Electric Dipole Moments in PseudoDirac Gauginos

Junji Hisano, Minoru Nagai, Tatsuya Naganawa and Masato Senami

Institute for Cosmic Ray Research,
University of Tokyo, Kashiwa 277-8582, Japan

Abstract

The SUSY CP problem is one of serious problems in construction of realistic supersymmetric standard models. We consider the problem in a framework in which adjoint chiral multiplets are introduced and gauginos have Dirac mass terms induced by a U(1) gauge interaction in the hidden sector. This is realized in hidden sector models without singlet chiral multiplets, which are favored from a recent study of the Polonyi problem. We find that the dominant contributions to electron and neutron electric dipole moments (EDMs) in the model come from phases in the supersymmetric adjoint mass terms. When the supersymmetric adjoint masses are suppressed by a factor of $\sim 100$ compared with the Dirac ones, the electron and neutron EDMs are suppressed below the experimental bound even if the SUSY particle masses are around 1 TeV. Thus, this model works as a framework to solve the SUSY CP problem.
1 Introduction

Origin of gaugino masses is one of the important issues in supersymmetric (SUSY) extension of the standard model (SM). In the minimal supergravity model [1], a hidden sector is introduced, and its interactions with the observable sector are assumed to be suppressed by the gravitational scale. Majorana gaugino mass terms in the minimal SUSY SM (MSSM) are derived by a singlet chiral multiplet with non-vanishing $F$-component in the hidden sector. However, introduction of the singlet chiral multiplet makes this model in trouble.

First, other $F$-term SUSY breaking terms in the scalar potential, the $A$ and $B$ terms, are also generated by the singlet chiral multiplet, and they have $O(1)$ phases relative to gaugino masses generically. They contribute to the electron and neutron electric dipole moments (EDMs) beyond the experimental bounds unless the SUSY particle masses are heavier than several TeV. This is called as the SUSY CP problem [2, 3].

Second is a cosmological problem called as the Polonyi problem [4]. This problem was considered not to be serious when the Polonyi field gets a mass of the dynamical SUSY breaking scale. However, it is recently pointed out that the Polonyi problem is more serious in hidden sector models in which a singlet chiral multiplet acquires $F$-term vacuum expectation value (VEV) to generate the gaugino masses [5]. The linear term of the singlet field is allowed in the Kähler potential. When the Hubble parameter during inflation is not small enough, the linear term destabilizes the minimum of the singlet scalar (Polonyi) field during the inflation, and the Polonyi problem is revived.

Therefore, it is one of the ways to construct a realistic SUSY SM to assume that the hidden sector has no singlet chiral field. Even if the hidden sector have no singlet chiral field, Majorana gaugino mass terms are generated by the pure supergravity (anomaly mediation) effect [6]. However, they are suppressed by one-loop factors compared with the gravitino mass. When the particle contents in the SUSY SM are minimal, the $B$ parameter in the Higgs potential is of the order of the gravitino mass. Thus, we need to extend this model furthermore.

An alternative way to generate gaugino mass terms is introduction of adjoint chiral superfields into the SUSY SM [7]. When the hidden sector has a $U(1)$ gauge multiplet whose $D$ component is non-vanishing, Dirac mass terms with gauginos and the adjoint fermions are generated [8, 9]. Their sizes can be comparable to the gravitino mass. Since the gaugino Majorana mass terms are also generated by the anomaly mediation effect, the gauginos are pseudoDirac fermions. We call this model as PseudoDirac Gaugino (PDG) model.

In this paper we discuss the electron and neutron EDMs in the PDG model. The Dirac gaugino mass term is $U(1)_R$ invariant, and it is pointed out that the EDMs are suppressed in a $U(1)_R$ symmetric limit [7], in which the Majorana gaugino mass terms, the supersymmetric Higgs mass term, and the $A$ terms are vanishing. The $U(1)_R$ breaking terms are introduced by the supergravity effects in the PDG model. We found that the EDMs are still suppressed when the supersymmetric adjoint multiplet mass is small. Thus, the PDG model is a framework to solve the SUSY CP problem.
This paper is organized as follows. In next section we review the PDG model in detail, and also discuss the SUSY particle mass spectrum. In Section 3 we discuss the CP violation and the EDMs derived from it. Section 4 is devoted to Conclusions and Discussion.

2 Model

As mentioned in Introduction, we assume that the hidden sector has no singlet chiral multiplet, and that the interactions with observable sector are suppressed by the gravitational scale \( M_\star \equiv 2.4 \times 10^{18} \text{ GeV} \). Furthermore, it is assumed for simplicity that the hidden sector does not have fields with VEV \( \sim M_\star \). In this case, the Majorana gaugino masses in the SUSY SM are derived by the anomaly mediation effect, and they are one-loop suppressed compared with the gravitino mass \( \frac{m_{3/2}}{2} \).

\[ \mu_{\tilde{g}_i} = \frac{\beta_i(g_i^2)}{2g_i^2} F_\phi. \]  

Here, \( \beta_i(g_i^2) \) \((i = Y, 2, 3)\) are the beta functions for \( U(1)_Y, \text{SU}(2)_L \) and \( \text{SU}(3)_C \) gauge coupling constants squared and \( F_\phi \) is proportional to \( m_{3/2} \).

On the other hand, the scalar masses in the SUSY SM are comparable to the gravitino mass unless the Kähler potential has a sequestering form between the hidden and observable sectors. We assume that the sfermion mass universality is imposed by some flavor symmetries or underlying physics.

In order to give sizable masses to the gauginos, we introduce adjoint chiral multiplets to each gauge multiplet in the SUSY SM. The adjoint fermions and the gauginos get Dirac mass terms by a VEV of the \( D \) component of a U(1) gauge field in the hidden sector. The kinetic terms of the adjoint and gauge multiplets and their interaction with the U(1) gauge field are given as follows,

\[ \mathcal{L} = \int d^2 \theta \left\{ \sum_{i=Y,2,3} \left( \text{Tr} \left[ \frac{1}{4k_i g_i^2} W_i^\alpha W_{i\alpha} \right] - \frac{\sqrt{2}c_i}{k_i g_i^2 M_\star} \text{Tr} \left[ W_i^\alpha W_\alpha^i \Sigma_i \right] \right) + h.c. \right\} \]

\[ + \int d^4 \theta \sum_{i=Y,2,3} \frac{1}{k_i g_i^2} \text{Tr} \left[ \Sigma_i^\dagger e^{-2V_i} \Sigma_i e^{2V_i} \right]. \]  

Here, \( W_i^\alpha \) and \( \Sigma_i \) are gauge field strengths and adjoint multiplets, respectively, and their interaction with the \( U(1) \) gauge field in the hidden sector, \( W_\alpha^i \), is suppressed by the gravitational scale \( M_\star \). When the \( U(1) \) gauge multiplet in the hidden sector get the \( D \)-component VEV \( \langle W_\alpha^i \rangle = D\theta_\alpha \), the Dirac masses for the gauginos and adjoint fermions are given as \( M_{\tilde{D}_i} = c_i D/M_\star \). We assume \( D/M_\star \sim m_{3/2} \) so that the gaugino masses are comparable to the scalar masses in the SUSY SM.

The \( U(1)_R \) charges for the adjoint multiplets are zero in this model, as expected from Eq. (2.2). Thus, their supersymmetric mass terms vanish in a \( U(1)_R \) symmetric limit,
similar to those of the Higgs multiplets in the SUSY SM, \( H_1 \) and \( H_2 \). The supergravity effect also may generate the mass terms when following terms are present in the Kähler potential \([10, 11]\),

\[
\mathcal{L} = -\int d^4\theta \frac{\phi_i^\dagger}{\phi} \left\{ c_H H_1 H_2 + \sum_{i=Y,2,3} \frac{c_{\Sigma_i}}{2k_i g_i^2} \text{Tr}[\Sigma_i^2] \right\} + h.c.. \tag{2.3}
\]

Here, \( \phi(=1+F\theta^2) \) is the compensator chiral multiplet in supergravity. From Eq. (2.3), the supersymmetric Higgs (\( \mu_H \)) and adjoint masses (\( \mu_{\Sigma_i} \)) and their \( B \) parameters (\( B_H \) and \( B_{\Sigma_i} \)) are given as

\[
\mu_H = c_H F_\phi^*, \quad B_H = -F_\phi, \\
\mu_{\Sigma_i} = c_{\Sigma_i} F_\phi^*, \quad B_{\Sigma_i} = -F_\phi. \tag{2.4}
\]

When the Higgs and adjoint chiral multiplets have \( M_\star \)-suppressed interactions with a non-singlet chiral multiplet \( Z \) in the hidden sector as

\[
\int d^4\theta \frac{Z^\dagger Z}{M_\star^2} \left\{ c_H' H_1 H_2 + \sum_{i=Y,2,3} \frac{c_{\Sigma_i}'}{2k_i g_i^2} \text{Tr}[\Sigma_i^2] \right\} + h.c., \tag{2.5}
\]

their \( B \) parameters and supersymmetric masses are changed from Eqs. (2.4). However, if the VEV for \( Z \) is not as large as \( \sim M_\star \), the supergravity contributions in Eq. (2.4) are still dominant in the supersymmetric masses.

The sizes of \( \mu_{\Sigma_i} \) and \( \mu_H \) depend on underlying physics. Even if \( c_{\Sigma_i} \) and \( c_H \) are zero at tree level, they may be induced radiatively. For example, when the massive chiral multiplets \( X \) and \( \bar{X} \) have an interaction with \( \Sigma_i \), \( fX\bar{X}\Sigma_i \), \( c_{\Sigma_i} \) is generated at one-loop level as \( c_{\Sigma_i} \sim f^2/(4\pi)^2 \), and it is not suppressed by the mass of \( X \) and \( \bar{X} \). Thus, \( c_{\Sigma_i} \) and \( c_H \) are expected to be larger than \( O(10^{-12.3}) \). In the following we assume that \( M_D, \mu_{\Sigma_i} \) so that the gauginos are pseudoDirac.

The successful gauge coupling unification in the MSSM is spoiled in the PDG model, since the adjoint chiral multiplets are introduced in the SUSY SM. In order to recover the unification, it is pointed out in Ref. \([8, 9]\) that the adjoint chiral multiplets are embedded in complete multiplets under the GUT gauge group and the additional fields in the multiplets, which are called as bachelor fields, are introduced. The candidates for the GUT gauge group are \( SU(5) \) and \( SU(3)^3 \). The low-energy gauge coupling constants favor the bachelor masses (\( M_\mu \)) are \( 10^{(5-9)} \) GeV and the GUT scale is \( 10^{(17-18)} \) GeV. It is economical if the GUT scale is close to the gravitational scale.

The low-energy mass spectrum in the PDG model is affected by the radiative correction, and it depends on detail of the models between the weak scale and the gravitational scale. We introduce the bachelor fields as a working hypothesis and take the GUT scale to be \( M_\star \).
in the following. The contribution of the adjoint and bachelor fields to the beta functions for gauge coupling constants \((b_B)\) is \(5(3)\) in the SU(5) model (the SU(3)\(^3\) model). The Dirac gaugino masses at low energy are given from the one-loop renormalization group (RG) equations as

\[
\frac{M_{D_Y}}{g_Y}(m_{\text{SUSY}}) = \frac{M_{D_Y}(M_\ast)}{g_Y},
\]

\[
\frac{M_{D_2}}{g_2}(m_{\text{SUSY}}) = \left( \frac{\alpha_2(m_{\text{SUSY}})}{\alpha_2(M_B)} \right)^{-\frac{2}{3}} \left( \frac{\alpha_2(M_B)}{\alpha_2(M_\ast)} \right)^{-\frac{2}{1+b_B}} \frac{M_{D_2}(M_\ast)}{g_2},
\]

\[
\frac{M_{D_3}}{g_3}(m_{\text{SUSY}}) = \left( \frac{M_B}{m_{\text{SUSY}}} \right)^{\frac{3\alpha_3(m_{\text{SUSY}})}{2\pi}} \left( \frac{\alpha_3(M_B)}{\alpha_3(M_\ast)} \right)^{-\frac{3}{3+b_B}} \frac{M_{D_3}(M_\ast)}{g_3}.
\]

When \(b_B = 3\), the SU(3) gaugino mass becomes

\[
\frac{M_{D_3}}{g_3}(m_{\text{SUSY}}) = \left( \frac{M_\ast}{m_{\text{SUSY}}} \right)^{\frac{3\alpha_3(m_{\text{SUSY}})}{2\pi}} \frac{M_{D_3}(M_\ast)}{g_3}.
\]

The gaugino mass differences among the gauge groups are larger than those in the MSSM under the minimal supergravity assumption (cMSSM), in which the gaugino mass ratios are given by those of squares of the gauge coupling constants. If \(\sqrt{\frac{3}{5}}M_{D_Y}/g_Y = M_{D_2}/g_2 = M_{D_3}/g_3\) at \(M_\ast\), the gaugino mass ratios in the SU(3)\(^3\) model are \(M_{D_2}/M_{D_Y} \simeq (2.3 - 2.6)\) and \(M_{D_3}/M_{D_Y} \simeq 11\), depending on \(M_B\). In the SU(5) model, the ratios are larger than those in the SU(3)\(^3\) model, \(M_{D_2}/M_{D_Y} \simeq (2.3 - 4.3)\) and \(M_{D_3}/M_{D_Y} \simeq (11 - 31)\). This is because the gauge coupling constants at \(M_\ast\) are quite large when \(M_B \sim 10^{(7-8)}\) GeV.

The adjoint supersymmetric masses \(\mu_{\Sigma_i}\) also receive sizable radiative corrections. The following combination including \(\mu_{\Sigma_i}\) is RG invariant at one-loop level,

\[
\frac{\mu_{\Sigma_i} B_{\Sigma_i}}{M_{D_i}^2},
\]

for \(i = Y, 2, 3\).

The light Higgs boson is predicted to be lighter in the PDG model, relatively to the MSSM. It is, unfortunately, contrary to null results of the Higgs boson searches. The first reason comes from the quartic coupling of the Higgs boson given from Eq. (2.2) as

\[
\frac{1}{8} \left( \frac{m_{A_2}^2}{m_{A_2}^2 + 4M_{D_2}^2} g_Y^2 + \frac{m_{A_Y}^2}{m_{A_Y}^2 + 4M_{D_Y}^2} g_Y^2 \right).
\]

Here, we take simply \(\mu_{\Sigma_i} B_{\Sigma_i} \simeq 0\) and \(\mu_{\Sigma_i} \simeq 0\), and \(m_{A_i}(i = Y, 2)\) are the SUSY breaking masses of the adjoint scalars for \(U(1)_Y\) and \(SU(2)_L\) [8]. The MSSM quartic coupling is recovered in a case of \(m_{A_i} \gg M_{D_i}\). The quartic coupling is suppressed when \(m_{A_i} \sim M_{D_i}\), and the Higgs boson mass becomes lighter. The second reason is that while the large \(A\) parameter for top squarks enhances the radiative correction to the light Higgs boson mass
it is one-loop suppressed in the PDG model. Since the hidden sector have no singlet fields, the $A$ parameters come from only the anomaly mediation similar to the Majorana gaugino masses. The $A$ parameters for the supersymmetric Yukawa coupling $\lambda_{flm}$ are given as

$$A_{flm} = -(\gamma_f + \gamma_l + \gamma_m)F_\phi$$

where $\gamma_f$ is the anomalous dimension for a field $f$. These implies from null results of the light Higgs boson searches that the large top squark masses ($\gtrsim 1$ TeV) are required so that the radiative correction to the Higgs boson mass is enhanced.

In the PDG model the radiative corrections to the scalar masses due to non-vanishing Dirac gaugino masses are finite and are not enhanced by a logarithm. This nature is called as “supersoft” [8]. The first and second-generation sfermions do not receive large radiative corrections to the mass terms when the gaugino masses are comparable to the sfermion masses and the $U(1)_Y$ D-term contribution to the RG equations is zero. The third generation sfermions and Higgs boson masses get RG effects by the large Yukawa interaction.

### 3 Electron and neutron EDMs

First, we review electron and neutron EDMs in the MSSM briefly, and show the constraints on the CP phases in the parameters in the model. In the MSSM the new sources of CP violation in the flavor-conserving terms are phases in the gaugino masses, $A$ and $B$ parameters, and the supersymmetric Higgs mass. The physical CP phases are following two due to degrees of freedom for rephasing, $U(1)_R$ and $U(1)_{PQ}$,

$$\phi_{\mu g} \equiv \arg(\mu_H \mu_{\tilde{g}} (\mu_H B_H)^*), \quad \phi_{A_f} \equiv \arg(\mu_{\tilde{g}} A_f^*) \quad (3.11)$$

Here the gaugino mass unification in GUTs and universality of $A$ parameters are assumed.

The relevant CP-violating operators for the electron and neutron EDMs are following up to dimension six [13],

$$\mathcal{L}_{(C)EDM} = - \sum_{f=e,u,d} i f \frac{d_f}{2} \bar{f}(F \cdot \sigma)\gamma_5 f + \frac{\alpha_s}{8\pi} G\tilde{G} - \sum_{q=u,d} i \frac{d_q}{2} \bar{q}(g_s G \cdot \sigma)\gamma_5 q + \frac{1}{3} \omega G\tilde{G}G \quad (3.12)$$

where $F_{\mu\nu}$ and $G_{\mu\nu}$ are the electromagnetic and the SU(3)$_C$ gauge field strengths, respectively. The terms in the second line are the flavor-conserving CP-violating terms in QCD, and the first term is the QCD theta term. Here, we simply assume Peccei-Quinn symmetry for it. The second and third terms are quark chromoelectric dipole moments (CEDMs) and the Weinberg operator, respectively.
The current experimental bound on the electron EDM is $|d_e| < 1.6 \times 10^{-27}$ ecm from the $^{205}\text{Tl}$ EDM measurement [14]. The neutron EDM bound is recently improved as $|d_n| < 2.9 \times 10^{-26}$ ecm [15]. The evaluation of the neutron EDM from the parton level is still a difficult task. Here, we use a formula derived from the QCD sum rule [16], in which the valence quark contribution to the EDM is evaluated\(^\#1\),

$$
\frac{d_n}{e} = (1 \pm 0.5) \frac{|\langle \bar{q}q \rangle|}{(225 \text{ MeV})^3} \left( 1.1 (d_d^c + 0.5d_u^c) + 1.4 \left( \frac{d_d}{e} - 0.25 \frac{d_u}{e} \right) \right).
$$

(3.13)

The Weinberg operator contribution to the neutron EDM is also evaluated as [19]

$$
\frac{d_n}{e} \simeq \omega \times 22 \text{ MeV},
$$

(3.14)

though it has more uncertainties.

In the MSSM the electron EDM is given as

$$
\frac{d_e}{e} = \sum_{i=Y,2} \text{Im} [\mu_{\tilde{g}_i} H] d_i^{(e)}(|\mu_{\tilde{g}_i}|^2) + \text{Im} [\mu_{\tilde{g}_Y} A_i^*] d_i^{(e)}(|\mu_{\tilde{g}_Y}|^2),
$$

(3.15)

while the quark EDMs and CEDMs are given as

$$
\frac{d_q}{e} = \sum_{i=Y,2,3} \text{Im} [\mu_{\tilde{g}_i} H] d_i^{(q)}(|\mu_{\tilde{g}_i}|^2) + \sum_{i=Y,3} \text{Im} [\mu_{\tilde{g}_Y} A_i^*] d_i^{(q)}(|\mu_{\tilde{g}_Y}|^2),
$$

$$
\frac{d_c}{e} = \sum_{i=Y,2,3} \text{Im} [\mu_{\tilde{g}_i} H] \tilde{d}_i^{(q)}(|\mu_{\tilde{g}_i}|^2) + \sum_{i=Y,3} \text{Im} [\mu_{\tilde{g}_Y} A_i^*] \tilde{d}_i^{(q)}(|\mu_{\tilde{g}_Y}|^2),
$$

(3.16)

where $q = u, d$ and we take a basis in which $\mu_H B_H$ is real. The mass functions for EDMs, $d_i^{(f)}(x)$ and $d_i^{(f)\prime}(x)$ ($f = e, u, d$), and those for CEDMs, $\tilde{d}_i^{(f)}(x)$ and $\tilde{d}_i^{(f)\prime}(x)$ ($f = u, d$), are defined for convenience to evaluate the EDMs and CEDMs in the PDG model later. The explicit forms in a limit $m_{\text{SUSY}} \gg m_Z$ are given in Appendix. More complete formulae can be derived from those in Ref. [3]. The Weinberg operator contribution to the neutron EDM is subdominant in the MSSM except for cases in some specific mass spectrum, and then it is ignored here.

### Footnote

\(^\#1\) The sea quarks, including strange quark, may contribute to the neutron EDM sizable [17]. In this paper we ignore them for simplicity since the evaluation still has uncertainties. The mercury EDM [18] is also sensitive to the quark CEDMs, and the bound is as stringent as the neutron one though it is also expected to suffer from larger hadronic uncertainties than neutron one due to the nuclear dynamics [13].
and

\[
\begin{align*}
\frac{d_d}{e^-} &= -\left(\frac{2\alpha_3}{27} + \frac{7\alpha_2}{24} - \frac{11\alpha_Y}{648}\right) \frac{m_d}{4\pi m^2_{\text{SUSY}}} \tan \beta \sin \phi_{\mu_3} - \left(\frac{2\alpha_3}{27} - \frac{\alpha_Y}{324}\right) \frac{m_d}{4\pi m^2_{\text{SUSY}}} \sin \phi_{A_d}, \\
\frac{d_u^c}{e^-} &= -\left(\frac{5\alpha_3}{18} + \frac{\alpha_2}{8} + \frac{11\alpha_Y}{216}\right) \frac{m_u}{4\pi m^2_{\text{SUSY}}} \tan \beta \sin \phi_{\mu_3} - \left(\frac{5\alpha_3}{18} + \frac{\alpha_Y}{108}\right) \frac{m_u}{4\pi m^2_{\text{SUSY}}} \sin \phi_{A_u}, \\
\frac{d_u}{e^-} &= +\left(\frac{4\alpha_3}{27} + \frac{\alpha_2}{4} - \frac{5\alpha_Y}{324}\right) \frac{m_u}{4\pi m^2_{\text{SUSY}}} \tan^{-1} \beta \sin \phi_{\mu_3} + \left(\frac{4\alpha_3}{27} + \frac{\alpha_Y}{81}\right) \frac{m_u}{4\pi m^2_{\text{SUSY}}} \sin \phi_{A_u}, \\
\frac{d_u^c}{e^-} &= -\left(\frac{5\alpha_3}{18} + \frac{\alpha_2}{8} + \frac{5\alpha_Y}{216}\right) \frac{m_u}{4\pi m^2_{\text{SUSY}}} \tan^{-1} \beta \sin \phi_{\mu_3} - \left(\frac{5\alpha_3}{18} - \frac{\alpha_Y}{54}\right) \frac{m_u}{4\pi m^2_{\text{SUSY}}} \sin \phi_{A_u}.
\end{align*}
\]

From these equations, the phases are constrained as

\[
| \sin \phi_{\mu_3} | < \left( \frac{m_{\text{SUSY}}}{6.1 \text{ TeV}} \right)^2 \left( \frac{\tan \beta}{10} \right)^{-1}, \quad | \sin \phi_{A_f} | < \left( \frac{m_{\text{SUSY}}}{0.65 \text{ TeV}} \right)^2, \quad (3.19)
\]

from the experimental bound on the electron EDM, and

\[
| \sin \phi_{\mu_3} | < \left( \frac{m_{\text{SUSY}}}{12 \text{ TeV}} \right)^2 \left( \frac{\tan \beta}{10} \right)^{-1}, \quad | \sin \phi_{A_f} | < \left( \frac{m_{\text{SUSY}}}{3.7 \text{ TeV}} \right)^2, \quad (3.20)
\]

from neutron one. Here, we use \(\alpha_3(m_Z) = 0.1176, \alpha^{-1}(m_Z) = 127.918, \sin^2 \theta_W(m_Z) = 0.23122\) \[^{20}\], \(m_d(m_Z) = 4.1 \text{ MeV}\) and \(m_u(m_Z) = 2.4 \text{ MeV}\) \[^{22}\] and the QCD corrections to the dipole operators are ignored. The center value in Eq. (3.13) is taken in evaluation of Eq. (3.20). The down quark contribution tends to dominate in the neutron EDM. Especially, it is significant when \(\sin \phi_{\mu_3} \neq 0\). It is found from these bounds that the SUSY particles should be larger than several TeV when the CP phase is \(O(1)\).

In the PDG model the complex parameters introduced are

\[
M_{D_i}, \ \mu_{\tilde{3}}, \ \mu_{\Sigma_i}, \ B_{\Sigma_i} \ (i = Y, 2, 3), \\
\mu_H, \ B_H, \ A_f \ (f = u, d, l).
\]

Degrees of freedom for rephasing come from phase transformations of adjoint fields in addition to \(U(1)_R\) and \(U(1)_{PQ}\) in the MSSM. Thus, the physical CP phases are four, including \(\phi_{\mu_3}\) and \(\phi_{A_f}\) in Eq. (3.11), when assuming GUT and universality of the \(A\) terms. The new phases in the PDG model are

\[
\phi_{\mu_3} \equiv \arg(\mu_H \mu_{\Sigma_i}^* (\mu_H B_H)^* M^2_{D_i}), \quad \phi_{B_{\Sigma}} \equiv \arg(\mu_{\Sigma_i} B_{\Sigma_i} M^2_{D_i}). \quad (3.22)
\]

In the following we take a basis in which \(M_{D_i}\) and \(\mu_H B_H\) are real for convenience.

\[^{22}\] These values are evaluated from \(m_d(1 \text{ GeV}) = 8.9 \text{ MeV}\) and \(m_u(1 \text{ GeV}) = 5.1 \text{ MeV} \[^{21}\].
When the hidden sector does not have singlet chiral multiplets, the $A$ parameters and
the Majorana gaugino masses are dominated by the anomaly mediation contribution as in
Eqs. (2.10) and (2.11), and their phases are automatically aligned. That is, $\phi_{A_f} = 0$. On
the other hand, the $B$ parameters still have other contributions even in the case, and $\phi_{B_S}$
and $\phi_{\mu_S}$ do not necessarily vanish. The adjoint scalar bosons are not directly coupled
to quarks and leptons. The CP phase $\phi_{B_S}$ contributes to the EDMs at three-loop level via
the Weinberg operator \[7\]. We evaluate the EDMs in $\phi_{\mu_q} \neq 0$ and $\phi_{\mu_S} \neq 0$, first. The
Weinberg operator contribution in $\phi_{B_S} \neq 0$ is discussed later.

The adjoint fermions are also coupled to quarks and leptons via the gauginos. Thus,
when $\mu_{\tilde{g}_i}$ and $\mu_{\Sigma_i} \ll M_{D_i}$, the one-loop contributions to the EDMs and CEDMs are derived
from those in the MSSM as

$$
\frac{d_e}{e} = \sum_{i=Y,2} \left\{ \text{Im} [\mu_{\tilde{g}_i},\mu_H] \left( 1 + M^2_D \frac{\partial}{\partial M^2_D} \right) d_i^{(e)}(M^2_D) + \text{Im} \left[ \mu_{\Sigma_i}^*,\mu_H \right] M^2_D \frac{\partial}{\partial M^2_D} d_i^{(e)}(M^2_D) \right\} 
+ \text{Im} [\mu_{\tilde{g}_i},\nu_{A_i}] \left( 1 + M^2_{D_Y} \frac{\partial}{\partial M^2_{D_Y}} \right) d_i^{(e)}(M^2_{D_Y}) - \text{Im} [\mu_{\Sigma_i},\nu_{A_i}] M^2_{D_Y} \frac{\partial}{\partial M^2_{D_Y}} d_i^{(e)}(M^2_{D_Y}),
$$

$$
\frac{d_q}{e} = \sum_{i=Y,2,3} \left\{ \text{Im} [\mu_{\tilde{g}_i},\mu_H] \left( 1 + M^2_D \frac{\partial}{\partial M^2_D} \right) d_i^{(q)}(M^2_D) + \text{Im} \left[ \mu_{\Sigma_i}^*,\mu_H \right] M^2_D \frac{\partial}{\partial M^2_D} d_i^{(q)}(M^2_D) \right\} 
+ \sum_{i=Y,3} \left\{ \text{Im} [\mu_{\tilde{g}_i},A_i^*] \left( 1 + M^2_D \frac{\partial}{\partial M^2_D} \right) d_i^{(q)}(M^2_D) - \text{Im} [\mu_{\Sigma_i},A_i] M^2_D \frac{\partial}{\partial M^2_D} d_i^{(q)}(M^2_D) \right\},
$$

$$
\frac{d_q}{e} = \sum_{i=Y,2,3} \left\{ \text{Im} [\mu_{\tilde{g}_i},\mu_H] \left( 1 + M^2_D \frac{\partial}{\partial M^2_D} \right) d_i^{(q)}(M^2_D) + \text{Im} \left[ \mu_{\Sigma_i}^*,\mu_H \right] M^2_D \frac{\partial}{\partial M^2_D} d_i^{(q)}(M^2_D) \right\} 
+ \sum_{i=Y,3} \left\{ \text{Im} [\mu_{\tilde{g}_i},A_i^*] \left( 1 + M^2_D \frac{\partial}{\partial M^2_D} \right) d_i^{(q)}(M^2_D) - \text{Im} [\mu_{\Sigma_i},A_i] M^2_D \frac{\partial}{\partial M^2_D} d_i^{(q)}(M^2_D) \right\},
$$

(3.23)

$(q = u, d)$. We include the contribution in $\phi_{A_f} \neq 0$ for completeness though it is subdom-
inant.

The one-loop contributions to the EDMs and CEDMs are suppressed by $\mu_{\tilde{g}_i}$ or $\mu_{\Sigma_i}$. When taking a common value for SUSY breaking parameters except for $\mu_{\tilde{g}_i}$ and $\mu_{\Sigma_i}$, the
We evaluate the EDMs in a case that assume the SU(5) model for the bachelor fields and the sfermion mass universality. Here, $M^0_{\text{Dirac gaugino mass}} \sim \text{experimental bounds.}$

In Fig. 1 we show the electron and neutron EDMs in the PDG model as functions of Dirac gaugino mass ($M_D$) and sfermion mass ($m_0$) at the gravitational scale $M_\ast$. We assume the SU(5) model for the bachelor fields and the sfermion mass universality. Here, $\sin \phi_{\mu_\Sigma} = 1$, $\sin \phi_{\mu_\tilde{g}} = 0$ and $|\mu_{\Sigma}/M_D| = 10^{-3}$ at $M_\ast$, $M_B = 10^9$ GeV, $\mu_H = 200$ GeV, and $\tan \beta = 10$.

The electron EDM and down quark EDM and CEDM are given as follows,

$$d_e = e = -\frac{\alpha_2 \delta_{\tilde{g}_i}}{30} \frac{m_\mu \tan \beta \sin \phi_{\mu_\tilde{g}}}{4 \pi m^2_{\text{SUSY}}} \sin \phi_{\mu_\Sigma},$$

$$d_d = e = -\left(\frac{2 \alpha_2 \delta_{\tilde{g}_i}}{135} + \frac{\alpha_2 \delta_{\tilde{g}_i}}{15} - \frac{\alpha_Y \delta_{\tilde{g}_i}}{162}\right) \frac{m_\mu \tan \beta}{4 \pi m^2_{\text{SUSY}}} \sin \phi_{\mu_\tilde{g}}$$

$$+ \left(\frac{8 \alpha_2 \delta_{\Sigma_i}}{135} + \frac{9 \alpha_2 \delta_{\Sigma_i}}{40} - \frac{7 \alpha_Y \delta_{\Sigma_i}}{648}\right) \frac{m_\mu \tan \beta}{4 \pi m^2_{\text{SUSY}}} \sin \phi_{\mu_\tilde{g}},$$

$$d_d = e = \left(\frac{2 \alpha_2 \delta_{\tilde{g}_i}}{45} - \frac{\alpha_2 \delta_{\tilde{g}_i}}{20} - \frac{\alpha_Y \delta_{\tilde{g}_i}}{54}\right) \frac{m_\mu \tan \beta}{4 \pi m^2_{\text{SUSY}}} \sin \phi_{\mu_\tilde{g}}$$

$$+ \left(\frac{29 \alpha_2 \delta_{\tilde{g}_i}}{90} + \frac{30 \alpha_2 \delta_{\Sigma_i}}{40} + \frac{7 \alpha_Y \delta_{\Sigma_i}}{216}\right) \frac{m_\mu \tan \beta}{4 \pi m^2_{\text{SUSY}}} \sin \phi_{\mu_\Sigma},$$

where $\delta_{\tilde{g}_i} = |\mu_{\tilde{g}_i}/m_{\text{SUSY}}|$ and $\delta_{\Sigma_i} = |\mu_{\Sigma_i}/m_{\text{SUSY}}|$. We keep terms proportional to $\tan \beta$ in the equations. Compared with Eqs. (3.17) and (3.18), it is found that $\delta_{\tilde{g}_i}$ and $\delta_{\Sigma_i}$ should be smaller than $\sim 0.01$ so that the electron and neutron EDMs are suppressed below the experimental bounds.

In Fig. 1 we show the electron and neutron EDMs in the PDG model as functions of the Dirac gaugino mass ($M_D$) and the sfermion mass ($m_0$) at the gravitational scale $M_\ast$. We assume the SU(5) model for the bachelor fields and the sfermion mass universality. Here, we evaluate the EDMs in a case that $|\mu_{\Sigma_i}/M_D| = 10^{-3}$, $\sin \phi_{\mu_\Sigma} = 1$, and $\sin \phi_{\mu_\tilde{g}} = 0$ at $M_\ast$. In this case, $|\mu_{\Sigma_3}/M_D| = 0.5 \times 10^{-2}$ and $|\mu_{\Sigma_3}/M_D| = 5.7 \times 10^{-2}$ at $m_{\text{SUSY}}$ due to the radiative corrections, and the phases in the adjoint fermion masses, $\phi_{\mu_\Sigma}$, are the dominant sources of CP violation in the EDMs. Other parameters are taken as $M_B = 10^9$ GeV, $\mu_H = 200$ GeV, and $\tan \beta = 10$. The EDMs are proportional to $\mu_{\Sigma_i}$ in $\mu_{\Sigma_i}/M_D \ll 1$. 

Figure 1: Electron and neutron EDMs in the PDG model as functions of Dirac gaugino mass ($M_D$) and sfermion mass ($m_0$) at the gravitational scale $M_\ast$. We assume the SU(5) model for the bachelor fields and the sfermion mass universality. Here, $\sin \phi_{\mu_\Sigma} = 1$, $\sin \phi_{\mu_\tilde{g}} = 0$ and $|\mu_{\Sigma}/M_D| = 10^{-3}$ at $M_\ast$. 

$M_B = 10^9$ GeV, $\mu_H = 200$ GeV, and $\tan \beta = 10$. The EDMs are proportional to $\mu_{\Sigma_i}$ in $\mu_{\Sigma_i}/M_D \ll 1$. 

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Figure 2: Electron and neutron EDMs in the PDG model as functions of Dirac gaugino mass and sfermion mass at the gravitational scale. Here, $F_\phi/M_D = 1$, $\mu_{\Sigma_i} = 0$ and $\sin \phi_{\tilde{g}} = 1$. The other assumptions and input parameters are the same as in Fig. 1.

It is found from the figure that if $|\mu_{\Sigma_i}/M_D| \lesssim 0.01$, the EDMs are suppressed below the experimental bounds even when $m_0$ and $M_D$ are smaller than 1 TeV.

In Fig. 2 the electron and neutron EDMs are shown in a case of $\mu_{\Sigma_i} = 0$, $|F_\phi/M_D| = 1$ and $\sin \phi_{\tilde{g}} = 1$. Other parameters are the same as in Fig. 1. In this case, the phases in the Majorana gaugino mass terms are the dominant sources of CP violation in the EDMs. Even if $\mu_{\Sigma_i}$ is suppressed, the Majorana gaugino masses are derived by the anomaly mediation effect. Thus, the evaluated EDMs are the lower bounds on the EDMs in the PDG model unless accidental cancellation suppress the EDMs.

The EDMs in the case of $\mu_{\Sigma_i} = 0$ and $\sin \phi_{\tilde{g}} = 1$ are much suppressed below the current experimental bounds. Especially, the electron EDM is suppressed by the cancellation when $m_0 \lesssim 2$ TeV and $M_D \lesssim 1.5$ TeV. However, the finite values are still predicted so that the future experiments may cover them.

The neutron EDM measurements will be improved at SNS and ILL in near future, and it is argued that their sensitivities may reach to $\sim 10^{-29}$ e cm [22]. The deuteron EDM is also planned to be measured in the storage ring with sensitivity $10^{-29}$ e cm, which corresponds to $d_n \sim 10^{-29}$ e cm [22]. The electron EDM measurements using polarizable paramagnetic molecules are aiming to the sensitivities of $10^{-29}$ e cm [23, 22].

Finally, we discuss the Weinberg operator induced by non-zero $\phi_{B_{\Sigma_i}}$. It is derived at two-loop diagrams, and roughly evaluated by dimensional counting as

$$\omega \sim \frac{\alpha_3^2}{(4\pi)^2} \frac{\text{Im}[\mu_{\Sigma_i} B_{\Sigma_i}]}{m_{\text{SUSY}}^4}.$$  \hfill (3.25)

#3 Notice that in the PDG model the one-loop beta function for the SU(3)$_C$ gauge coupling constant is zero. Thus, non-vanishing $\mu_{\tilde{g}}$ is derived at two-loop level, and the gluino contribution to the quark (C)EDM in $\sin \phi_{\mu_S} = 0$ is subdominant.
Thus, the neutron EDM is derived from it as

\[ d_n \sim 4 \times 10^{-26} \text{ e cm} \times \left( \frac{\text{Im}[\mu_{3}\Sigma B_{3}]}{m_{SUSY}^{2}} \right) \left( \frac{m_{SUSY}}{1 \text{ TeV}} \right)^{-2}, \]  

(3.26)

though it also suffers from large hadronic uncertainties. The Weinberg operator contribution is marginal to the current bound on the neutron EDM if \( \mu_{3}\Sigma B_{3} \) is not suppressed. When the supersymmetric adjoint mass terms and the SUSY breaking term associated with them come from independent terms, such as in Eqs. (2.3) and (2.5), \( \mu_{3}\Sigma B_{3} \) may be comparable to \( m_{3/2}^{2} \) while \( \mu_{3} \ll m_{3/2} \). In such models, the Weinberg operator contribution would be dominant in the neutron EDM.

4 Conclusions and Discussion

The SUSY CP problem is one of the serious problems in construction of realistic SUSY SMs. We consider the problem in a framework in which the gaugino mass terms are Dirac-type ones induced by a U(1) gauge interaction in the hidden sector. The adjoint chiral multiplets are introduced to have the Dirac mass terms. This model is called as the PDG model in this paper. The setup is also preferable from a viewpoint of the Polonyi problem, since we do not need to introduce singlet chiral superfields in the hidden sector in order to derive the gaugino masses as in the minimal supergravity model.

We find that the dominant contributions to the electron and neutron EDMs in the PDG model come from the phases in the supersymmetric adjoint mass terms. When the supersymmetric adjoint masses are suppressed by a factor of \( \sim 100 \) compared with the Dirac ones, the neutron and electron EDMs are suppressed below the experimental bounds even if the SUSY particle masses are around 1 TeV. Thus, this model works as a framework to solve the SUSY CP problem.

The electron and neutron EDMs are suppressed in the PDG model, however, the predicted values may be covered in the near future experiments. Even if the supersymmetric adjoint masses are much suppressed, the Majorana gaugino mass terms, which are induced by the anomaly mediation effect, may induce large enough neutron EDM to be observed in near future experiments. In addition to it, the phase in the \( B \) term of the adjoint scalar boson contributes to the Weinberg operator. It is not necessarily suppressed even if the supersymmetric adjoint mass is. In the case, the Weinberg operator would be dominant in the neutron EDM.

In this paper, we consider only the CP-violating terms which are flavor-conserving. The flavor-violating ones also contribute to the EDMs in the SUSY SM. Especially, when both the left- and right-handed sfermions have the flavor-mixing mass terms, the phases in them contribute to the EDMs at one-loop level. Even if the sfermion mixing angles are suppressed similar to the CKM ones, the EDMs can reach to the experimental bounds in the MSSM [25]. These contributions are also suppressed in the PDG model, since the contributions are proportional to the Majorana bino or gluino masses.
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A Appendix

Here, we present mass functions for the electron and quark EDMs and quark CEDMs in the MSSM. Here, we assume $|\tilde{\mu}_{\tilde{g}_2/\gamma} \pm \mu_H| \gg m_Z$. Furthermore, we take a common value $m$ for the sfermion masses.

The mass functions for the electron EDM in Eq. (3.15) are given as follows,

$$d_{2e}(|\mu_{\tilde{g}_2}|^2) = \frac{\alpha_2}{4\pi} m_e \tan \beta \times \hat{D}_{M^2}(|\mu_{\tilde{g}_2}|^2,|\mu_H|^2) \left[ -\frac{1}{4m^2} f_0(M^2/m^2) - \frac{1}{2m^2} f_1(M^2/m^2) \right],$$

$$d_{Ye}(|\mu_{\tilde{g}_Y}|^2) = \frac{\alpha_Y}{4\pi} m_e \tan \beta \times \hat{D}_{M^2}(|\mu_{\tilde{g}_Y}|^2,|\mu_H|^2) \left[ -\frac{1}{4m^2} f_0(M^2/m^2) \right]$$

$$+ \frac{\alpha_Y}{4\pi} m_e \tan \beta \times \frac{\partial}{\partial_m} \left[ \frac{1}{2m^2} f_0(|\mu_{\tilde{g}_Y}|^2/m^2) \right],$$

$$d_{Ye'}(|\mu_{\tilde{g}_Y}|^2) = \frac{\alpha_Y}{4\pi} m_e \times \frac{\partial}{\partial_m} \left[ \frac{1}{2m^2} f_0(|\mu_{\tilde{g}_Y}|^2/m^2) \right]. \quad (A.27)$$

Here, a finite-difference operator and mass functions $f_0$ and $f_1$ are defined as

$$\hat{D}_{z}^{(x,y)}[f(z)] \equiv \frac{f(x) - f(y)}{x - y} \quad (A.28)$$

and

$$f_0(x) = \frac{1}{(1-x)^3} (1 - x^2 + 2x \log x),$$

$$f_0(x) = \frac{1}{(1-x)^3} (3 - 4x + x^2 + 2 \log x). \quad (A.29)$$
For down quark, the functions in Eq. (3.16) are

\[
\begin{align*}
    d_3^{(d)}(\mu_{\bar{g}3}) &= \frac{\alpha_3}{4\pi} m_d \tan \beta \times \frac{\partial}{\partial m^2} \left[ \frac{4}{9m^2} f_0(\mu_{\bar{g}3}^2/m^2) \right], \\
    d_2^{(d)}(\mu_{\bar{g}2}) &= \frac{\alpha_2}{4\pi} m_d \tan \beta \times \hat{D}_{M^2}^{(\mu_{\bar{g}2}^2,\mu_H^2)} \left[ \frac{1}{4m^2} f_0(M^2/m^2) - \frac{1}{2m^2} f_1(M^2/m^2) \right], \\
    d_4^{(d)}(\mu_{\bar{g}y}) &= \frac{\alpha_y}{4\pi} m_d \tan \beta \times \hat{D}_{M^2}^{(\mu_{\bar{g}y}^2,\mu_H^2)} \left[ -\frac{1}{12m^2} f_0(M^2/m^2) \right], \\
    d_Y^{(d)}(\mu_{\bar{g}y}) &= \frac{\alpha_y}{4\pi} m_d \tan \beta \times \frac{\partial}{\partial m^2} \left[ -\frac{1}{54m^2} f_0(\mu_{\bar{g}y}^2/m^2) \right], \\
    d_3^{(d)'}(\mu_{\bar{g}3}) &= \frac{\alpha_3}{4\pi} m_d \tan \beta \times \frac{\partial}{\partial m^2} \left[ \frac{4}{9m^2} f_0(\mu_{\bar{g}3}^2/m^2) \right], \\
    d_Y^{(d)'}(\mu_{\bar{g}y}) &= \frac{\alpha_y}{4\pi} m_d \times \frac{\partial}{\partial m^2} \left[ -\frac{1}{54m^2} f_0(\mu_{\bar{g}y}^2/m^2) \right],
\end{align*}
\]

(A.30)

and

\[
\begin{align*}
    \bar{d}_3^{(d)}(\mu_{\bar{g}3}) &= \frac{\alpha_3}{4\pi} m_d \tan \beta \times \frac{\partial}{\partial m^2} \left[ \frac{1}{6m^2} f_0(\mu_{\bar{g}3}^2/m^2) - \frac{3}{2m^2} f_1(\mu_{\bar{g}3}^2/m^2) \right], \\
    \bar{d}_2^{(d)}(\mu_{\bar{g}2}) &= \frac{\alpha_2}{4\pi} m_d \tan \beta \times \hat{D}_{M^2}^{(\mu_{\bar{g}2}^2,\mu_H^2)} \left[ \frac{3}{4m^2} f_0(M^2/m^2) \right], \\
    \bar{d}_4^{(d)}(\mu_{\bar{g}y}) &= \frac{\alpha_y}{4\pi} m_d \tan \beta \times \hat{D}_{M^2}^{(\mu_{\bar{g}y}^2,\mu_H^2)} \left[ \frac{1}{4m^2} f_0(M^2/m^2) \right] \\
    &\quad + \frac{\alpha_y}{4\pi} m_d \tan \beta \times \frac{\partial}{\partial m^2} \left[ \frac{1}{18m^2} f_0(\mu_{\bar{g}y}^2/m^2) \right], \\
    \bar{d}_3^{(d)'}(\mu_{\bar{g}3}) &= \frac{\alpha_3}{4\pi} m_d \times \frac{\partial}{\partial m^2} \left[ \frac{1}{6m^2} f_0(\mu_{\bar{g}3}^2/m^2) - \frac{3}{2m^2} f_1(\mu_{\bar{g}3}^2/m^2) \right], \\
    \bar{d}_Y^{(d)'}(\mu_{\bar{g}y}) &= \frac{\alpha_y}{4\pi} m_d \times \frac{\partial}{\partial m^2} \left[ \frac{1}{18m^2} f_0(\mu_{\bar{g}y}^2/m^2) \right].
\end{align*}
\]

(A.31)
For up quark,

\[ d_3^{(u)} (|\tilde{\mu}_3|^2) = \frac{\alpha_3}{4\pi} m_u \tan^{-1} \beta \times \frac{\partial}{\partial m^2} \left[ -\frac{8}{9m^2} f_0(|\mu_{\tilde{\mu}_3}|^2/m^2) \right], \]

\[ d_2^{(u)} (|\mu_3|^2) = \frac{\alpha_2}{4\pi} m_u \tan^{-1} \beta \times \frac{\partial}{\partial m^2} \left[ \frac{1}{2m^2} f_0(M^2/m^2) \right], \]

\[ d_Y^{(u)} (|\mu_3|^2) = \frac{\alpha_Y}{4\pi} m_u \tan^{-1} \beta \times \frac{\partial}{\partial m^2} \left[ \frac{1}{6m^2} f_0(M^2/m^2) \right] + \frac{\alpha_Y}{4\pi} m_u \tan^{-1} \beta \times \frac{\partial}{\partial m^2} \left[ -\frac{2}{27m^2} f_0(|\mu_3|^2/m^2) \right], \]

\[ d_3^{(u)'} (|\mu_3|^2) = \frac{\alpha_3}{4\pi} m_u \times \frac{\partial}{\partial m^2} \left[ -\frac{8}{9m^2} f_0(|\mu_{\tilde{\mu}_3}|^2/m^2) \right], \]

\[ d_Y^{(u)'} (|\mu_3|^2) = \frac{\alpha_Y}{4\pi} m_u \times \frac{\partial}{\partial m^2} \left[ -\frac{2}{27m^2} f_0(|\mu_3|^2/m^2) \right], \quad (A.32) \]

and

\[ \tilde{d}_3^{(u)} (|\mu_{\tilde{\mu}_3}|^2) = \frac{\alpha_3}{4\pi} m_u \tan^{-1} \beta \times \frac{\partial}{\partial m^2} \left[ \frac{1}{6m^2} f_0(|\mu_{\tilde{\mu}_3}|^2/m^2) - \frac{3}{2m^2} f_1(|\mu_{\tilde{\mu}_3}|^2/m^2) \right], \]

\[ \tilde{d}_2^{(u)} (|\mu_{\tilde{\mu}_3}|^2) = \frac{\alpha_2}{4\pi} m_u \tan^{-1} \beta \times \frac{\partial}{\partial m^2} \left[ \frac{3}{4m^2} f_0(M^2/m^2) \right], \]

\[ \tilde{d}_Y^{(u)} (|\mu_{\tilde{\mu}_3}|^2) = \frac{\alpha_Y}{4\pi} m_u \tan^{-1} \beta \times \frac{\partial}{\partial m^2} \left[ \frac{1}{4m^2} f_0(M^2/m^2) \right] + \frac{\alpha_Y}{4\pi} m_u \tan^{-1} \beta \times \frac{\partial}{\partial m^2} \left[ -\frac{1}{9m^2} f_0(|\mu_{\tilde{\mu}_3}|^2/m^2) \right], \]

\[ \tilde{d}_3^{(u)'} (|\mu_{\tilde{\mu}_3}|^2) = \frac{\alpha_3}{4\pi} m_u \times \frac{\partial}{\partial m^2} \left[ \frac{1}{6m^2} f_0(|\mu_{\tilde{\mu}_3}|^2/m^2) - \frac{3}{2m^2} f_1(|\mu_{\tilde{\mu}_3}|^2/m^2) \right], \]

\[ \tilde{d}_Y^{(u)'} (|\mu_{\tilde{\mu}_3}|^2) = \frac{\alpha_Y}{4\pi} m_u \times \frac{\partial}{\partial m^2} \left[ -\frac{1}{9m^2} f_0(|\mu_{\tilde{\mu}_3}|^2/m^2) \right]. \quad (A.33) \]
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