Neutron and electron electric dipole moments in supersymmetric theories

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The electric dipole moments (EDMs) of the neutron and the electron in the supersymmetric standard models are discussed in detail taking systematically all the one-loop contributions into account. In the framework of grand unified theories coupled to $N=1$ supergravity, the contribution from the chargino-loop diagram is larger than those from the gluino- and/or neutralino-loop diagrams in wide ranges of the supersymmetric parameter space. We assume that *CP*-violating phases of the basic parameters are not suppressed in an unnatural way. The experimental limits of the EDMs could give constraints on the values of the supersymmetric mass parameters. From the neutron EDM, the lower bounds of ~ 1 TeV are obtained for the squark masses, whereas the masses of ~ 100 GeV are allowed for the charginos and the neutralinos. The electron EDM similarly gives the lower bounds of ~ 1 TeV on the slepton masses.

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I. INTRODUCTION

A nonvanishing value for the electric dipole moment (EDM) of an elementary particle could be observed only if *CP* invariance is violated [1]. Various searches for the EDMs have been made up to now to give limits on their values. In particular, the experimental bounds on the EDMs of the neutron d_n and the electron d_e are fairly small at present as $|d_n| < 10^{-25} e$ cm [2] and $|d_e| < 10^{-26} e$ cm [3], which makes it possible to impose constraints on some theories.

In the supersymmetric standard models it has been know from early days [4] that, in general, the EDM of the neutron is predicted to have a large value, typically of order 10^{-22}–$10^{-24} e$ cm, if the masses of the supersymmetric particles, such as the gauginos and the squarks, are of the order of 100 GeV [4–7]. Because the predicted values are larger than the experimental upper bounds, it was claimed that *CP*-violating phases of the parameters inherent in the models would be quite small or would vanish at the unification scale. However, this conclusion seems to be premature, because the supersymmetric particles might not be as light as 100 GeV. If they are much heavier than 100 GeV, say 1 TeV, the EDM could become smaller than the experimental upper bounds [8]. In fact, no signal for the supersymmetric particles has been found by now in the high-energy collider experiments at the CERN e^+e^- collider LEP or Fermilab Tevatron. The lower mass limits of the gluino and squarks are now set at about 150 GeV by the Collider Detector at Fermilab (CDF) Collaboration [9]. The supersymmetric particles are surely heavier than what was thought before. We must have an open mind regarding the mass values of the supersymmetric particles.

We reconsider in this paper the EDMs of the neutron [4–7] and the electron [10,11] in the supersymmetric standard models. Our approach to the problem of the EDMs is the following: We do not impose an unnatural assumption that the *CP*-violating phases are quite small or vanish. Instead the supersymmetric particles have to be as heavy as the predictions of the EDMs do not conflict with experiments. Taking into account all one-loop contributions, we discuss the allowed regions of the supersymmetric parameter space, and show how the masses of the supersymmetric particles are constrained [12].

In Sec. II, we shall review the EDMs from the one-loop effects of the supersymmetric particles, and give the expressions without any approximation. The EDM of a quark (lepton) receives contributions at the one-loop level from the diagrams in which propagate the squarks $\tilde{q}$ (sleptons $\tilde{l}$) and one of the charginos ω_i, neutralinos χ_j, and gluinos $\tilde{g}$ as shown in Fig. 1. These diagrams exist because the charginos and the neutralinos are the mixed states of the gauginos and the Higgsinos, and the squarks (sleptons) are the mixed states of the scalar left-handed quark $\tilde{q}_L$ (lepton $\tilde{l}_L$) and the scalar right-handed quark $\tilde{q}_R$ (lepton $\tilde{l}_R$). As has been discussed recently [13], there are also two other operators which could sizably contribute to the EDM of the neutron: Weinberg's gluonic operator of dimension six and the quark chromoelectric dipole moment operator of dimension five. In this paper, however, we consider only the EDM operator of dimension five stated above, since this operator could give the largest contribution in the supersymmetric standard models [14].

In order to precisely obtain the mixing angles, and correctly evaluate the effects of the parameters, we have

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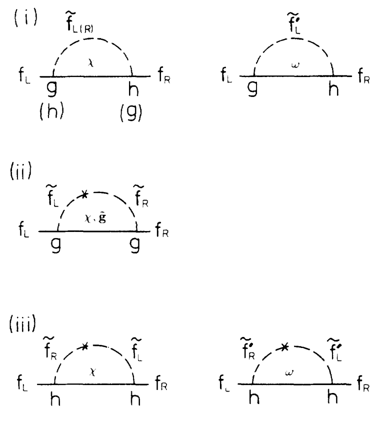

FIG. 1. The Feynman diagrams of the SUSY contributions to the EDM of a quark, g and h being coupling constants for gauge and Yukawa interactions. The photon line should be attached to the internal line of the squarks or the charginos.

to diagonalize the complex mass matrices for the charginos, the neutralinos, and the squarks (sleptons) explicitly without any assumption on the relative magnitudes of the parameters. This is done in Sec. III. For the EDMs due to the charginos and the gluinos, the formulas will be given in closed forms. We shall also briefly discuss in this section how the EDMs from the supersymmetric (SUSY) sector can be suppressed by the heavy masses of the squarks and the sleptons. Their $\tan\beta$ dependence shall be also discussed.

In Sec. IV, the EDMs of the neutron and the electron will be numerically evaluated. The experimental limits give constraints on the supersymmetric parameters, from which we obtain the allowed mass ranges for the supersymmetric particles. In particular, we explore the possibility for these particles to have masses less than 1 TeV which would be accessible at the next e^+e^- and/or pp colliders.

Section V shall be devoted to the conclusion.

II. ELECTRIC DIPOLE MOMENT

The supersymmetric standard models based on $N=1$ supergravity [15] contain generally several complex parameters in addition to the Yukawa coupling constants, which are possible new sources of CP violation. For instance, in the models with minimal particle content, complex could be the mass parameters of SU(3), SU(2), and U(1) gauginos $\tilde{m}_3, \tilde{m}_2$, and $\tilde{m}_1$, the mass parameter m_H in the bilinear term of Higgs superfields, and the dimensionless parameters in the trilinear and bilinear terms of scalar fields A's and B. (For our conventions, see Ref. [16]). These lead to complex mass terms for the supersymmetric particles. After the SU(2)×U(1) gauge symmetry is broken, the mass matrices M^- and M^0 for the charginos and the neutralinos become

$$M^- = \begin{pmatrix} \tilde{m}_2 & -gv_1^*/\sqrt{2} \\ -gv_2^*/\sqrt{2} & m_H \end{pmatrix}, \tag{1}$$

$$M^0 = \begin{pmatrix} \tilde{m}_1 & 0 & g'v_1^*/2 & -g'v_2^*/2 \\ 0 & \tilde{m}_2 & -gv_1^*/2 & gv_2^*/2 \\ g'v_1^*/2 & -gv_1^*/2 & 0 & -m_H \\ -g'v_2^*/2 & gv_2^*/2 & -m_H & 0 \end{pmatrix}, \tag{2}$$

where v_1 and v_2 are the vacuum expectation values of the two Higgs doublets with U(1) hypercharges $-\frac{1}{2}$ and $\frac{1}{2}$, respectively. The masses of the charginos and the neutralinos are obtained by diagonalizing M^- and M^0 as

$$C_R^\dagger M^- C_L = \mathrm{diag}(m_{\omega 1}, m_{\omega 2}), \tag{3}$$

$$N^t M^0 N = \mathrm{diag}(m_{\chi 1}, m_{\chi 2}, m_{\chi 3}, m_{\chi 4}). \tag{4}$$

The mass of the gluino is given by

$$m_{\tilde{g}} = |\tilde{m}_3|. \tag{5}$$

The mass-squared matrix $\tilde{M}_f^2$ for the scalar partners of a quark or a lepton with flavor f becomes

$$\tilde{M}_f^2 = \begin{pmatrix} m_f^2 + \cos2\beta(T_{3f} - Q_f \sin^2\theta_W) M_Z^2 + M_{\tilde{f}L}^2 & m_f(R_f m_H + A_f^* m_{3/2}) \\ m_f(R_f^* m_H^* + A_f m_{3/2}) & m_f^2 + \cos2\beta Q_f \sin^2\theta_W M_Z^2 + M_{\tilde{f}R}^2 \end{pmatrix}, \tag{6}$$

where m_f represents a mass of the fermion f, Q_f an electric charge of f, and T_{3f} the third component of the weak isospin of left-handed f. $m_{3/2}$ is the gravitino mass, and $M_{\tilde{f}L}^2$ and $M_{\tilde{f}R}^2$ are the mass-squared parameters for $\tilde{f}_L$ and $\tilde{f}_R$. R_f and $\tan\beta$ are defined by

$$R_f = \begin{cases} \dfrac{v_1}{v_2^*} & (\text{for } T_{3f} = \frac{1}{2}), \\ \dfrac{v_2}{v_1^*} & (\text{for } T_{3f} = -\frac{1}{2}), \end{cases}$$

$$\tan\beta = \left|\frac{v_2}{v_1}\right|. \tag{7}$$

Each mass matrix (6) for the squarks and the sleptons is diagonalized by the unitary matrix S_f as

$$S_f^\dagger \tilde{M}_f^2 S_f = \mathrm{diag}(M_{\tilde{f}1}^2, M_{\tilde{f}2}^2). \tag{8}$$

We have neglected flavor mixing here, but it is straightforward to incorporate flavor mixing by interpreting each element of the matrix (6) as the matrix in flavor space. Since all the complex phases of the parameters cannot be absorbed by redefining the fields, several interactions have complex coupling constants when expressed in terms of the mass eigenstates, leading to the violation of CP invariance.

Before we calculate the EDM from the SUSY sector, it should be noted that the EDM operator needs to flip the chirality of the quark or the lepton. In the supersymmetric standard models the gauginos couple the quark (lepton) to the squark (slepton) with the same chiral assignment via the gauge interactions, while the Higgsinos couple the quark (lepton) to the squark (slepton) with the opposite chiral assignment via the Yukawa interactions. Therefore the EDM can arise at the one-loop level from three origins.

(i) One vertex of the loop diagram is due to the gauge interaction and the other is Yukawa. The gaugino and the Higgsino are mixed.

(ii) The two vertices are both due to the gauge interactions. The scalars $\tilde{q}_L(\tilde{l}_L)$ and $\tilde{q}_R(\tilde{l}_R)$ are mixed.

(iii) The two vertices are both due to the Yukawa interactions. $\tilde{q}_L(\tilde{l}_L)$ and $\tilde{q}_R(\tilde{l}_R)$ are mixed.

The chargino diagram originates in (i) and (iii), the gluino diagram in (ii), and the neutralino diagram in all of these. Taking into account all the possibilities, the EDM of a quark or a lepton at the one-loop level is given by the following equations.

The chargino contribution:

$$d_f^C/e=\frac{\alpha_{\rm EM}}{4\pi\sin^2\theta_W}\sum_{k=1}^{2}\sum_{i=1}^{2}{\rm Im}(F_{f_i}^k)\frac{m_{\omega i}}{M_{\tilde{f}'k}^2}\left[Q_{\tilde{f}'}I\left[\frac{m_{\omega i}^2}{M_{\tilde{f}'k}^2},\frac{m_f^2}{M_{\tilde{f}'k}^2}\right]+(Q_f-Q_{\tilde{f}'})J\left[\frac{m_{\omega i}^2}{M_{\tilde{f}'k}^2},\frac{m_f^2}{M_{\tilde{f}'k}^2}\right]\right], \tag{9}$$

$$F_{f_i}^k=\begin{cases}(C_{L1i}S_{\tilde{f}'1k}^*-\kappa_{f'}C_{L2i}S_{\tilde{f}'2k}^*)\kappa_fC_{R2i}^*S_{f'1k} & (\text{for } T_{3f}=\frac{1}{2}),\\ (C_{R1i}^*S_{\tilde{f}'1k}^*-\kappa_{f'}C_{R2i}^*S_{\tilde{f}'2k}^*)\kappa_fC_{L2i}S_{f'1k} & (\text{for } T_{3f}=-\frac{1}{2}).\end{cases} \tag{10}$$

In the above equations f' represents a flavor of a quark or a lepton whose left-handed component forms an SU(2) doublet with the left-handed component of f, and κ_f is defined by

$$\kappa_f=\frac{m_f}{\sqrt{2}M_W}\frac{\sqrt{|v_1|^2+|v_2|^2}}{v_a}, \tag{11}$$

where $a=2$ for $T_{3f}=\frac{1}{2}$, and $a=1$ for $T_{3f}=-\frac{1}{2}$.

The neutralino contribution:

$$d_f^N/e=\frac{\alpha_{\rm EM}}{4\pi\sin^2\theta_W}\sum_{k=1}^{2}\sum_{j=1}^{4}{\rm Im}(G_{f_j}^k)\frac{m_{\chi j}}{M_{\tilde{f}k}^2}Q_{\tilde{f}}I\left[\frac{m_{\chi j}^2}{M_{\tilde{f}k}^2},\frac{m_f^2}{M_{\tilde{f}k}^2}\right], \tag{12}$$

$$G_{f_j}^k=\{\sqrt{2}[\tan\theta_W(Q_f-T_{3f})N_{1j}+T_{3f}N_{2j}]S_{\tilde{f}1k}^*+\kappa_fN_{bj}S_{\tilde{f}2k}^*\}(\sqrt{2}\tan\theta_WQ_fN_{1j}S_{f2k}-\kappa_fN_{bj}S_{f1k}), \tag{13}$$

where $b=4$ for $T_{3f}=\frac{1}{2}$, and $b=3$ for $T_{3f}=-\frac{1}{2}$.

The gluino contribution:

$$d_q^G/e=\frac{2\alpha_S}{3\pi}\sum_{k=1}^{2}{\rm Im}(H_q^k)\frac{m_{\tilde{g}}}{M_{\tilde{q}k}^2}Q_{\tilde{q}}I\left[\frac{m_{\tilde{g}}^2}{M_{\tilde{q}k}^2},\frac{m_q^2}{M_{\tilde{q}k}^2}\right], \tag{14}$$

$$H_q^k=e^{-i\arg(\tilde{m}_3)}S_{q1k}^*S_{q2k}. \tag{15}$$

The functions $I(r,s)$ and $J(r,s)$ come from the one-loop integrals. The experimental lower bounds on the squark and slepton masses are much larger than the quark and lepton masses except the top quark. Thus, as far as the mixing between the first generation and the third generation can be neglected, $s(=m_f^2/M_{\tilde{f}k}^2)$ almost vanishes. Putting $s=0$, these functions are expressed as

$$\begin{aligned}I(r,0)&=\frac{1}{2(1-r)^2}\left[1+r+\frac{2r}{1-r}\ln r\right]\quad[\equiv I(r)],\\ J(r,0)&=\frac{1}{2(1-r)^2}\left[3-r+\frac{2}{1-r}\ln r\right]\quad[\equiv J(r)].\end{aligned} \tag{16}$$

To get the EDM of the neutron itself from the above quark EDMs, we have to invoke some model on the structure of the neutron. Although there might be a more sophisticated treatment of its structure, we use the nonrelativistic quark model in this paper as the first-order approximation of the neutron to get

$$d_n=\tfrac{1}{3}(4d_d-d_u). \tag{17}$$

As discussed before, the supersymmetric contributions to the EDM can be attributed to the three origins (i), (ii), and (iii). This is easily seen if we carefully look at the mixing factors (10), (13), and (15). Since the suppression factors coming from the gaugino-Higgsino mixing and the $\tilde{f}_L$-$\tilde{f}_R$ mixing are of $\sim M_W/m_\omega$ and $\sim m_f/M_{\tilde{f}}$, the products of the coupling constants and the suppression factors in (i), (ii), and (iii) become roughly of $\sim\alpha m_f/m_\omega$, $\sim\alpha m_f/M_{\tilde{f}}$, and $\sim\alpha m_f^2/M_Wm_\omega$, respectively, α being an appropriate fine structure constant. Therefore the contribution from (iii) is much smaller than those from (i) and (ii) and can be safely neglected, whereas both contributions from (i) and (ii) could have the same magnitude or dominate over the other depending on the masses of ω and $\tilde{f}$. It will be shown later that, in the grand unified theory (GUT) scheme, $|d_q^C|$ coming from mainly (i) becomes larger than $|d_q^G|$ coming from (ii) in some interesting parameter regions, although $|d_q^G|$ was considered to be dominant in much of the literature. Since the coupling strengths of the neutralinos are smaller than those of the gluinos and the charginos, $|d_f^N|$ coming from mainly (i) and (ii) is almost always the smallest among all contributions.

III. SIMPLIFIED AND CLOSED FORMS OF EDMs

To get an overview of the dependence of the supersymmetric contributions to the EDMs on the parameters, let us obtain simpler and closed expressions of the EDMs with sensible reasonings and reasonable approximations. For the phases of the fields there are some degrees of freedom by which the phases of the parameters can be altered. As discussed in Refs. [6,7], in the minimal supersymmetric standard model there remain two physical phases in addition to the phase of the (super) Kobayashi-Maskawa matrix [17]. Thus, in order to avoid unphysical redundancy we adopt hereafter the phase convention that $\tilde{m}_3$, $\tilde{m}_2$, $\tilde{m}_1, v_1$, and v_2 are real and positive. Then the remaining two complex parameters relevant to our discussion are m_H and A_f, which we parametrize as

$$m_H=|m_H|e^{i\theta}\ , \quad A_f=|A_f|e^{i\alpha_f}\ . \tag{18}$$

The diagonalization of the matrices M^- and $\tilde{M}_f^2$ in Eqs. (1) and (6) can be made analytically without any approximation. (See Appendix A.) Neglecting the tiny contribution from the origin (iii), the chargino contribution to the EDM is finally written as

$$\begin{aligned} d_f^C/e &= \frac{\alpha_{EM}}{4\pi\sin^2\theta_W}\sin\theta R_f \frac{\tilde{m}_2|m_H|}{(m_{\omega 2}^2-m_{\omega 1}^2)}\frac{m_f}{M_{\tilde{f}'}^2}\sum_{i=1}^{2}(-1)^i\left[Q_{\tilde{f}'}I\left[\frac{m_{\omega i}^2}{M_{\tilde{f}'}^2}\right]+(Q_f-Q_{\tilde{f}'})J\left[\frac{m_{\omega i}^2}{M_{\tilde{f}'}^2}\right]\right]\ ,\\ m_{\omega 1(2)}^2 &= \tfrac{1}{2}[\tilde{m}_2{}^2+|m_H|^2+2M_W^2-(+)\sqrt{D}\ ]\ ,\\ D &= (\tilde{m}_2{}^2+|m_H|^2+2M_W^2)^2-4(\tilde{m}_2{}^2|m_H|^2+\sin^2 2\beta M_W^4-2\sin 2\beta M_W^2\tilde{m}_2|m_H|\cos\theta)\ . \end{aligned} \tag{19}$$

Here we have approximated the masses of two mass eigenstates of the squarks or the sleptons to be $M_{\tilde{f}1}\simeq M_{\tilde{f}2}(\equiv M_{\tilde{f}})$, which is given by

$$M_{\tilde{f}}^2=\tfrac{1}{2}\mathrm{Tr}(\tilde{M}_f^2)\ . \tag{20}$$

The gluino contribution is also written in a simpler form if we substitute the average squark mass (see Appendix B):

$$\begin{aligned} d_q^G/e &= \frac{2\alpha_s}{3\pi}\left[\sin\alpha_q|A_q|-\sin\theta R_q\frac{|m_H|}{m_{3/2}}\right]\\ &\quad\times\frac{m_{3/2}}{M_{\tilde{q}}}\frac{m_q}{M_{\tilde{q}}^2}Q_{\tilde{q}}\frac{m_{\tilde{g}}}{M_{\tilde{q}}}K\left[\frac{m_{\tilde{g}}^2}{M_{\tilde{q}}^2}\right]\ ,\\ K(r) &= \frac{-1}{2(1-r)^3}\left[1+5r+\frac{2r(2+r)}{1-r}\ln r\right]\ . \end{aligned} \tag{21}$$

Concerning the neutralino contribution we have to numerically diagonalize the mass matrix M^0 of Eq. (2), although an analytical calculation may be possible in some interesting regions of the parameter space [18].

From the above simplified expressions, we can roughly see that the supersymmetric contributions to the EDMs can lie under the experimental upper bounds if the squarks and sleptons are heavier than 1 TeV. For example, the factor in the chargino contribution of Eq. (19) $(\alpha_{EM}/4\pi\sin^2\theta_W)(m_f/M_{\tilde{f}'}^2)$ represents a typical value of d_f^C/e, which is written as

$$\frac{\alpha_{\mathrm{EM}}}{4\pi\sin^2\theta_W}\frac{m_f}{M_{\tilde{f}'}^2}=5.0\times10^{-25}\left[\frac{1\ \mathrm{TeV}}{M_{\tilde{f}'}}\right]^2\left[\frac{m_f}{10\ \mathrm{MeV}}\right]\mathrm{cm}\ . \tag{22}$$

Of course, to know the precise value of d_f^C, we have to evaluate the functions $I(r)$ and $J(r)$. The values of these functions vary $I(r)=5\times10^{-1}$–5×10^{-3}, $J(r)=(-3)$–(-5×10^{-3}) for $r=10^{-2}$–10^2. Thus, for the squark masses $M_{\tilde{q}}\gtrsim 1$ TeV, the quark EDM d_q^C can become smaller than the experimental bound on $|d_n|$ of $1.2\times10^{-25}e$ cm even if the *CP* violation is maximal. Though the experimental bound on $|d_e|$ is smaller than those on $|d_n|$ by one order of magnitude, the EDM of the electron similarly constrains the slepton masses as $M_{\tilde{l}}\gtrsim 1$ TeV, since $m_e/m_q\sim10^{-1}$. One might worry about the gluino contribution, because the factor $2\alpha_s/3\pi$ is about 10 times larger than $\alpha_{\mathrm{EM}}/4\pi\sin^2\theta_W$. However, the function $\sqrt{r}K(r)$ takes value $-1\times10^{-1}<\sqrt{r}K(r)<-2\times10^{-3}$ for $r=10^{-2}$–10^2, so that the squark masses $M_{\tilde{q}}\gtrsim 1$ TeV reduce the gluino contribution to d_n below the experimental upper bound. Thus the maximal *CP* violation can be allowed for the TeV squarks.

Before closing this section, we comment on the $\tan\beta$ dependence of the EDMs from the SUSY origin. The chargino contribution d_f^C, Eq. (19), has a factor R_f, which is $1/\tan\beta$ for the u quark and $\tan\beta$ for the d quark and the electron. Also d_f^C is affected by $\tan\beta$ indirectly through the chargino masses and the squark or slepton masses, but these effects are not so strong. Therefore the electron EDM due to the charginos is approximately proportional to $\tan\beta$. On the other hand, the neutron EDM consists of the EDMs of the u quark and the d quark as shown in Eq. (17), which might make the $\tan\beta$ dependence of d_n^C more complicated. However, since m_d is about twice larger than m_u, the multiplicative factor of d_d^C in Eq. (17) is four times larger than d_u^C. Moreover, $\tan\beta$ is larger than $1/\tan\beta$ (i.e., $\tan\beta>1$) if the $SU(2)\times U(1)$ symmetry is broken through radiative corrections [19]. Then d_n^C is dominated by d_d^C, and d_n^C is approximately given by $\frac{4}{3}d_d^C$. Thus the neutron EDM due to the charginos is roughly proportional to $\tan\beta$. As shown later, both the electron and the neutron EDMs get the largest contribution from the chargino-loop diagrams in wide parameter regions allowed by the experiments. In these regions, d_n and d_e due to the SUSY contributions are roughly proportional to $\tan\beta$, and the larger $\tan\beta$ is, the bigger d_n and d_e are predicted to be in the supersymmetric standard models.

IV. NUMERICAL RESULTS

We now discuss numerically the supersymmetric contributions to the EDM of the neutron. To make the numerical calculations definite, we need to set several theoretically and experimentally undetermined parameters. As a typical example of the natural magnitudes for the *CP*-violating phases, we simply take $\theta=\alpha_u=\alpha_d=\pi/4$. The dimensionless parameters A_f are fixed as $|A_u|=|A_d|=1$. For the mass parameters let us first make some assumptions based on theoretically plausible GUT models with $N=1$ local supersymmetry broken spontaneously. In this scheme the mass parameters of the gauginos satisfy the relation $(g^2/g_s^2)\tilde{m}_3=\tilde{m}_2=(3g^2/5g'^2)\tilde{m}_1$. The mass parameters $M_{\tilde{f}L}$ and $M_{\tilde{f}R}$ appearing in Eq. (6) could be related to the gravitino mass and the gaugino masses. In the ordinary scheme for the mass generation, the gaugino masses are smaller than or around the gravitino mass, so that the scale characteristic of $M_{\tilde{f}L}$ and $M_{\tilde{f}R}$ is given by $m_{3/2}$. For simplicity, we take $M_{\tilde{u}L}=M_{\tilde{u}R}=M_{\tilde{d}L}=M_{\tilde{d}R}=m_{3/2}$. Then the masses of the squarks can be estimated approximately by $m_{3/2}$. The magnitude of the Higgsino mass parameter $|m_H|$ should be at most of $\sim m_{3/2}$ for correctly breaking the $SU(2)\times U(1)$ symmetry.

In Fig. 2 we show the chargino contribution to the EDM of the neutron as a function of $m_{3/2}$ (1 TeV $\leq m_{3/2} \leq 10$ TeV) for (a) $\tan\beta=2$ and (b) $\tan\beta=10$. The values of the mass parameters are shown in Table I [20]. In these parameter regions, except $M_{3/2}\simeq 1$ TeV for case (iii.b) [21], the gluino and the neutralino contributions are smaller than the chargino one, giving the EDM of the neutron by $d_n \simeq d_n^C$. Thus, not to conflict with the experimental bound of $|d_n|<10^{-25}e$ cm, we must have $m_{3/2}\gtrsim 3$ TeV if $\tan\beta=2$, and $m_{3/2}\gtrsim 7$ TeV if $\tan\beta=10$. This difference caused by $\tan\beta$ is easily understood since d_n^C is roughly proportional to $\tan\beta$ as discussed in the preceding section. So when $\tan\beta$ increases, the lower bounds on the squark masses become larger. The difference between the three contributions due to the gluinos, the charginos, and the neutralinos can be seen in Fig. 3, where $|d_n^G|$, $|d_n^C|$, and $|d_n^N|$ are shown as functions of $m_{3/2}$ for $\tan\beta=2$ and $\tilde{m}_2=|m_H|=0.5$ TeV. This example clearly shows that the charginos really give the largest contribution to d_n. Since d_n^C is dominant, the EDM of the neutron is roughly proportional as a whole to $\sin\theta$ as seen from Eq. (19), and does not depend much on α_u and α_d.

TABLE I. Four sets of $\tilde{m}_2$ and $|m_H|$.

	(i)	(ii)	(iii)	(iv)
$\tilde{m}_2$ (TeV)	0.2	1.0	0.2	1.0
$\|m_H\|$ (TeV)	0.2	0.2	1.0	1.0

Although the squarks have to be heavier than ~ 1 TeV, the charginos and the neutralinos could be lighter than ~ 1 TeV [22]. Even for relatively small values of $\tilde{m}_2$ and $|m_H|$, such as 200 GeV, the neutron EDM can lie within the experimental bounds. This is more clearly shown in Fig. 4, where we show the parameter region which leads to $|d_n|<10^{-25}e$ cm in the $(\tilde{m}_2,|m_H|)$ plane for $m_{3/2}=1,3,5$ TeV and $\tan\beta=2,10$. The curves are the contours of $|d_n^C|=10^{-25}e$ cm, and the inside of each curve is excluded. A sizable part in the region of $\tilde{m}_2,|m_H|<1$ TeV is allowed in the case of $\tan\beta=2$ and $m_{3/2}=3$ TeV. Thus searching for the charginos and the neutralinos seems hopeful in the near future e^+e^- and/or pp experiments, while it will be quite difficult to find the squarks. The gluino mass should be larger than about 0.5 TeV according to the GUT relation among the gaugino mass parameters.

For examining further possibilities of light supersymmetric particles, let us relax the assumptions on the mass parameters mentioned before. If we do not assume the GUT relation among the gaugino masses, the gluino mass is constrained by comparing the gluino contribution d_n^G with the experimental upper bound on d_n. In Fig. 5 $|d_n^G|$

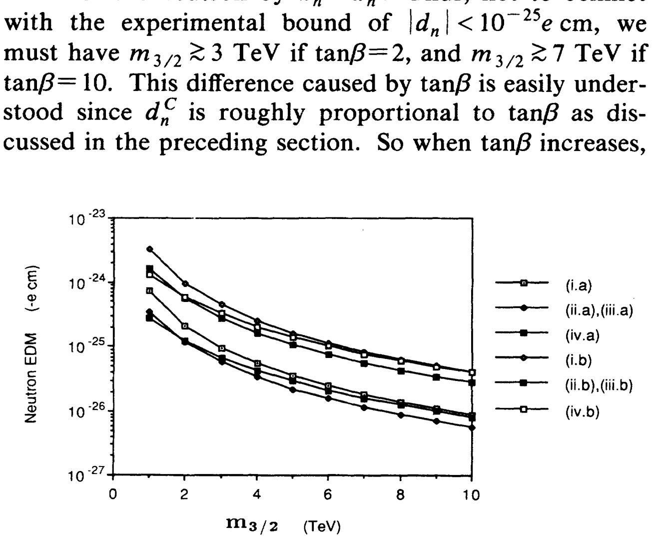

FIG. 2. The neutron EDM from the chargino contribution for the parameter values shown in Table I with (a) $\tan\beta=2$ and (b) $\tan\beta=10$.

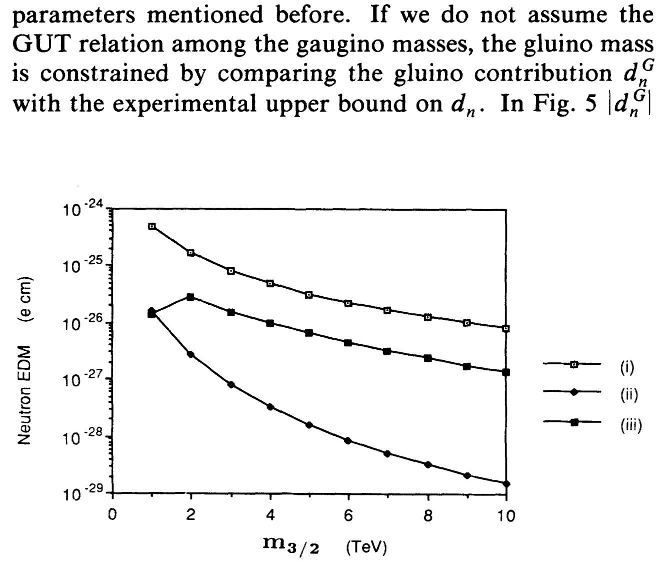

FIG. 3. (i) The chargino, (ii) neutralino, and (iii) gluino contributions to the neutron EDM for $\tan\beta=2, \tilde{m}_2=|m_H|=0.5$ TeV. The absolute values of the EDMs are shown. The signs are negative for curve (i), and positive for curves (ii) and (iii).

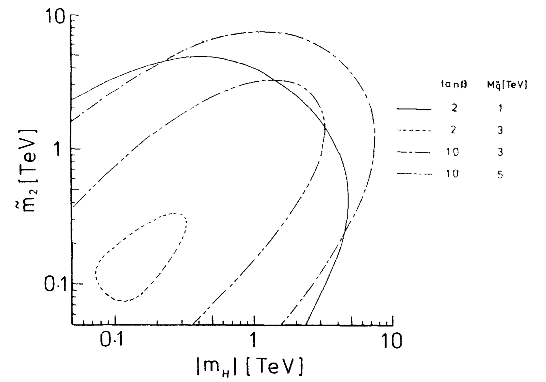

FIG. 4. The parameter regions compatible with the present experimental upper bound on the neutron EDM.

is shown as a function of $m_{3/2}$ for $m_{\tilde{g}}=0.2$ TeV, $|m_H|=0.2,1$ TeV, and $\tan\beta=2,10$. The ~ 100 GeV gluinos become allowed for the squarks heavier than a few TeV. We should note that the gluino contribution is affected by two *CP*-violating phases α_q and θ as seen in Eq. (21). If the two terms proportional to $\sin\alpha_q$ and $\sin\theta$ cancel each other, the gluino contribution to the neutron EDM could be very small without suppressed *CP*-violating phases nor heavy supersymmetric particles. This, however, would be unnatural unless there is some profound symmetry forcing the two terms to cancel.

If $\tilde{m}_2$ and $|m_H|$ can have any values independently of $m_{3/2}$ or the squark masses, the squarks could be relatively light for large $\tilde{m}_2$ and $|m_H|$. We show the largest value among $|d_n^C|$, $|d_n^N|$, and $|d_n^G|$ for $m_{3/2}=0.2,1$ TeV as a function of $|m_H|$ ($\tilde{m}_2=1,3$ TeV) in Fig. 6 (top), and as a function of $\tilde{m}_2$ ($|m_H|=1,3$ TeV) in Fig. 6 (bottom). Since a smaller value of $\tan\beta$ gives a smaller value of $|d_n^C|$, we have taken $\tan\beta=2$ to look for the light squark possibility. In Fig. 6 (top), e.g., for $|m_H|=3$ TeV, the largest contribution to the EDM is made by d_n^G in $|m_H|\gtrsim 5$ TeV for $m_{3/2}=0.2$ TeV and in $|m_H|\gtrsim 7$ TeV for $m_{3/2}=1$ TeV, but by d_n^C in other regions. The gluino contribution increases with $|m_H|$ when $\tilde{m}_2$ is fixed. This is because the mixings of $\tilde{q}_L$ and $\tilde{q}_R$ become larger while the squark masses are not much changed and the gluino mass remains the same. In Fig. 6 (bottom), $|d_n^C|$ is the largest except in $\tilde{m}_2\lesssim 2$ TeV for $|m_H|=3$ TeV and $m_{3/2}=0.2$ TeV. These examples show that squarks as light as 0.2 TeV could be possible for sufficiently large $|m_H|$ and $\tilde{m}_2$. In these regions the squarks and the sleptons become lighter than the lightest neutralino, which might be disfavored from cosmology.

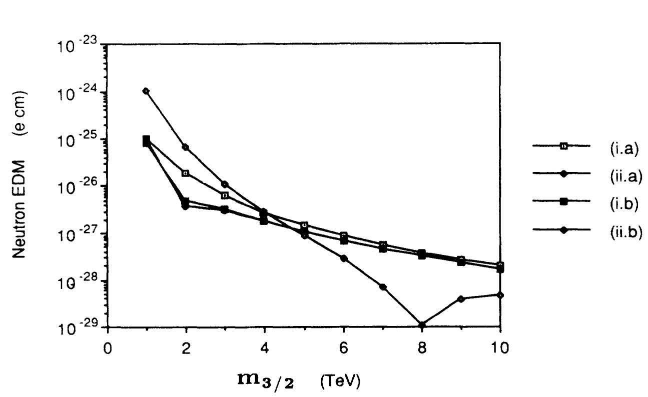

FIG. 5. The neutron EDM from the gluino contribution for $m_{\tilde{g}}=0.2$ TeV: (i) $|m_H|=0.2$ TeV, (ii) $|m_H|=1$ TeV; (a) $\tan\beta=2$, (b) $\tan\beta=10$. The absolute value of the EDM is shown. The sign is positive in $m_{3/2}\gtrsim 1$ TeV (i.a.), $m_{3/2}\gtrsim 2$ TeV (ii.a, i.b), and $m_{3/2}\gtrsim 8$ TeV (ii.b).

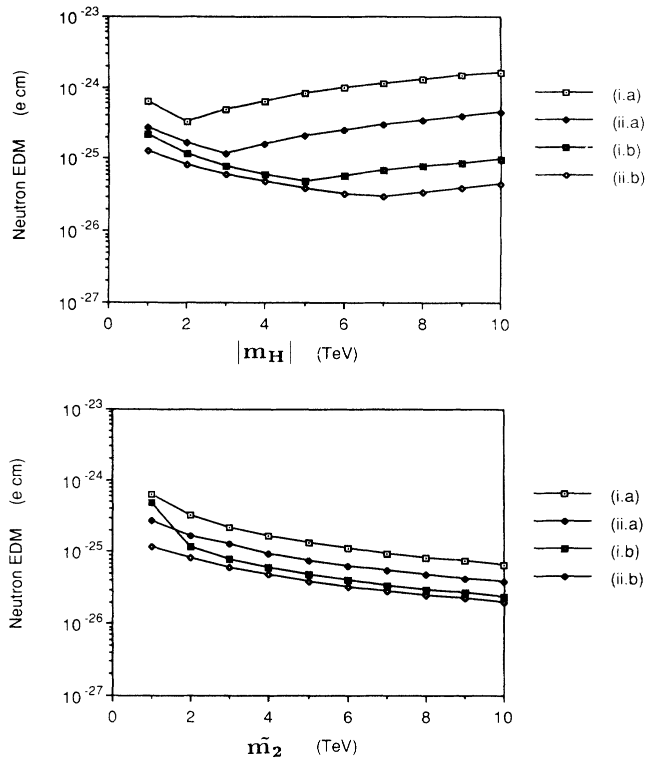

FIG. 6. (Top) The neutron EDM as a function of $|m_H|$ for $\tan\beta=2$; (a) $\tilde{m}_2=1$ TeV, (b) $\tilde{m}_2=3$ TeV; (i) $m_{3/2}=0.2$ TeV, (ii) $m_{3/2}=1$ TeV. (Bottom) The neutron EDM as a function of $\tilde{m}_2$ for $\tan\beta=2$: (a) $|m_H|=1$ TeV, (b) $|m_H|=3$ TeV; (i) $m_{3/2}=0.2$ TeV, (ii) $m_{3/2}=1$ TeV.

We finally discuss the EDM of the electron. In Fig. 7 the chargino contribution is shown assuming $\theta=\alpha_e=\pi/4$ and $|A_e|=1$. We have taken

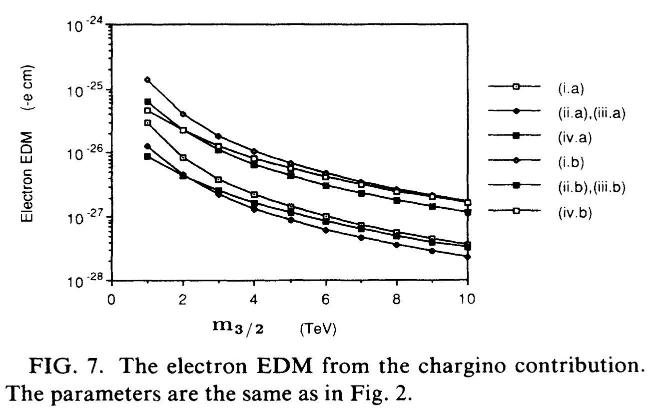

FIG. 7. The electron EDM from the chargino contribution. The parameters are the same as in Fig. 2.

$M_{\tilde{e}L} = M_{\tilde{e}R} = M_{\tilde{\nu}L} = m_{3/2}$, so that the slepton masses are approximately given by $m_{3/2}$. The values of the other parameters are the same as those for Fig. 2. In these parameter regions the neutralino contribution is much smaller than the chargino one. The electron EDM becomes $|d_e| < 10^{-26}e$ cm if $m_{3/2} \gtrsim 1$ TeV for $\tan\beta=2$ and $m_{3/2} \gtrsim 4$ TeV for $\tan\beta=10$. The values of a few hundred GeV are allowed for $\tilde{m}_2$ and $|m_H|$. We can see that similar results are obtained from the EDMs of the neutron and the electron.

V. CONCLUSION

Unless there is some symmetry, it would be natural to consider that the magnitude of possible *CP* violation is not much suppressed. Assuming that the arguments of the *CP*-violating phases are of order unity, we have studied constraints on the values of the parameters derived from the EDMs of the neutron and the electron. All the one-loop contributions to the EDMs due to the supersymmetric particles have been analyzed in detail. Under the assumption of GUTs the chargino contribution is dominant in wide ranges of the parameter space compatible with the experiments. Our study shows that the masses of the squarks and sleptons are likely to be larger than a few TeV, while $\tilde{m}_2$ and $|m_H|$ can be smaller by one order of magnitude. These results can be derived both from the neutron EDM and the electron EDM separately. The value of $\tan\beta$ has also a sizable effect on the EDMs. If $\tilde{m}_2$ and $|m_H|$ are less than ~ 1 TeV, the charginos and the neutralinos could have masses accessible at near future colliders. Then the *CP* violation may be observed in these experiments as, e.g., a *T*-odd asymmetry [12,16,23]. On the other hand, if $\tilde{m}_2$ and $|m_H|$ are larger than about 1 TeV, it will be hard to find supersymmetric particles in near future. In this case the supersymmetric standard models could only be indirectly examined by observing nonvanishing values for the neutron and electron EDMs and/or by detecting Higgs bosons [24].

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APPENDIX A

The chargino mass matrix M^- is diagonalized as follows: The phases of $\tilde{m}_2$, v_1, and v_2 are extracted as

$$M^- = P_R M_C P_L^\dagger ,$$
$$M_C = \begin{bmatrix} |\tilde{m}_2| & g|v_1|/\sqrt{2} \\ g|v_2|/\sqrt{2} & |m_H| e^{i(\theta_H+\theta_g+\theta_1+\theta_2)} \end{bmatrix} ,$$
$$P_R = \begin{bmatrix} e^{i\theta_g} & 0 \\ 0 & -e^{-i\theta_2} \end{bmatrix} , \tag{A1}$$
$$P_L = \begin{bmatrix} 1 & 0 \\ 0 & -e^{i(\theta_1+\theta_g)} \end{bmatrix} ,$$

where $\theta_H = \arg(m_H)$, $\theta_g = \arg(\tilde{m}_2)$, $\theta_1 = \arg(v_1)$, and $\theta_2 = \arg(v_2)$. The complex 2×2 matrix M_C can be readily diagonalized by unitary matrices $U_{R,L}$ as

$$U_R^\dagger M_C U_L = \begin{bmatrix} m_{\omega 1} & 0 \\ 0 & m_{\omega 2} \end{bmatrix} ,$$
$$U_R = \begin{bmatrix} \cos\theta_R & -e^{-i\beta_R}\sin\theta_R \\ e^{i\beta_R}\sin\theta_R & \cos\theta_R \end{bmatrix} \begin{bmatrix} e^{i\gamma_1} & 0 \\ 0 & e^{i\gamma_2} \end{bmatrix} , \tag{A2}$$
$$U_L = \begin{bmatrix} \cos\theta_L & -e^{-i\beta_L}\sin\theta_L \\ e^{i\beta_L}\sin\theta_L & \cos\theta_L \end{bmatrix} .$$

The angles are calculated to be

$$\tan 2\theta_R = \sqrt{2}g \frac{\sqrt{|v_2\tilde{m}_2|^2 + |v_1 m_H|^2 + 2|v_1 v_2 \tilde{m}_2 m_H|\cos\theta}}{|\tilde{m}_2|^2 - |m_H|^2 + g^2(|v_1|^2 - |v_2|^2)/2} ,$$
$$\tan 2\theta_L = \sqrt{2}g \frac{\sqrt{|v_1\tilde{m}_2|^2 + |v_2 m_H|^2 + 2|v_1 v_2 \tilde{m}_2 m_H|\cos\theta}}{|\tilde{m}_2|^2 - |m_H|^2 + g^2(|v_2|^2 - |v_1|^2)/2} ,$$
$$\tan\beta_R = \frac{\sin\theta}{\cos\theta + |v_2\tilde{m}_2|/|v_1 m_H|} ,$$
$$\tan\beta_L = -\frac{\sin\theta}{\cos\theta + |v_1\tilde{m}_2|/|v_2 m_H|} , \tag{A3}$$
$$\tan\gamma_1 = -\frac{\sin\theta}{\cos\theta + \frac{|\tilde{m}_2|}{|m_H|}\frac{2(m_{\omega 1}^2 - |m_H|^2)}{g^2|v_1 v_2|}} ,$$
$$\tan\gamma_2 = \frac{\sin\theta}{\cos\theta + \frac{|\tilde{m}_2|}{|m_H|}\frac{g^2|v_1 v_2|}{2(m_{\omega 2}^2 - |\tilde{m}_2|^2)}} ,$$

where $\theta = \theta_H + \theta_g + \theta_1 + \theta_2$. The unitary matrices $C_{R,L}$, which diagonalize M^- in Eq. (1) to give real positive eigenvalues, are now given by $U_{R,L}$ and $P_{R,L}$ as

$$C_{R,(L)} = P_{R,(L)} U_{R,(L)} . \tag{A4}$$

If we neglect in Eq. (10) the tiny contribution from the origin (iii), which is doubly proportional to the fermion masses, the imaginary part of F_{fi}^k is proportional to $\mathrm{Im}(C_{L1i}C_{R2i}^* e^{-i\theta_2})$ for $T_{3f} = \frac{1}{2}$ and $\mathrm{Im}(C_{R1i}^* C_{L2i} e^{-i\theta_1})$ for $T_{3f} = -1/2$, which can be explicitly given as

$$\mathrm{Im}(C_{L11}C_{R21}^* e^{-i\theta_2}) = -\frac{g|v_1\tilde{m}_2 m_H|}{\sqrt{2}(m_{\omega 2}^2 - m_{\omega 1}^2)m_{\omega 1}}\sin\theta ,$$
$$\mathrm{Im}(C_{L12}C_{R22}^* e^{-i\theta_2}) = \frac{g|v_1\tilde{m}_2 m_H|}{\sqrt{2}(m_{\omega 2}^2 - m_{\omega 1}^2)m_{\omega 2}}\sin\theta ,$$
$$\mathrm{Im}(C_{L21}C_{R11}^* e^{-i\theta_1}) = -\frac{g|v_2\tilde{m}_2 m_H|}{\sqrt{2}(m_{\omega 2}^2 - m_{\omega 1}^2)m_{\omega 1}}\sin\theta , \tag{A5}$$
$$\mathrm{Im}(C_{L22}C_{R12}^* e^{-i\theta_1}) = \frac{g|v_2\tilde{m}_2 m_H|}{\sqrt{2}(m_{\omega 2}^2 - m_{\omega 1}^2)m_{\omega 2}}\sin\theta .$$

From these results Eq. (19) has been obtained. It should be noted that $\mathrm{Im}(F_{\tilde{f}i}^k)$ is proportional to the sine of the only one linear combination of the phases, $\sin\theta$, despite keeping all the phases of the mass parameters nonvanishing in the above derivation. This justifies our phase convention adopted in the text.

The scalar mass-squared matrix $\tilde{M}_f^2$, which is Hermitian, can be readily diagonalized by a unitary matrix S_f as

$$S_f^\dagger \tilde{M}_f^2 S_f = \begin{pmatrix} M_{\tilde{f}1}^2 & 0 \\ 0 & M_{\tilde{f}2}^2 \end{pmatrix},$$
$$S_f = \begin{pmatrix} \cos\theta_f & -e^{-i\beta_f}\sin\theta_f \\ e^{i\beta_f}\sin\theta_f & \cos\theta_f \end{pmatrix}. \quad \text{(A6)}$$

The diagonalized squark masses are

$$M_{\tilde{f}1(\tilde{f}2)}^2 = \frac{A-(+)\sqrt{B^2+C^2}}{2},$$
$$A = 2m_f^2 + \cos2\beta T_{3f} M_Z^2 + M_{\tilde{f}L}^2 + M_{\tilde{f}R}^2,$$
$$B = \cos2\beta(T_{3f} - 2Q_f \sin^2\theta_W) M_Z^2 + M_{\tilde{f}L}^2 - M_{\tilde{f}R}^2, \quad \text{(A7)}$$
$$C = 2m_f |R_f m_H + A_f^* m_{3/2}|.$$

The angles θ_f and β_f are calculated to be

$$\tan2\theta_f = \frac{C}{B},$$

$$\tan\beta_f = \begin{cases} -\dfrac{|m_H|\cot\beta\sin(\theta_H+\theta_1+\theta_2) - |A_f|m_{3/2}\sin\theta_{Af}}{|m_H|\cot\beta\cos(\theta_H+\theta_1+\theta_2) + |A_f|m_{3/2}\cos\theta_{Af}} & (T_{3f}=\frac{1}{2}), \\ -\dfrac{|m_H|\tan\beta\sin(\theta_H+\theta_1+\theta_2) - |A_f|m_{3/2}\sin\theta_{Af}}{|m_H|\tan\beta\cos(\theta_H+\theta_1+\theta_2) + |A_f|m_{3/2}\cos\theta_{Af}} & (T_{3f}=-\frac{1}{2}), \end{cases} \quad \text{(A8)}$$

where $\theta_{Af} = \arg(A_f)$.

If $\tilde{m}_3$ has the same complex phase as $\tilde{m}_2$, the imaginary part of H_q^k in Eq. (15) becomes

$$\mathrm{Im}(H_q^1) = -\mathrm{Im}(H_q^2) = -\left(\sin\alpha_q |A_q| - \sin\theta |R_q| \frac{|m_H|}{m_{3/2}}\right) \frac{m_{3/2} m_q}{M_{\tilde{q}2}^2 - M_{\tilde{q}1}^2}, \quad \text{(A9)}$$

where $\alpha_q = \theta_{Aq} - \theta_g$. We can see that the EDM due to the gluino contribution depends on two linear combinations of the phases α_q and θ.

APPENDIX B

Expressing the imaginary part of H_q^k in terms of the parameters, Eq. (14) is written as

$$d_q^G/e = \frac{2\alpha_s}{3\pi}\left(\sin\alpha_q|A_q| - \sin\theta R_q \frac{|m_H|}{m_{3/2}}\right) \frac{m_{3/2}m_q}{M_{\tilde{q}2}^2 - M_{\tilde{q}1}^2} Q_{\tilde{q}} \times \frac{1}{m_{\tilde{g}}} \sum_{k=1}^{2} (-1)^k I\left(\frac{M_{\tilde{q}k}^2}{m_{\tilde{g}}^2}\right). \quad \text{(B1)}$$

Here the identity $rI(r) = I(1/r)$ has been employed. Since $M_{\tilde{q}1}^2 \simeq M_{\tilde{q}2}^2$,

$$\frac{1}{M_{\tilde{q}2}^2 - M_{\tilde{q}1}^2} \sum_{k=1}^{2} (-1)^k I\left(\frac{M_{\tilde{q}k}^2}{m_{\tilde{g}}^2}\right) = \frac{1}{m_{\tilde{g}}^2} I'\left(\frac{M_{\tilde{q}1}^2}{m_{\tilde{g}}^2}\right) = \frac{m_{\tilde{g}}^2}{M_{\tilde{q}1}^4} K\left(\frac{m_{\tilde{g}}^2}{M_{\tilde{q}1}^2}\right), \quad \text{(B2)}$$

$K(r)$ being defined in Eq. (21). Substituting $M_{\tilde{q}}^2$ for $M_{\tilde{q}1}^2$, we then obtain Eq. (21).

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