Hadronic electric dipole moments, the Weinberg operator, and light gluinos

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We reexamine questions concerning the contribution of the three-gluon Weinberg operator to the electric dipole moment of the neutron, and provide several QCD sum-rule-based arguments that the result is smaller than—but nevertheless consistent with—estimates which invoke naive dimensional analysis. We also point out a regime of the minimal supersymmetric standard model parameter space with light gluinos for which this operator provides the dominant contribution to the neutron electric dipole moment due to enhancement via the dimension five color electric dipole moment of the gluino.

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I. INTRODUCTION AND SUMMARY

New sources of CP violation in supersymmetric extensions of the standard model are highly constrained by the null experimental results for the electric dipole moments (EDMs) of neutrons and heavy atoms [1,2]. Typically, when the superpartners have an electroweak scale mass, Λ_W, the additional CP violating phases are constrained to be of $O(10^{-2})$. When confronted with the natural expectation that the size of these phases in the soft-breaking sector should be of order one, this creates a problem for weak-scale supersymmetry (SUSY).

The interactions which generate EDMs are described by a CP-odd effective Lagrangian, induced at 1 GeV by integrating out heavy standard model particles and superpartners, which contains a series of operators of increasing dimension. The leading θ term,

$$\mathcal{L}_{\rm eff}^{[4]}=\frac{g_s^2}{32\pi^2}\bar{\theta}G^a_{\mu\nu}\tilde{G}^a_{\mu\nu}, \tag{1}$$

has dimension four, and an arbitrary value for $\bar{\theta}$ constitutes the usual strong CP problem as its contribution to EDMs is unsuppressed by any heavy scale. Moreover, the existence of additional CP-odd phases in the soft-breaking sector of the minimal supersymmetric standard model (MSSM) aggravates this problem by inducing a large additive renormalization of $\bar{\theta}$ that survives in the decoupling limit. The conventional ‘‘cure’’—the Peccei-Quinn mechanism—eliminates $\bar{\theta}$ and leaves the dimension five quark EDMs and color EDMs (CEDMs),

$$\mathcal{L}_{\rm eff}^{[5]}=-\frac{i}{2}\sum_{i=e,u,d,s} d_i\bar{\psi}_i(F\sigma)\gamma_5\psi_i -\frac{i}{2}\sum_{i=u,d,s}\tilde{d}_i\bar{\psi}_i(G\sigma)\gamma_5\psi_i, \tag{2}$$

and the Weinberg operator [3],

$$\mathcal{L}_{\rm eff}^{[6]}=\frac{1}{3}w\, f^{abc}G^a_{\mu\nu}\tilde{G}^b_{\nu\beta}G^c_{\beta\mu}, \tag{3}$$

as the dominant mediators of CP violation from the soft breaking sector to the observables. Note that although the quark (C)EDMs have dimension five, chiral symmetry requires that the corresponding coefficients are proportional to a light quark mass, and thus d_i, $\tilde{d}_i$, and w generically scale in the same way with the overall SUSY breaking scale.

Extracting constraints on the underlying CP-odd phases thus requires quantitative knowledge of the dependence of observable EDMs on d_i, $\tilde{d}_i$, and w normalized at the hadronic scale. Recently, the dependence on d_i and $\tilde{d}_i$ has been determined more precisely using QCD sum rules [5], and now we turn our attention to the Weinberg operator. Although rather intractable within the standard framework, we will present several sum-rule-based estimates. The resulting preferred range for the neutron EDM,

$$d_n(w)=e(10\text{–}30)\ \text{MeV}\, w(1\ \text{GeV}), \tag{4}$$

is a factor of two smaller than conventional estimates [3,4] using ‘‘naive dimensional analysis’’ (NDA) [6]. This moderate suppression can be understood through the appearance of combinatoric factors which are not accounted for within NDA. However, while our result for $d_n(w)$ is smaller than the NDA estimate, and thus $d_n(d_i,\tilde{d}_i)$ generally dominates the contributions to d_n, there is a regime in which $d_n(w)$ is important as it is generated rather differently from the quark (C)EDMs within the MSSM.

In order to explain this point recall, first of all, that there are several generic ‘‘strategies’’ for curing the SUSY CP problem. The first is to require that the superpartners are heavy enough to suppress all operators of dim$\geqslant$5 generated

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at the SUSY threshold. This decoupling is usually applied to sfermions of the first two generations only, in order to avoid problems with fine tuning in the Higgs sector. However, this approach is only partially successful as relatively large EDMs may be generated through higher loops [7] or through four-fermion operators induced by Higgs exchange [8]. Secondly, one could conceive of a universal conspiracy leading to cancellations between different contributions [9], but this is difficult to reconcile with the null results for all three types of EDM measurement (neutron, paramagnetic and diamagnetic atoms) that *a priori* have different phase dependence [10]. A third, perhaps more elegant, option is to invoke an exact CP or parity at some high-energy scale and specify the mechanisms that break supersymmetry in such a way that all the relevant soft breaking parameters are rendered real. This could also be one way of obviating the need for axion relaxation [11]. However, some of these scenarios may face problems when confronted with the large CP violation that is by now well documented in the B-meson system [12].

Given these difficulties, one may pursue another option which is to suppress the SUSY contributions by creating some (mild) hierarchies between the soft breaking parameters in order to suppress the EDMs generated at one loop. Notably, in the limit where gauginos are much lighter than the sfermions, all one-loop contributions to the EDMs of light quarks and the electron take the following form:

$$d_i(\text{one loop}) \sim (\text{loop factor}) \times \frac{m_i}{m_{\text{sf}}^4} \text{Im}(m_\lambda A), \tag{5}$$

with a similar expression for d_i induced by the relative phase of μ and m_λ. Here $i=e, u, d, s$, and m_{sf} stands for a generic sfermion mass. It is easy to see that as $m_\lambda \to 0$ the expression (5) for d_i vanishes. Thus a mild hierarchy $m_\lambda \sim (10^{-3}$–$10^{-2}) m_{\text{sf}}$ would appear to be sufficient to evade the SUSY CP problem [13,14]. In slightly different language, it follows from Eq. (5) that in this regime the quark EDMs are demoted to dimension seven operators and thus are relatively harmless.

While the quark EDMs are suppressed by this hierarchy, we emphasize that sending m_λ down to hadronic scales actually *enhances* the neutron EDM via the generation of the Weinberg operator. The main point is that the gluino CEDM,

$$\mathcal{L}_\lambda = \frac{1}{4} \tilde{d}_\lambda f^{abc} \overline{\lambda^b \sigma} \cdot G^a \gamma_5 \lambda^c, \tag{6}$$

can be induced by a top-stop loop [15], leading to

$$\tilde{d}_\lambda(\text{one loop}) \sim (\text{loop factor}) \times \frac{m_t^2}{m_{\text{sf}}^4} \text{Im}(A_t - \mu^* \cot\beta) \tag{7}$$

in a basis in which the gluino mass is real. Thus $\tilde{d}_\lambda$ is a genuine dimension five operator, $\tilde{d}_\lambda \sim 1/\Lambda_W$, for $A_t \sim \mu \sim m_{\text{sf}} \sim \Lambda_W$. It follows that for $m_\lambda \sim \Lambda_{\text{hadr}}$, the gluino takes part in the strong interactions and contributes to the energy density of all hadrons. Consequently the neutron EDM is unsuppressed by any additional scale, and at a crude level $d_n \sim e\tilde{d}_\lambda \sim 1/\Lambda_W$.

This enhancement by the gluino CEDM (6) persists in the intermediate hierarchical regime $\Lambda_{\text{hadr}} \ll m_\lambda \ll \Lambda_W$ where, on integrating out the gluino, one generates a contribution to the Weinberg operator alluded to above that scales as $1/(m_\lambda \Lambda_W)$. At a critical scale $m_\lambda = m_\lambda^{\text{int}} < \Lambda_W$ these contributions will dominate over $d_i \sim \Lambda_{\text{hadr}} m_\lambda / \Lambda_W^3$ and d_n will start increasing while m_λ decreases. As we will determine below, the scale

$$m_\lambda^{\text{int}} \sim (6-12) \ \text{GeV} \tag{8}$$

sets an effective threshold for the maximal suppression of EDMs possible with this superpartner hierarchy.[1]

Our results suggest that at this scale the neutron EDM is still considerably larger than the experimental bound,

$$\Delta d_n(m_\lambda^{\text{int}}) \sim (40-80) d_n^{\text{exp}}, \tag{9}$$

unless the SUSY CP phases are fine tuned. Note that both m_λ^{int} and $d_n(m_\lambda^{\text{int}})$ depend, in addition, on possible intergenerational hierarchies for the squark masses. When the first generation of sfermions is taken to be heavier than Λ_W, m_λ^{int} increases while $d_n(m_\lambda^{\text{int}})$ decreases.

Therefore, the Weinberg operator has an important role to play in minimizing the suppression possible within the light gluino regime. Note that for $m_\lambda \ll \Lambda_{\text{hadr}}$, the CP-violating phase can be rotated to m_λ itself leading to a suppression of d_n by $m_\lambda / \Lambda_{\text{hadr}}$ as one approaches the super-Yang-Mills limit. A schematic plot of the behavior of $d_n(m_\lambda)$ is shown in Fig. 1.

In the next section we turn to the problem of estimating the contribution to d_n induced by the Weinberg operator, justifying the result (4). Then, in Sec. III we describe in more detail the calculation justifying the argument outlined above which uses the Weinberg operator to limit the suppression of EDMs for light gluinos.

II. NEUTRON EDM INDUCED BY THE WEINBERG OPERATOR

Unlike the case of d_n induced by the θ term, or the EDMs and CEDMs of quarks, where chiral loop [18] and QCD sum-rule-based calculations [5] are available, the matrix element that relates d_n with the Weinberg operator is unknown. The standard estimate, first obtained by Weinberg [3], makes use of "naive dimensional analysis" [6,19] which keeps track of dimensions, in terms of the generic hadronic scale Λ_{hadr}, and Goldstone-mediated interactions through the effective dimensionless coupling $\Lambda_{\text{hadr}}/f_\pi$. One finds [3,4]

[1]As an aside, we note that (perhaps surprisingly) a gluino mass of order m_λ^{int} is still not ruled out by direct constraints, and indeed has recently been revived [16] in relation to the enhanced hadronic b-quark production observed at the Collider Detector at Fermilab (CDF) and DØ [17].

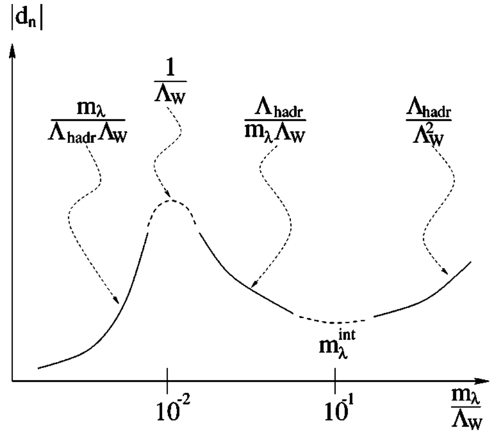

FIG. 1. Schematic behavior of the neutron EDM d_n as a function of the gluino mass. Lowering m_λ from the SUSY threshold there is an initial suppression of d_n due to the decrease of $d_i(m_\lambda)$ as m_λ decreases from Λ_W to the intermediate value $m_\lambda^{\rm int}$. A further decrease of m_λ in the interval from $m_\lambda^{\rm int}$ to $\Lambda_{\rm hadr}$ leads to the increase of d_n due to the contribution of the Weinberg operator, induced by the gluino CEDM. When m_λ is smaller than $\Lambda_{\rm hadr}$, d_n receives a linear suppression by m_λ.

$$d_n \sim e\frac{\Lambda_{\rm hadr}}{4\pi}w(\mu)\sim e \;\; 90 \;\; {\rm MeV} \;\; w(\mu), \tag{10}$$

at a low-energy normalization point μ, taking $\Lambda_{\rm hadr}\sim 4\pi f_\pi \sim 1.2$ GeV. The large value $\sim 4\pi$ for the coupling amounts to demanding that loop corrections are qualitatively similar to the tree-level terms at the matching scale. In the gluonic sector, which is important here, this means that within the UV quark and gluon description the relevant value of the gauge coupling is necessarily very large and consequently the inferred matching scale does not mesh easily with expectations from the chiral sector [6,19]. In the present context Weinberg [3], and many papers since [4], have, for the purpose of evaluating the gauge coupling, chosen a specific matching scale corresponding to $g_s=4\pi/\sqrt{6}$, or $\alpha_s\simeq 2$ [cf. $\alpha_s(1\text{ GeV})\simeq 0.4$]. If we adopt this normalization scale in Eq. (10), and use (somewhat optimistically) the one-loop anomalous dimension for w [20], the relation $w[\mu(g_s=4\pi/\sqrt{6})]\simeq 0.4w(1\text{ GeV})$ leads to the most commonly used estimate for $d_n(w)$:

$$d_n^{(1)}\sim e\; 40 \text{ MeV } w(\mu=1 \text{ GeV}). \tag{11}$$

We will avoid quoting a result for the dependence of d_n on $w(\Lambda_W)$, as there are additional threshold contributions from $\tilde{d}_b$ and $\tilde{d}_c$ generated by top-quark–top-squark–gluino loops, which are in general model dependent [21].

To get some intuition regarding the estimate (11), we can consider more carefully the loop factors which are effectively set to unity in Eq. (10). For illustration, consider reducing the Weinberg operator to the EDM by ''integrating out'' the gluons. This leads to an effective loop factor of $g_s^3/(4\pi)^4$ which reproduces Eq. (10) provided we take $g_s\sim 4\pi$. One obtains a similar conclusion for the effective scale by considering the gauge kinetic term itself [6]. As a consistent matching condition we might then choose $\mu=\mu(g_s=4\pi)$, leading to a result

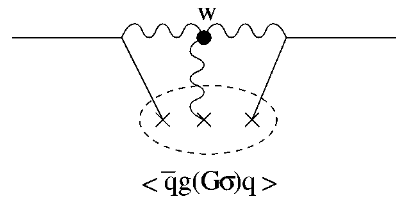

FIG. 2. Perturbative insertion of the Weinberg operator into a quark line. The resulting correction to the propagator is proportional to $w\gamma_5\langle\bar{q}g_s(G\sigma)q\rangle$.

$$d_n^{(2)}\sim e\; 18 \text{ MeV } w(\mu=1 \text{ GeV}), \tag{12}$$

which is half the size of Eq. (11). Although both results (11) and (12) are consistent within the expected precision of the NDA, it is clear that independent quantitative calculations are needed to determine $d_n(w)$ to better than an order of magnitude.

As a quantitative test of the NDA estimates, we will now revisit the calculation of $d_n(w)$ using QCD sum rules, leading to a result that is a factor of 2 *smaller* than Eq. (11) and consistent with Eq. (12). To proceed, we note first that the leading contribution to the EDM from the operator product expansion (OPE) of the nucleon current correlator in the presence of the source (3) exhibits a logarithmic infrared divergence. This signals [22] the presence of additional operators, required to resolve the divergence, whose contributions are generally rather difficult to calculate directly. Therefore, we will be content to regulate the logarithmic-divergent contributions with an IR cutoff. These terms will then form the basis of our estimates as they are correspondingly enhanced and thus provide the dominant contributions to the EDM.

We begin by noting that the Weinberg operator allows for a perturbative insertion into the quark propagator. The leading CP-odd correction is described by the diagram shown in Fig. 2, and standard manipulations [23] lead to the following result:

$$iS(p)=\frac{i\not{p}}{p^2}+\frac{ig_s w}{8p^4}\gamma_5\langle\bar{q}g_s(G\sigma)q\rangle, \tag{13}$$

where the value of the quark-gluon condensate is given by [24]

$$\langle\bar{q}g_s(G\sigma)q\rangle=m_0^2|\langle\bar{q}q\rangle|\simeq 0.8 \;\; {\rm GeV}^2\langle\bar{q}q\rangle, \tag{14}$$

with $\langle\bar{q}q\rangle=-(230\text{ MeV})^3$. It is the $1/p^4$ momentum dependence in the second term of Eq. (13) which leads to the logarithmic infrared divergence alluded to above in the correlator of two nucleon currents. This signals the breakdown of the OPE, but also singles out this insertion as providing the dominant effect which we will use in calculating $d_n(w)$. The ambiguity of the infrared logarithm does of course render the result less reliable than the corresponding determination of $d_n(d_i,\tilde{d}_i)$ [5], but nonetheless sufficient for our estimates.

The insertion present in the second term in Eq. (13) behaves as a ''soft γ_5 mass.'' Indeed, while irrelevant for large p^2, at hadronic scale momenta it mimics the existence of an effective CP-odd mass of order

$$m_5^{\rm eff} \simeq \frac{g_s w}{8\Lambda_{\rm hadr}^2}\langle \overline{q} g_s (G\sigma) q\rangle \sim (120 \ \ {\rm MeV} - 160 \ \ {\rm MeV})^3 w \tag{15}$$

where $\Lambda_{\rm hadr}$ is the effective hadronic scale, which we take to lie in a range from m_ρ up to $4\pi f_\pi$, with previous results [5] suggesting that the lower end of this range is most relevant for EDM observables. This determines the neutron EDM according to the scaling relation [5],

$$d_n \sim e\frac{|m_5^{\rm eff}|}{\Lambda_{\rm hadr}^2} \sim e(1.5-7) \ \ {\rm MeV} \ \ w(\Lambda_{\rm hadr}), \tag{16}$$

where the large range in this estimate arises from the allowed variation in $\Lambda_{\rm hadr}$.

This result is 5–10 times smaller than the conventional NDA estimate (11). This is actually not too surprising once we recall that *a priori* $d_n(w)$ should be of $O(\langle \overline{q}q\rangle^0)$ in the chiral limit, while the contribution in Eq. (16) is $O(\langle \overline{q}q\rangle)$ and thus may indeed be subleading. To test this one can consider an explicit sum-rules-based estimate [25] utilizing the insertion (13). One finds that for the natural chirally invariant Lorentz structure, $\{\not p,(F\sigma)\gamma_5\}$ [5], the tractable contributions are of $O(\langle \overline{q}q\rangle^2)$ and render a result for d_n within the range (16). Previous experience [5] would suggest that the terms of $O(\langle \overline{q}q\rangle^2)$ are sub-dominant, but unfortunately the (*a priori*) leading contributions of $O(\langle \overline{q}q\rangle^0)$ for $d_n(w)$ are intractable in this direct approach due to the presence of unknown condensates.

This analysis suggests that the range (16) might represent an underestimate of $d_n(w)$. A natural path to follow is to consider the sum rules in chirally variant channels such as $(F\sigma)$ or $\not p(F\sigma)\not p$ from which one can still extract $d_n(w)$ along the lines considered previously for $d_n(\theta)$ [26]. A convenient means of estimating $d_n(w)$ in this vein is to calculate the γ_5 rotation of the nucleon wave function induced by the Weinberg operator and determine d_n in terms of the corresponding rotation of the neutron anomalous magnetic moment μ_n:

$$d_n \sim \mu_n \frac{\langle N|\frac{w}{3}(GG\tilde{G})|N\rangle}{m_n \overline{N} i\gamma_5 N}. \tag{17}$$

This approach was considered previously by Bigi and Uraltsev [27] who estimated $\langle N|(GG\tilde{G})|N\rangle$ in terms of $\langle N|GG|N\rangle$ and the corresponding vacuum condensates.

We can follow this route and perform a more explicit calculation by evaluating the ''γ_5'' term in the standard mass sum-rule correlator of the two nucleon currents. For the conventional choice of the Ioffe interpolating current for the neutron η [28], we obtain at leading order,

$$\begin{aligned}\int d^4x e^{ip\cdot x}\langle \eta(0)\overline{\eta}(x)\rangle &= \frac{1}{16\pi^2}p^2 \ln(-\Lambda_{\rm UV}^2/p^2)\langle \overline{q}q\rangle \\ &\times\left[1 + i\gamma_5\frac{3g_s w}{32\pi^2}m_0^2\ln(-p^2/\mu_{IR}^2)\right] + \cdots .\end{aligned} \tag{18}$$

It is the relative coefficient between the structures **1** and $i\gamma_5$ that determines the chiral rotation and consequently enters into the estimate of [27]. From Eqs. (17) and (18) we obtain

$$d_n \simeq \mu_n \frac{3g_s m_0^2}{32\pi^2} w \ln(M^2/\mu_{IR}^2) \simeq e \ \ 22 \ \ {\rm MeV} \ \ w(1 \ \ {\rm GeV}), \tag{19}$$

where we took $M/\mu_{IR}=2$ and $g_s=2.1$. It is important to note that the estimate (19) arises at $O(\langle \overline{q}q\rangle^0)$, which we would expect to be dominant, and is indeed considerably larger than the estimate (16). A more involved calculation of the nucleon current correlator in an external electromagnetic field [25] reveals additional contributions to $d_n(w)$, but the overall result remains quite close to Eq. (19). Additional induced corrections, from Peccei-Quinn relaxation, would also be subleading [27] as they cannot contribute at $O(\langle \overline{q}q\rangle^0)$.

The only other QCD sum-rules estimate of $d_n(w)$ that we are aware of was made by Khatsimovsky [29] who considered a high order term in the OPE proportional to the dimension-eight operator $F(GG\tilde{G})$. An estimate of the nonlocal correlator, $\int d^4x\langle 0|T\{(GG\tilde{G})(0),(GG\tilde{G})(x)\}|0\rangle$ produced a result for $d_n(w)$ similar to Eq. (10). However, combinatoric factors were ignored in this calculation which clearly reduce the result to a value consistent with—or somewhat smaller than—Eqs. (16),(19). In practice a precise calculation along these lines does not appear feasible, as multiple perturbative insertions of the gluon field strength into a quark line generally leads to power-like infrared divergences [22], signifying the breakdown of the OPE.

Putting these results together, and ignoring the lower range of Eq. (16) for the reasons discusssed above, we find the preferred range for $d_n(w)$,

$$d_n(w) \sim e(10\text{–}30) \ {\rm MeV}\, w(1 \ {\rm GeV}), \tag{20}$$

which is a factor of two smaller than the conventional NDA estimate (11), and consistent with Eq. (12). This result will be discussed in more detail elsewhere [25], but we turn now to a regime of the SUSY parameter space for which this contribution to d_n is nonetheless very significant.

III. ENHANCEMENT VIA GLUINO COLOR EDM

As described and schematically illustrated in Sec. I, the neutron EDM is particularly enhanced in the domain $\Lambda_{\rm hadr} \ll m_\lambda \ll \Lambda_W$ where the gluino develops a color EDM via a top-quark–top-squark loop [15],

$$\tilde{d}_\lambda(\Lambda_W) = -\frac{g_s^3(\Lambda_W)}{32\pi^2}\frac{m_t}{M^2_{\tilde{t}_1}}\sin(2\theta_{\tilde{t}})\sin\delta_t \times\left[f_g\left(\frac{m_t^2}{M^2_{\tilde{t}_1}}\right) - \frac{M^2_{\tilde{t}_1}}{M^2_{\tilde{t}_2}} f_g\left(\frac{m_t^2}{M^2_{\tilde{t}_2}}\right)\right], \quad (21)$$

with $\delta_t = \mathrm{Arg}[A_t - \mu^\star \cot\beta]$ and $f_g(y) = [1-y+\ln(y)]/(1-y)^2$. Note that this expression is independent of m_λ and scales as $1/\Lambda_W$. The corresponding contribution to the Weinberg operator [3,4,30,31],

$$\Delta w(m_\lambda) = -\frac{3g_s^2(m_\lambda)}{32\pi^2}\frac{\tilde{d}_\lambda(m_\lambda)}{m_\lambda} \quad (22)$$

scales as $1/m_\lambda\Lambda_W$. It is worth noting that in addition to the obvious enhancement by a factor of Λ_W/m_λ relative to the standard scenario [30], the gluino CEDM-induced shift of the Weinberg operator is also enhanced relative to that induced by c or b quarks which is of order $1/\Lambda_W^2$ [31,21].

The normalization of Δw at the hadronic scale involves running the gluino CEDM from Λ_W down to the gluino mass threshold, and subsequent running of w down to Λ_{hadr}. For completeness, we give the one-loop β function coefficient, $\beta_0 = 11 - 2n_\lambda - 2n_q/3$, where n_x stands for the number of light x particles at the scale under concern. Besides this, the anomalous dimensions of $\tilde{d}_\lambda$ and w are given, respectively, by $\gamma^\lambda = -18 + \beta_0$ and $\gamma^W = -36 + 3\beta_0$. The latter has been computed in [20], and the computation of the former is similar to that of the quark color EDMs [32].

We now illustrate numerically the impact of light gluinos on the neutron EDM using the range for $d_n(w)$ in Eq. (4). Using Eqs. (21) and (22), we can write

$$\Delta d_n \sim 100\sin\delta_t \frac{(4\text{–}12)\ \mathrm{GeV}}{m_\lambda} d_n^{\exp}, \quad (23)$$

where we have taken $M_{\tilde{t}_1} = 200$ GeV, $M_{\tilde{t}_2} = 700$ GeV, $\theta_{\tilde{t}} = \pi/4$, and the current experimental bound on the neutron EDM is $d_n^{\exp} < 6\times 10^{-26}$ e cm [1]. The final results are presented in Fig. 3, where we have chosen the mid-value $d_n = 20$ MeV w in Eq. (4). The solid curve stands for $m_\lambda = 1$ GeV (with $d_n/d_n^{\exp} \approx 980$ at $\delta_t = \pi/2$), the dashed curve for $m_\lambda = m_b$ (with $d_n/d_n^{\exp} \approx 170$ at $\delta_t = \pi/2$), and the dot-dashed curve for $m_\lambda = 20$ GeV (with $d_n/d_n^{\exp} \approx 30$ at $\delta_t = \pi/2$). Thus, for light gluinos, where $m_\lambda \sim (1-4)$ GeV, one finds that $d_n(w)$ exceeds the experimental bound by at least two orders of magnitude throughout the entire preferred range in Eq. (4) unless the SUSY phases are tuned such that $\delta_t \lesssim 10^{-2}$.

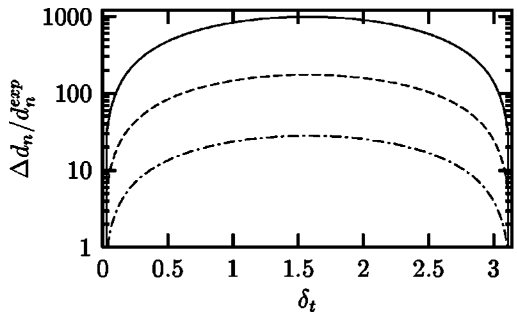

FIG. 3. The δ_t dependence of the gluino contribution to d_n for $m_\lambda = 1$ GeV (solid line), $m_\lambda = m_b$ (dashed line), and $m_\lambda = 20$ GeV (dot-dashed line).

Of particular interest is the maximal suppression that one can achieve for the EDM in this hierarchical regime with light gluinos. We denote by m_λ^{int} the critical scale at which the one-loop contribution induced by quark EDMs (and CEDMs) is approximately equal to the contribution associated with the Weinberg operator discussed here. Choosing the soft-breaking parameters in the first generation of squarks to be $\mathcal{O}(200\ \mathrm{GeV})$, and assuming no accidental cancellations, we find

$$m_\lambda^{\mathrm{int}} \sim (6-12)\ \mathrm{GeV} \quad (24)$$

accounting for the range in Eq. (4), for which the (minimal) correction to the EDM is approximately,

$$\Delta d_n(m_\lambda^{\mathrm{int}}) \sim (40\text{–}80) d_n^{\exp}, \quad (25)$$

which still exceeds the experimental bound by at least an order of magnitude unless the CP-odd phases are small.

It is interesting to compare our estimates for d_n with those one obtains when the gluino is heavy, $m_\lambda \sim \Lambda_W$. In this case, the Weinberg operator is first generated at the weak scale at two-loop order [30]. On including the contributions arising at the b-quark and c-quark thresholds, one finds that d_n obtained via Eq. (10) only exceeds $d_n^{\exp}$ by at most one order of magnitude [21]. Consequently, the light gluino scenario actually induces a larger contribution to d_n via the color EDM of the gluino. Thus, while it is possible to suppress the one-loop contributions to the EDMs of leptons and hadrons by taking light gauginos [14], the induced contribution to the Weinberg operator means that the constraints on the SUSY CP-odd phases are not correspondingly relaxed.

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