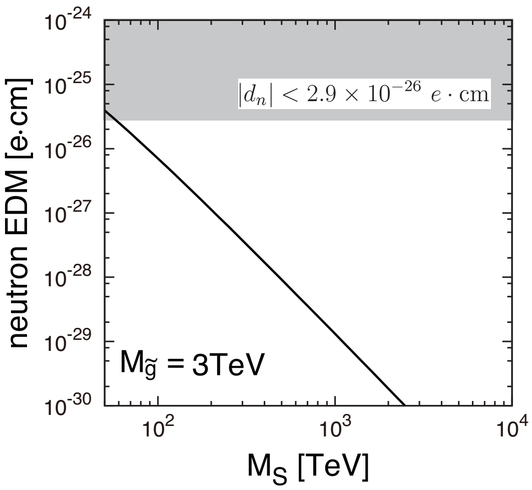

Figure 4: Neutron EDM d_n as a function of M_S. The same parameters are used as those exploited in Fig. 3. Shaded region represents the current experimental limit $|d_n| < 2.9 \times 10^{-26}\ e \cdot \mathrm{cm}$ [56].

yet. We strongly anticipate that the lattice simulations will evaluate the neutron EDM induced by the quark EDMs and CEDMs with high accuracy. In the present calculation, we use the result presented in Ref. [54]:[3]

$$d_n = 0.79 d_d - 0.20 d_u + e(0.30\tilde{d}_u + 0.59\tilde{d}_d) \ , \tag{22}$$

where we assume the Peccei-Quinn mechanism. In Fig. 4, we plot the resultant neutron EDM as a function of M_S. In this figure, we use the same parameters as those exploited in Fig. 3. The shaded region represents the current experimental limit $|d_n| < 2.9\times 10^{-26}\ e\cdot$cm [56]. As seen from this figure, the present experimental limit has already excluded the squark mass scale nearly up to 10^2 TeV. Future experiments of the neutron EDM are expected to reach $\sim 10^3$ TeV, which covers most of the region favored from the high-scale SUSY scenario compatible with the 126 GeV Higgs mass and the existence of $\mathcal{O}(1)$ TeV gauginos. Hence, the EDM experiments are quite promising, and may be about to grasp the signature of supersymmetry.

In the case of the minimal flavor violation discussed in the previous section, on the other hand, the predicted neutron EDM lies around $d_n \simeq 10^{-30}\ e \cdot$cm for $M_S = 10^2$ TeV, which is much below the current experimental limit.

[3] The numerical values presented here are in fact different from those in Ref. [54] by nearly a factor of two. The difference results from the use of different values for the quark condensate; We use $\langle \bar{q}q \rangle = -m_\pi^2 f_\pi^2/(m_u + m_d) \simeq -(262\ \mathrm{MeV})^3$ [3] while $\langle \bar{q}q \rangle = -(225\ \mathrm{MeV})^3$ is used in Ref. [54].

5 Conclusion and discussion

In this paper, we discuss the QCD corrections to the dimension-five CP-violating operators in the case of the high-scale supersymmetry. To appropriately evaluate the radiative corrections in the presence of a large hierarchy between the squark and gluino mass scales, we exploit the RGEs (including CP violating gluino-quark four-Fermi operators) in an effective theory where only the SM particles and gluinos are taken into account. As a result, the values of the low-energy quark EDMs and CEDMs may differ from those evaluated in previous works by $\mathcal{O}(100)$ % and $\mathcal{O}(10)$ %, respectively.

In the high-scale SUSY scenario, similar calculations based on the RGEs may have significant consequences for the prediction of other low-energy observables, such as gluino decay rates [57–59], particle-antiparticle mixing, rare and CP-violating decays, and so on. Even though these processes are often induced by the flavor-changing operators, a lot of our results are applicable to the cases since gluinos as well as photons and gluons do not distinguish quark flavors.

In the above calculation, we have only included the leading order effects, though the one-loop short-distance correction is also discussed. Before concluding this article, let us discuss possible ways of improvement of the above calculation. A straightforward improvement is achieved if one uses the two-loop RGEs as well as a complete set of one-loop threshold corrections. In addition, to go beyond the leading order analysis, we also need to include the operators which we neglect in our calculation; the gluino CEDM, four-quark operators, four-gluino operators, and the Weinberg operator. These operators mix with each other as well as with the quark EDMs and CEDMs during the RGE flow. A complete calculation beyond the leading order will be carried out on another occasion [60].

Note Added: While this work was being finalized, we realized the authors in Ref. [61] estimated the anomalous dimensions for the quark-gluino four-Fermi operators in a similar context. The results presented in the reference are, however, inconsistent with ours. Especially, the authors insist that they have not found the mixing among the four-Fermi operators by their explicit calculation, though we do as shown in Eq. (8).

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