CP violation in beta decay and electric dipole moments

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The T-odd correlation coefficient D in nuclear β decay probes CP violation in many theories beyond the standard model. We provide an analysis for how large D can be in light of constraints from electric dipole moment (EDM) searches. We argue that the neutron EDM d_n currently provides the strongest constraint on D, which is 10–10^3 times stronger than current direct limits on D (depending on the model). In particular, contributions to D in leptoquark models (previously regarded as ''EDM safe'') are more constrained than previously thought. Bounds on D can be weakened only by fine-tuned cancellations or if theoretical uncertainties are larger than estimated in d_n. We also study implications for D from mercury and deuteron EDMs.

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I. INTRODUCTION

The search for CP violation beyond the standard model (SM) remains an open question at the forefront of nuclear physics, particle physics, and cosmology.[1] New CP violation is a generic feature of physics beyond the SM [1], and is likely required to explain the baryon asymmetry of the Universe [2]. Furthermore, unlike the SM Kobayashi-Maskawa (KM) phase [3], new CP violation may be unconnected with flavor and can be probed in systems of ''ordinary matter'' through searches for T violation in nuclear β decay and electric dipole moments (EDMs) of atoms, nucleons, and nuclei.

CP violation in β decay is manifested through T-odd triple product correlations [4]. (See Refs. [5–8] for reviews of fundamental symmetry tests in β decay.) In this work, we study the so-called D correlation, corresponding to the triple product $\langle \mathbf{J} \rangle \cdot \mathbf{p}_e \times \mathbf{p}_\nu$, where $\langle \mathbf{J} \rangle$ is nuclear polarization, and $\mathbf{p}_e$ ($\mathbf{p}_\nu$) is the $e^\pm$ (ν) momentum. It is useful to write $D \equiv D_t + D_f$ to delineate fundamental T violation (D_t) from T-even final state effects (D_f) [6]. In the SM, the KM phase contribution to D_t is vanishingly small [9]. Therefore, to the extent that D_f is computable or negligible, measurements of D directly probe CP violation beyond the SM.

To date, D has been measured for the neutron [10–15] and ^{19}Ne [16,17]. The best neutron D measurement has been obtained recently by the emiT Collaboration [15]:[2]

$$D_n = (-1.0 \pm 2.1) \times 10^{-4}. \tag{1}$$

Final state interactions give $D_f = \mathcal{O}(10^{-5})$ [18], and have been computed to an accuracy better than 1% [19]. Although D_n measurements so far agree with SM expectations, there remains (in principle) a discovery window for future experiments down to $D_n \sim 10^{-7}$. For ^{19}Ne, an average of previous measurements [16,17] gives

$$D_{\mathrm{Ne}} = (1 \pm 6) \times 10^{-4}, \tag{2}$$

which has reached a level comparable to final state interaction effects $D_f \sim 10^{-4}$ [17].

Measurements of EDMs (denoted d) are also sensitive to CP violation in and beyond the SM [20]. No EDM has yet been observed, but many future experiments await [21]. Currently, the most significant EDM bounds are for the neutron [22], atomic mercury (^{199}Hg) [23], atomic thallium (^{205}Tl) [24], and recently molecular YbF [25]. These null results provide important constraints on CP violation in the SM due to the θ_{QCD} phase associated with the strong interaction (present limits on d_n require $\theta_{\mathrm{QCD}} < 10^{-10}$ [26]), and on CP violation beyond the SM, such as in the minimal supersymmetric standard model (MSSM) [27,28]. On the other hand, these observables are rather insensitive to the KM phase, requiring many orders of magnitude increases in sensitivities (see Ref. [20] and references therein).

In this work, we compare D vs EDMs (in particular, d_n and d_{Hg}) as probes of CP violation beyond the SM. For a given model, any CP-odd phase contributing to D generates an ''irreducible'' EDM that can only be avoided by fine-tuned cancellations with other phases in the model. We compute the resulting bounds on D from EDMs in several new physics models: left-right symmetric models [29], MSSM with R-parity violation [30], models with exotic fermions [31], and leptoquark (LQ) models [32]. Most of these scenarios, and the resulting constraints from

[1]The discrete symmetries discussed herein are charge conjugation (C), parity (P), and time reversal (T) symmetries. Assuming CPT invariance, T violation is equated with CP violation.

[2]We have added in quadrature statistical and systematic errors quoted in Ref. [15].

EDMs, have been studied previously [6,33,34]. Here, we provide several improvements:

(i) We take into account recent improved computations of d_n [35] and $d_{\rm Hg}$ [36]. Large uncertainties in the sensitivity of $d_{\rm Hg}$ to the CP-odd isovector pion-nucleon coupling [36] have weakened this constraint, and the d_n bound currently provides the strongest limit on D_t.

(ii) In the literature, LQ contributions to D_t are regarded as being safe from EDM constraints [6,34]. We argue that D_t is in fact more constrained than previously thought. We also study implications for D from LQ searches at hadron colliders.

(iii) We compute for the first time D_t in the R-parity violating MSSM (with baryon-number violation), arising at one-loop order.

(iv) We provide a (partially) model-independent analysis that applies to all the aforementioned models except LQs, for which the current limit on d_n implies $D_t < 3 \times 10^{-7}$.

We emphasize that D is much cleaner theoretically than the EDMs constraining it, which rely on hadronic and nuclear computations. Moreover, any realistic model may contain many different CP-odd phases, to which D_t and EDMs are sensitive to different linear combinations. The bounds we derive may be negated if there exist accidental cancellations between phases entering EDMs, and we neglect this possibility in our analysis.

Our work is organized as follows. In Sec. II, we review CP violation in β decay. We also summarize theoretical computations of neutron, mercury, and deuteron EDMs from underlying CP-violating operators most relevant for constraining D_t. In Secs. III and IV, we study constraints on D_t from EDM bounds in several scenarios beyond the SM, focusing, in particular, on LQ models. We present our conclusions in Sec. V.

II. CP-VIOLATING OBSERVABLES

A. Beta decay

The most general set of β decay interactions can be parametrized at the quark level by an effective Lagrangian [5]

$$\mathcal{L}_\beta = -\frac{4G_F V_{ud}}{\sqrt{2}} \sum_{\alpha,\beta,\gamma} a^\gamma_{\alpha\beta} \bar{e}_\alpha \Gamma^\gamma \nu_e \bar{u} \Gamma_\gamma d_\beta + \text{h.c.} \tag{3}$$

The chiralities (L, R) of the electron and down quark are labeled by α, β. The index $\gamma = S, V, T$ labels whether the interaction is scalar ($\Gamma^S \equiv 1$), vector ($\Gamma^V \equiv \gamma^\mu$), or tensor ($\Gamma^T \equiv \sigma^{\mu\nu}/\sqrt{2}$). CP invariance is preserved in β decay if all ten complex coefficients

$$\begin{matrix} a^S_{LL}, & a^S_{LR}, & a^S_{RL}, & a^S_{RR}, & a^V_{LL}, \\ a^V_{LR}, & a^V_{RL}, & a^V_{RR}, & a^T_{LR}, & a^T_{RL} \end{matrix} \tag{4}$$

have a common phase (a^T_{LL}, a^T_{RR} terms are identically zero). At leading order in the SM, all parameters vanish except $a^V_{LL} = 1$. SM radiative corrections and new physics contributions to a^V_{LL} can play an important role in the extraction of V_{ud} (see, e.g., Ref. [37]), but for CP-violating observables they can be neglected as subleading effects. We also hereafter set $V_{ud} = 1$; correlations between D and EDMs depend on $|V_{ud}|$, but the $\mathcal{O}$(few%) deviation from $|V_{ud}| = 1$ is irrelevant compared to other theoretical uncertainties. We neglect possible flavor constraints by considering only couplings between first generation fermions. Last, we assume that β decay processes involve a single neutrino flavor eigenstate ν_e, and we allow for both L, R chiralities. Coefficients involving (sterile) right-handed neutrinos are only relevant provided these states are kinematically allowed in β decay.[3]

In terms of the parametrization in Eq. (3), D is given by [4]

$$D_t = \kappa\, \text{Im}(a^V_{LR} a^{V*}_{LL} + a^V_{RL} a^{V*}_{RR}) + \kappa \frac{g_S g_T}{g_V g_A} \text{Im}(a^S_{L+} a^{T*}_{LR} + a^S_{R+} a^{T*}_{RL}), \tag{5}$$

where $a^S_{L+} \equiv (a^S_{LL} + a^S_{LR})$ and $a^S_{R+} \equiv (a^S_{RL} + a^S_{RR})$. For initial (final) state nucleus of spin J (J'), the coefficient κ is

$$\kappa \equiv \frac{4 g_V g_A M_{\rm F} M_{\rm GT}}{g_V^2 M_{\rm F}^2 + g_A^2 M_{\rm GT}^2} \sqrt{\frac{J}{J+1}} \delta_{JJ'} \simeq \begin{cases} 0.87 & \text{for } n \\ -1.03 & \text{for } ^{19}\text{Ne}, \end{cases} \tag{6}$$

where $g_V = 1$, $g_A \approx 1.27$ [40], and $M_{\rm F}(M_{\rm GT})$ is the Fermi (Gamow-Teller) matrix element. Scalar and tensor form factors $g_{S,T}$, originally estimated in Ref. [41], have been computed using lattice techniques (see Ref. [42] and references therein). In this work, we neglect the scalar-tensor term in Eq. (5). The R coefficient, corresponding to the T-odd β decay correlation $\langle \mathbf{J} \rangle \cdot \boldsymbol{\sigma}_e \times \mathbf{p}_e$ where $\boldsymbol{\sigma}_e$ is $e^\pm$ polarization, is also sensitive to scalar- and tensor-type CP violation [5,6]. However, these couplings are correlated with CP-odd tensor and scalar electron-nucleon couplings, which are strongly constrained by ^{199}Hg [23] and ^{205}Tl [24] EDM bounds, respectively [43–47].

B. Electric dipole moments

EDM searches are sensitive to a wide class of CP-violating operators that can arise beyond the SM: CP-odd quark and lepton dipole moments, Weinberg's three-gluon operator [48], and four-fermion operators. Here, the most relevant one is a CP-odd four-quark operator $\mathcal{O}_{LR}$, given by

[3]Sterile neutrinos with eV-scale mass have been studied recently in connection with various neutrino anomalies (see, e.g., Ref. [38]), and important constraints are provided by cosmology [39]. We do not attempt to accommodate these issues here.

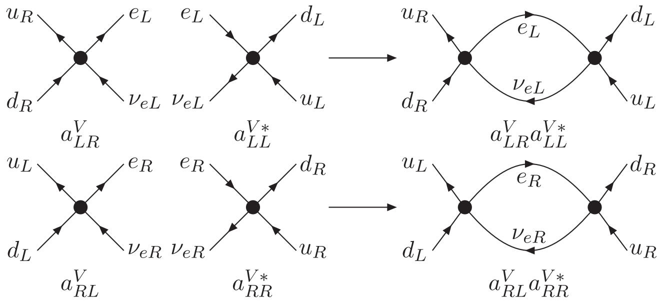

FIG. 1. *CP* violation entering $D_t = k\,\mathrm{Im}(a^V_{LR}a^{V*}_{LL} + a^V_{RL}a^{V*}_{RR})$ automatically generates the four-quark operator $\mathcal{O}_{LR} \equiv i(\bar{u}_L\gamma^\mu d_L \bar{d}_R\gamma_\mu u_R - \bar{d}_L\gamma^\mu u_L\bar{u}_R\gamma_\mu d_R)$, which contributes to neutron, mercury, and deuteron EDMs.

$$\mathcal{L}_{\rm eff} = -\frac{4G_F}{\sqrt{2}}k_{LR}\mathcal{O}_{LR},$$
$$\mathcal{O}_{LR} \equiv i(\bar{u}_L\gamma^\mu d_L\bar{d}_R\gamma_\mu u_R - \bar{d}_L\gamma^\mu u_L\bar{u}_R\gamma_\mu d_R), \quad (7)$$

where k_{LR} is the operator coefficient (normalized to $4G_F/\sqrt{2}$). Within the context of left-right symmetric models, this effective interaction arises from *CP*-violating W-W' mixing and has been studied previously [49–52]. We show in Fig. 1 that, by connecting the leptonic legs in a one-loop diagram, the same interference terms $a^V_{LR}a^{V*}_{LL}$ and $a^V_{RL}a^{V*}_{RR}$ contributing to D_t also generate $\mathcal{O}_{LR}$. Moreover, this diagram does not involve any chirality-changing mass insertions, and therefore is not suppressed by any light fermion masses. Other *CP*-odd operators (e.g., quark EDMs) also arise from new physics entering D_t, but are suppressed by light masses and will not be considered here.

The most significant EDM constraints on $\mathcal{O}_{LR}$ are for the neutron [22] and mercury atom [23]:

$$|d_n| < 2.9\times 10^{-26}\,e\,\mathrm{cm} \qquad (90\%\ \mathrm{CL}),$$
$$|d_{\rm Hg}| < 3.1\times 10^{-29}\,e\,\mathrm{cm} \qquad (95\%\ \mathrm{CL}). \quad (8)$$

Future measurements of the deuteron EDM d_D, expected at the level of $10^{-27}\,e\,$cm or better [53], will also provide important constraints on $\mathcal{O}_{LR}$.

Reference [35] has performed a systematic computation of d_n from *CP*-odd four-fermion operators, using a combination of chiral effective theory techniques and quark model estimates for the hadronic matrix elements. Using their results, we take

$$d_n = -1\times 10^{-19}k_{LR}\,e\,\mathrm{cm}, \quad (9)$$

with an $\mathcal{O}(1)$ uncertainty on the numerical prefactor [35].[4] Earlier results [44,49,51,54] are consistent at the order-of-magnitude level, but according to Ref. [35] are not as reliable in that they take into account different subsets of the full set of contributions to d_n.

Diamagnetic atoms (e.g., ^{199}Hg) are also sensitive to *CP*-odd four-quark interactions. Interpretation of these measurements is a three step process (see, e.g., Refs. [44,45]). First, atomic calculations relate the measured EDM to the nuclear Schiff moment S. For the case of mercury, we take [55]

$$d_{\rm Hg} = -2.6\times 10^{-17}\,e\,\mathrm{cm}\times\left(\frac{S_{\rm Hg}}{e\ \mathrm{fm}^3}\right). \quad (10)$$

This numerical value (2.6) agrees with an earlier result (2.8) by two of those authors [56], while another recent computation found a larger value (5.1) [57]. Second, the Schiff moment is computed in terms of P-, T-odd nucleon-pion couplings, of which only the isovector coupling $\bar{g}_1$ is relevant since $\mathcal{O}_{LR}$ is isovector [58]. Previous nuclear computations found (keeping only $\bar{g}_1$ terms) $S_{\rm Hg} = -0.071g\bar{g}_1\,e\,\mathrm{cm}^3$ [59] and $S_{\rm Hg} = -0.055g\bar{g}_1\,e\,\mathrm{cm}^3$ [60], where $g\approx 13.5$ is the (*CP*-even) pion-nucleon strong coupling. However, a recent and improved computation by Ref. [36] found that the $\bar{g}_1$ coefficient is very sensitive to the model-dependent nuclear potential inputs and may be suppressed by an order-of-magnitude (or more) and may have opposite sign compared to Refs. [59,60]. These nuclear physics uncertainties are crucial for constraining D_t using $d_{\rm Hg}$. In light of this unresolved issue, we take $|S_{\rm Hg}| = 0.01g|\bar{g}_1|\,e\,\mathrm{fm}^3$, remaining agnostic as to the sign (see Ref. [61] for additional discussion). Third, following Ref. [49], we conservatively take $\bar{g}_1 = 2\times 10^{-6}k_{LR}$. Reference [62] found a larger numerical prefactor by a factor of 7. Putting all these pieces together, we take

$$|d_{\rm Hg}| = 7\times 10^{-24}|k_{LR}|\,e\,\mathrm{cm}, \quad (11)$$

with an uncertainty at the order-of-magnitude level.

The deuteron EDM provides a much cleaner probe of $\bar{g}_1$ compared to $d_{\rm Hg}$. Following the recent computation of Ref. [63] (in good agreement with earlier results [64–66]), we take

$$|d_D| \approx 1.9\times 10^{-14}|\bar{g}_1|\,e\,\mathrm{cm} \approx 4.5\times 10^{-20}|k_{LR}|\,e\,\mathrm{cm}, \quad (12)$$

with $\mathcal{O}(20\%–30\%)$ uncertainty on the numerical factor (1.9) [63,67].

III. MODEL-INDEPENDENT BOUNDS ON D

New physics contributions to β decay can be organized in terms of a hierarchy of nonrenormalizable operators characterized by mass scale $\Lambda > G_F^{-1/2}$. Naively, the leading contributions to D_t will be those suppressed by the fewest powers of $(G_F\Lambda^2)^{-1}$, namely, from dimension-six operators contributing to a^V_{LR} that interfere with the SM amplitude a^V_{LL}. There is only one such operator [68,69]:

[4]This value is consistent with a naive estimate $d_n \sim eM_{\rm QCD}/\Lambda^2$, where $M_{\rm QCD}\sim 1$ GeV is the QCD scale and Λ is the scale of *CP* violation. Taking $\Lambda^{-2}\sim G_F k_{LR}$, we have $|d_n|\sim 2|k_{LR}|\times 10^{-19}\,e\,$cm. Also, it is useful to note $\mathcal{O}_{LR} = (\bar{u}u\bar{d}i\gamma_5 d - \bar{u}i\gamma_5 u\bar{d}d + 6\bar{u}t^a u\bar{d}i\gamma_5 t^a d - 6\bar{u}i\gamma_5 t^a u\bar{d}t^a d)/3$ using a Fierz transformation, where t^a is the $SU(3)_c$ generator, to make contact with the notation of Ref. [35].

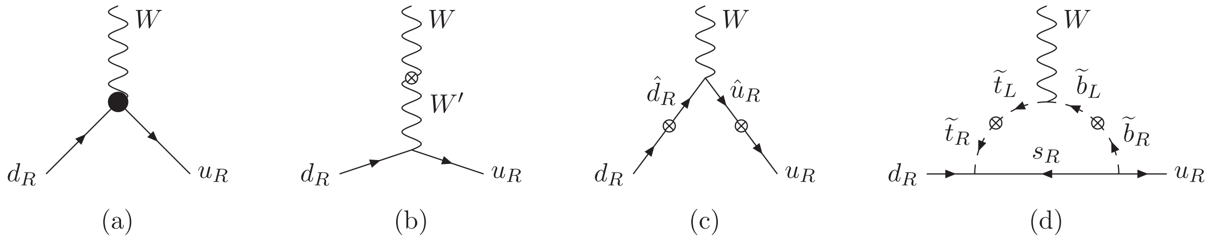

FIG. 2. (a) Effective $\bar{u}_R\gamma^\mu d_R W^+_\mu$ vertex arising beyond the SM, e.g., (b) left-right symmetric model with W-W' mixing; (c) exotic quarks $\hat{u}_R$, $\hat{d}_R$ with nonstandard $SU(2)_L \times U(1)_Y$ gauge couplings that mix with SM quarks u_R, d_R; (d) R-parity violating MSSM with baryon-number violation and squark left-right mixing. In each case, mixing insertions (involving the Higgs vacuum expectation value v) are denoted by ⊗.

$$\mathcal{L}_{\rm dim6} = \frac{c}{\Lambda^2}\bar{u}_R\gamma^\mu d_R i H^T \epsilon D_\mu H + \text{h.c.}, \tag{13}$$

where c is a complex coefficient. H is the Higgs doublet, D_μ is its covariant derivative, ϵ is the antisymmetric tensor ($\epsilon_{12} = -1$), and T denotes transpose acting on $SU(2)_L$ indices. Setting the Higgs field equal to its vacuum expectation value, Eq. (13) generates a coupling of the W boson to the right-handed charge current $\bar{u}_R\gamma^\mu d_R$, shown in Fig. 2. Integrating out the W boson, we obtain (recall we set $V_{ud} = 1$)

$$\mathcal{L}_{\rm dim6} = -\frac{c}{\Lambda^2}(\bar{u}_R\gamma^\mu d_R \bar{e}_L\gamma_\mu \nu_{eL} + \bar{u}_R\gamma^\mu d_R \bar{d}_L\gamma_\mu u_L) + \text{h.c.} \tag{14}$$

The operator of Eq. (13) generates at order $(G_F\Lambda^2)^{-1}$ contributions to both a^V_{LR} and k_{LR}:

$$\text{Im}\,(a^V_{LR}) = k_{LR} = \frac{\text{Im}(c)}{2\sqrt{2}G_F\Lambda^2}. \tag{15}$$

For all models that contribute to D_t via Eq. (13), EDMs are correlated with D_t in an otherwise model-independent way:

$$|d_n| = 1\times 10^{-19}\,e\,\text{cm}\times |D_t/\kappa|, \tag{16a}$$
$$|d_{\rm Hg}| = 7\times 10^{-24}\,e\,\text{cm}\times |D_t/\kappa|, \tag{16b}$$
$$|d_D| = 4.5\times 10^{-20}\,e\,\text{cm}\times |D_t/\kappa|. \tag{16c}$$

The current bound $|d_n| < 2.9\times 10^{-26}\,e\,\text{cm}$ [22] implies $|D_t/\kappa| < 3\times 10^{-7}$, far below present sensitivities.

This indirect limit on D_t applies to the following models:

(i) Left-right symmetric models with a W' boson that mixes with the W and couples to the right-handed quark charge current.

(ii) Models with exotic fermions with nonstandard gauge quantum numbers, e.g., exotic $SU(2)_L$-doublet vector quarks $\hat{u}$ and $\hat{d}$ that mix with the usual u and d quarks.

(iii) R-parity violating (RPV) MSSM with baryon-number violation, described below [70].

The relevant diagrams are shown in Fig. 2. The first two models were studied previously in connection with D in Refs. [6,33,34], and we do not describe them here.

The RPV MSSM is defined by adding to the MSSM superpotential gauge-invariant and renormalizable terms that violate either baryon or lepton number (but not both, to avoid proton decay) [30]. Contributions to D_t are generated by the baryon-number-violating terms[5]

$$W_{RPV} = \lambda''_{ijk} U^c_i D^c_j D^c_k, \tag{17}$$

where U^c_i, D^c_j are superfields corresponding to the (charge-conjugate) u^i_R and d^j_R quarks of generation i, j, respectively. Shown in Fig. 2, the leading contributions to D_t arise at one-loop from diagrams involving third generation squarks $\tilde{t}_{L,R}$ and $\tilde{b}_{L,R}$. This contribution relies on mixing between gauge eigenstates, described by (see, e.g., Ref. [72])

$$\mathcal{L}_{\rm mix} = -m_t(A^*_t \sin\beta + \mu\cos\beta)\tilde{t}^\dagger_L\tilde{t}_R - m_b(A^*_b\cos\beta + \mu\sin\beta)\tilde{b}^\dagger_L\tilde{b}_R + \text{h.c.}, \tag{18}$$

where $\tan\beta$ is the ratio between up- and down-type Higgs vacuum expectation values, $A_{t,b}$ and μ are MSSM mass parameters, and m_t (m_b) is the top (bottom) quark mass. For $\tan\beta \gg 1$, we have

$$a^V_{LR} = \frac{\lambda''_{123}\lambda''^*_{312}V_{tb}m_t m_b\tan\beta\,\mu A_t}{24\pi^2 m^4_{\tilde{q}}}, \tag{19}$$

assuming degenerate squarks with mass $m_{\tilde{q}}$ and treating Eq. (18) perturbatively by mass insertion. Bounds on n-$\bar{n}$ oscillations constrain $|\lambda_{312}| \lesssim 10^{-2}$ if all squarks have mass $m_{\tilde{q}} = 200$ GeV [73], but this bound is relaxed if only third generation squarks are light; $|\lambda_{123}|$ is unconstrained [30]. In principle Eq. (19) can be as large as $\mathcal{O}(10^{-3})$ for $m_{\tilde{q}}$, A_t, $\mu \sim 200$ GeV, $\lambda'' \sim 1$, and $\tan\beta \sim 50$ (perturbativity of the bottom Yukawa coupling requires

[5]Lepton number-violating terms have been studied previously in connection with the R coefficient [71].

$\tan\beta \lesssim 60$). However, the neutron EDM bound constrains $\mathrm{Im}(a_{LR}^V) < 3 \times 10^{-7}$, as per our previous discussion.

Reference [74] previously studied the RPV MSSM in connection with EDMs, focusing on contributions from quark and electron CP-odd dipole moments arising at two-loop. For the combination of RPV couplings λ'' in Eq. (19) entering D_t, quark EDM and chromo-EDM operators are suppressed by $m_{u,d}$. Here, the CP-odd four-quark operator gives a much stronger bound.

IV. LEPTOQUARK MODELS

Leptoquarks (LQs), fractionally charged colored states carrying baryon and lepton number, arise in many extensions of the SM, e.g., grand unification [75] and compositeness [76]. Here, we consider a phenomenological model of LQs coupled to first generation quarks and leptons [32]. LQ models have a rich phenomenology for β decay, potentially giving large contributions to D and other observables through tree-level processes [6].

In the literature, LQ models have been regarded as an "EDM safe" source of CP violation that might generate D as large as present experimental limits, without fine-tuning [6]. These previously considered models (dubbed the "usual scenarios") rely on LQ mixing to generate a dimension-eight operator contributing to a_{LR}^V at tree-level, which interferes with the SM amplitude a_{LL}^V [6]. In addition, scenarios involving LQs coupled to right-handed neutrinos can also generate D_t via the interference of two new physics amplitudes a_{RL}^V and a_{RR}^V.

In this section, we study in detail these cases (i.e., with or without right-handed neutrinos). We show that radiative corrections generate contributions to EDMs (in the spirit of Refs. [77,78]) sensitive to the same phases entering D_t. In both cases, the resulting bounds from the neutron EDM are stronger than the direct experimental limit.

A. Usual LQ scenarios: No right-handed neutrinos

There are two cases to consider: scalar and vector LQ exchange, both considered previously in Ref. [6]. Since both cases are similar, we treat them simultaneously. The relevant LQs are

$$\text{scalar case: } R = \begin{pmatrix} R_+ \\ R_- \end{pmatrix} \sim (3, 2, 7/6)$$

$$\tilde{R} = \begin{pmatrix} \tilde{R}_+ \\ \tilde{R}_- \end{pmatrix} \sim (3, 2, 1/6), \tag{20a}$$

$$\text{vector case: } V = \begin{pmatrix} V_+ \\ V_- \end{pmatrix} \sim (\bar{3}, 2, 5/6)$$

$$\tilde{V} = \begin{pmatrix} \tilde{V}_+ \\ \tilde{V}_- \end{pmatrix} \sim (\bar{3}, 2, -1/6), \tag{20b}$$

where $\pm$ states are weak isospin components, and the $SU(3)_C \times SU(2)_L \times U(1)_Y$ quantum numbers are given in parentheses.[6] In both cases, the most general renormalizable interactions to first generation fermions (including ν_{eR}) are

$$\text{scalar: } \mathcal{L}_{\rm int} = h_L \bar{u}_R L_L^T \epsilon R + h_R \bar{Q}_L e_R R + \tilde{h}_L \bar{d}_R L_L^T \epsilon \tilde{R} + \tilde{h}_R \bar{Q}_L \nu_{eR} \tilde{R} + \text{h.c.}, \tag{21a}$$

$$\text{vector: } \mathcal{L}_{\rm int} = g_L \bar{d}_R^c \gamma^\mu L_L^T \epsilon V_\mu + g_R \bar{Q}_L^c \gamma^\mu e_R \epsilon V_\mu + \tilde{g}_L \bar{u}_R^c \gamma^\mu L_L^T \epsilon \tilde{V}_\mu + \tilde{g}_R \bar{Q}_L^c \gamma^\mu \nu_{eR} \epsilon \tilde{V}_\mu + \text{h.c.}, \tag{21b}$$

with quark and lepton doublets $Q_L = (u_L, d_L)$ and $L_L = (\nu_{eL}, e_L)$. Here, $h_{L,R}$, $\tilde{h}_{L,R}$, $g_{L,R}$, $\tilde{g}_{L,R}$ are couplings (with L, R denoting lepton chirality). The presence of both L-, R-type couplings will lead to lepton universality violation in $\pi^+ \to e^+ \nu$; to avoid this constraint, we set R-type couplings to zero [32]. The relevant mass terms are

$$\text{scalar: } -\mathcal{L}_{\rm mass} = m_R^2 R^\dagger R + m_{\tilde{R}}^2 \tilde{R}^\dagger \tilde{R} + (\lambda_R (R^\dagger H)(\tilde{R} H) + \text{h.c.}), \tag{22a}$$

$$\text{vector: } \mathcal{L}_{\rm mass} = m_V^2 V_\mu^\dagger V^\mu + m_{\tilde{V}}^2 \tilde{V}_\mu^\dagger \tilde{V}^\mu + (\lambda_V (V_\mu^\dagger H)(\tilde{V}^\mu H) + \text{h.c.}). \tag{22b}$$

Through electroweak symmetry breaking, the quartic interactions (with couplings $\lambda_{R,V}$) give rise to R_--$\tilde{R}_+$ mixing and V_--$\tilde{V}_+$ mixing by generating off-diagonal mass terms proportional to $\lambda_{R,V} v^2$, where $v \equiv \langle H^0 \rangle$. Diagonalizing the R_--$\tilde{R}_+$ and V_--$\tilde{V}_+$ mass matrices, we can express the mass eigenstates, denoted $\mathcal{R}_{1,2}$ and $\mathcal{V}_{1,2}$, as

$$\text{scalar: } \mathcal{R}_1 \equiv \cos\theta_R R_- + \sin\theta_R e^{i\phi_R} \tilde{R}_+, \quad \mathcal{R}_2 \equiv \cos\theta_R \tilde{R}_+ - \sin\theta_R e^{-i\phi_R} R_-, \tag{23a}$$

$$\text{vector: } \mathcal{V}_1 \equiv \cos\theta_V V_- + \sin\theta_V e^{i\phi_V} \tilde{V}_+, \quad \mathcal{V}_2 \equiv \cos\theta_V \tilde{V}_+ - \sin\theta_V e^{-i\phi} V_- \tag{23b}$$

with mixing angles $\theta_{R,V}$ and mass eigenvalues given by

$$\text{scalar: } \tan 2\theta_R = \frac{2|\lambda_R| v^2}{m_R^2 - m_{\tilde{R}}^2}, \quad m_{\mathcal{R}_{1,2}}^2 = \frac{1}{2}(m_R^2 + m_{\tilde{R}}^2 \pm \sqrt{(m_R^2 - m_{\tilde{R}}^2)^2 + 4|\lambda_R|^2 v^4}), \tag{24a}$$

$$\text{vector: } \tan 2\theta_V = \frac{2|\lambda_V| v^2}{m_V^2 - m_{\tilde{V}}^2}, \quad m_{\mathcal{V}_{1,2}}^2 = \frac{1}{2}(m_V^2 + m_{\tilde{V}}^2 \pm \sqrt{(m_V^2 - m_{\tilde{V}}^2)^2 + 4|\lambda_V|^2 v^4}), \tag{24b}$$

[6] We follow the notation of Ref. [32] for LQ states, except we omit an additional subscript identifying the $SU(2)_L$ representation.

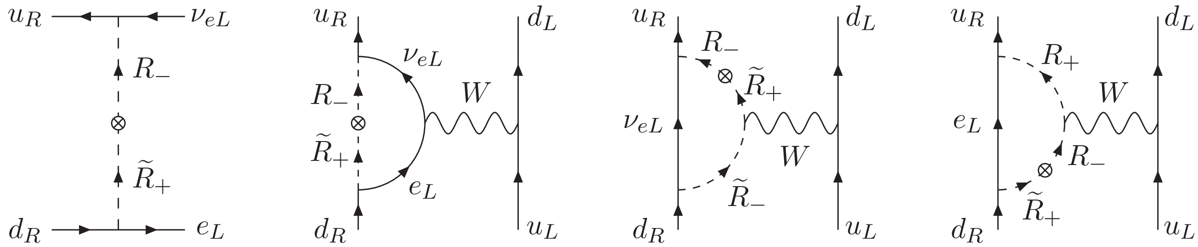

FIG. 3. Scalar LQ case: Tree-level exchange generates β decay amplitude a_{LR}^V (left), while $\mathcal{O}_{LR}$ is generated by one-loop vertex corrections (right), contributing to EDMs d_n, $d_{\rm Hg}$, d_D. Diagrams are shown in weak-eigenstate LQ basis to illustrate that the same CP-violating phases from LQ mixing (denoted ⊗) and couplings enter both D_t and EDMs.

and phases $\phi_{R,V} = \arg(\lambda_{R,V})$, defined such that $m^2_{\mathcal{R}_1,\mathcal{V}_1} < m^2_{\mathcal{R}_2,\mathcal{V}_2}$. The remaining (unmixed) LQ states R_+, V_+ and $\tilde{R}_-$, $\tilde{V}_-$ have masses $m_{R,V}$ and $m_{\tilde{R},\tilde{V}}$, respectively.

For β decay, this model gives $D_t = \kappa\,{\rm Im}(a_{LR}^V)$, where [6]

$$\text{scalar case:}\, a_{LR}^V = \frac{h_L \tilde{h}_L^* \sin 2\theta_R e^{i\phi_R}}{8\sqrt{2}G_F}\left(\frac{1}{m^2_{\mathcal{R}_1}} - \frac{1}{m^2_{\mathcal{R}_2}}\right), \quad (25a)$$

$$\text{vector case:}\, a_{LR}^V = \frac{g_L \tilde{g}_L^* \sin 2\theta_V e^{i\phi_V}}{4\sqrt{2}G_F}\left(\frac{1}{m^2_{\mathcal{V}_1}} - \frac{1}{m^2_{\mathcal{V}_2}}\right). \quad (25b)$$

The relevant Feynman diagrams are shown in Figs. 3 and 4.

Next, we consider implications for EDMs. Radiative corrections involving the W boson, shown in Figs. 3 and 4, generate the CP-odd four-quark operator $\mathcal{O}_{LR}$ given in Eq. (7) which contributes to d_n and $d_{\rm Hg}$. The resulting coefficient k_{LR} is proportional to the same CP-violating phases in Eq. (25) entering D. For each case, we find

$$\begin{aligned} \text{scalar: } k_{LR} &= \frac{8G_F m^2_{\mathcal{R}_1}}{\sqrt{2}(4\pi)^2} F_R\, {\rm Im}(a_{LR}^V), \\ \text{vector: } k_{LR} &= -\frac{8G_F m^2_{\mathcal{V}_1}}{\sqrt{2}(4\pi)^2} F_V\, {\rm Im}(a_{LR}^V). \end{aligned} \quad (26)$$

The loop functions $F_{R,V}$ are given by

$$F_R \equiv \frac{m^2_{\mathcal{R}_2}}{2(m^2_{\mathcal{R}_1} - m^2_{\mathcal{R}_2})}(f(m^2_{\mathcal{R}_1}, m^2_{\mathcal{R}_2}, m^2_R) + f(m^2_{\mathcal{R}_1}, m^2_{\mathcal{R}_2}, m^2_{\tilde{R}}) + f(m^2_{\mathcal{R}_1}, m^2_{\mathcal{R}_2}, 0)), \quad (27a)$$

$$F_V \equiv \frac{m^2_{\mathcal{V}_2}}{2(m^2_{\mathcal{V}_1} - m^2_{\mathcal{V}_2})}(3f(m^2_{\mathcal{V}_1}, m^2_{\mathcal{V}_2}, m^2_V) + 3f(m^2_{\mathcal{V}_1}, m^2_{\mathcal{V}_2}, m^2_{\tilde{V}}) - f(m^2_{\mathcal{V}_1}, m^2_{\mathcal{V}_2}, 0)), \quad (27b)$$

where

$$f(m_1^2, m_2^2, m_3^2) \equiv \frac{m_1^2 m_2^2 \log(m_1^2/m_2^2) + m_2^2 m_3^2 \log(m_2^2/m_3^2) + m_3^2 m_1^2 \log(m_3^2/m_1^2)}{(m_1^2 - m_3^2)(m_2^2 - m_3^2)}. \quad (28)$$

Defined in this way, we have $F_{R,V} \geq 1$, with equality in the limit $m^2_{R,V} = m^2_{\tilde{R},\tilde{V}} \gg \lambda_{R,V} v^2$. Equation (26) provides the leading contributions to EDMs from CP violation entering D; there is no suppression by light quark masses $m^2_{u,d}$ as previously argued [6]. Using the results of Sec. II, we have (both cases give the same numerical factors)

$$|d_n| > 4 \times 10^{-21}\, e\,{\rm cm} \times |D_t/\kappa| \left(\frac{m_{LQ}}{300\ {\rm GeV}}\right)^2, \quad (29a)$$

$$|d_{\rm Hg}| > 3 \times 10^{-25}\, e\,{\rm cm} \times |D_t/\kappa| \left(\frac{m_{LQ}}{300\ {\rm GeV}}\right)^2, \quad (29b)$$

$$|d_D| > 1.7 \times 10^{-21}\, e\,{\rm cm} \times |D_t/\kappa| \left(\frac{m_{LQ}}{300\ {\rm GeV}}\right)^2, \quad (29c)$$

where $m_{LQ} = m_{\mathcal{R}_1}$ ($m_{\mathcal{V}_1}$) corresponds to the lightest LQ state entering β decay for the scalar (vector) LQ case. ($\kappa \approx 0.87, -1.03$ for n, ^{19}Ne, respectively.)

Recent searches at hadron colliders [79–82] provide constraints on the mass of the lightest LQ ($\mathcal{R}_1$, $\mathcal{V}_1$) involved in β decay. These bounds depend on the branching ratio $\beta_e \equiv {\rm BR}(LQ \to je) = 1 - {\rm BR}(LQ \to j\nu) = \sin^2\theta_{R,V}$, where j is a jet. For the scalar case, the strongest limits have been obtained at the Large Hadron Collider by combining $jjee$ and $jje\nu$ channels [79,80]

$$m_{\mathcal{R}_1} > \begin{cases} 340\ {\rm GeV} & ({\rm CMS}) \\ 319\ {\rm GeV} & ({\rm ATLAS}) \end{cases} \qquad (\beta_e > 0.5) \quad (30)$$

with stronger limits (384 and 376 GeV, respectively) for $\beta_e \to 1$. Additionally, recent ATLAS searches for jets with

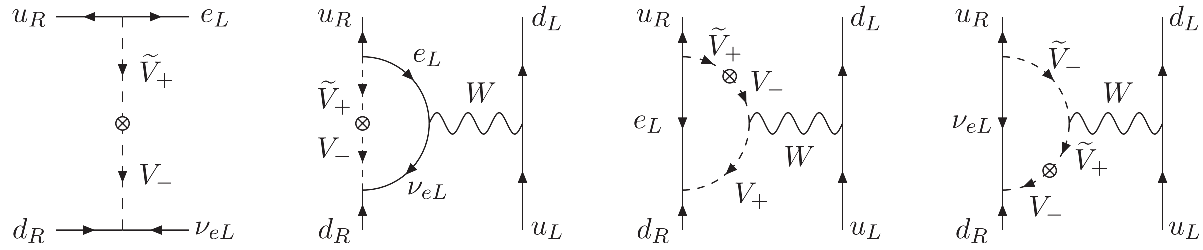

FIG. 4. Vector LQ case: Tree-level exchange generates β decay amplitude a_{LR}^V (left), while $\mathcal{O}_{LR}$ is generated by one-loop vertex corrections (right), contributing to EDMs d_n, $d_{\rm Hg}$, d_D. Diagrams are shown in weak-eigenstate LQ basis to illustrate that the same *CP*-violating phases from LQ mixing (denoted $\otimes$) and couplings enter both D_t and EDMs.

missing energy from squark pair production, within a simplified SUSY context [83], apply to $jj\nu\nu$ final states from $\mathcal{R}_1$ pair production. To translate the SUSY model into our framework, one must rescale the SUSY cross section by a factor $(1-\beta_e)^2/4$ and take the gluino to be massive.[7] The resulting limits are $m_{\mathcal{R}_1} \gtrsim 500$ GeV, for $\beta_e < 0.5$, with stronger bounds in the limit $\beta_e \to 0$. In the vector case, the D0 Collaboration has obtained [82,84]

$$m_{\mathcal{V}_1} > \begin{cases} 302\ \text{GeV} & (jjee + jje\nu,\, \beta_e > 0.1) \\ 144\ \text{GeV} & (jj\nu\nu,\, \beta_e < 0.1) \end{cases} \tag{31}$$

with stronger bounds for $\beta_e \to 1$ or for different choices of anomalous gluon-LQ couplings considered therein. Within the context of our model, for $\beta_e = \sin^2\theta_V < 0.1$, the lightest vector LQ $\mathcal{V}_1 \approx V_-$ is nearly degenerate with V_+ ($m_{\mathcal{V}_1} \approx m_V$). Since $\text{BR}(V_+ \to je) = 1$, we have $m_V > 367$ [82], and therefore $m_{\mathcal{V}_1}$ is constrained indirectly to be much heavier than 144 GeV. Additional constraints have been obtained by the H1 Collaboration at HERA.[8] These limits depend on the LQ-fermion couplings, and provide stronger bounds than those from hadron colliders if the relevant e-q-LQ coupling is larger than $\sim$few $\times 10^{-1}$ [85].

Atomic parity violation experiments [86] are sensitive to e-q contact interactions and provide important constraints on LQ models. For the case of cesium (^{133}Cs), the measured weak charge $Q_w(\text{Cs}) = -73.20 \pm 0.35$ is in good agreement with the SM prediction $Q_w^{\rm SM}(\text{Cs}) = -73.15 \pm 0.02$ [40]. Using results from Ref. [87], the LQ contribution to $Q_w(\text{Cs})$ for the two scenarios considered here is given by

$$Q_w^{\rm LQ}(\text{Cs}) \simeq \begin{cases} 5.7\ \text{TeV}^2|h_L|^2/m_R^2 + 6.4\ \text{TeV}^2|\tilde{h}_L|^2/m_{\tilde{R}}^2 & \text{scalar case} \\ -11.4\ \text{TeV}^2|g_L|^2/m_V^2 - 12.8\ \text{TeV}^2|\tilde{g}_L|^2/m_{\tilde{V}}^2 & \text{vector case,} \end{cases} \tag{32}$$

neglecting additional $\mathcal{O}(\lambda v_{R,V}^2/m_{\rm LQ}^2)$ corrections. Requiring $-0.75 < Q_w^{\rm LQ}(\text{Cs}) < 0.65$ (at 95% CL), we have

$$\text{scalar: } m_R > 3.0|h_L|\text{TeV}, \qquad m_{\tilde{R}} > 3.1|\tilde{h}_L|\text{TeV}, \tag{33a}$$

$$\text{vector: } m_V > 3.9|g_L|\text{TeV}, \qquad m_{\tilde{V}} > 4.1|\tilde{g}_L|\text{TeV}. \tag{33b}$$

These constraints are stronger than the aforementioned collider bounds for LQ couplings greater than $\mathcal{O}(0.1)$, although a cancellation is possible if both scalars and vectors are present.

Given the current limit $|d_n| < 2.9 \times 10^{-26}\, e\,\text{cm}$ [22], and conservatively taking $m_{LQ} > 300$ GeV, we conclude that $|D_t/\kappa| < 7 \times 10^{-6}$. *CP* violation in LQ models cannot saturate present experimental sensitivities in D_n—unless the hadronic uncertainties associated with the d_n computation of Ref. [35] are larger by an order of magnitude, or unless there is a cancellation with other *CP*-odd phases contributing to d_n to $\sim$10% (or a combination thereof). On the other hand, the mercury EDM does not strongly constrain D_t in this model, especially in light of its large hadronic uncertainties, although this situation may change with future refinements in the nuclear computations.

B. LQ scenarios with right-handed neutrinos

LQ models can contribute to D_t through the interference between two new physics amplitudes involving right-handed neutrinos. The relevant Feynman diagrams are shown in Fig. 5. To begin, we consider the model of the preceding section involving scalars R, $\tilde{R}$ and vectors V, $\tilde{V}$, with mixing defined in Eqs. (22)–(24) and couplings to SM fermions given in Eq. (21). Here, we set to zero L-type couplings in Eq. (21) and keep only R-type ones. For each case, the amplitude a_{RL}^V is

[7]The factor 4 counts the number of first and second generation squarks in the simplified SUSY model considered in Ref. [83].

[8]To translate between the notation used here and that in Ref. [85], we note $S_{1/2}^L \equiv R$, $\tilde{S}_{1/2}^L \equiv \tilde{R}$, $V_{1/2}^L \equiv V$, and $\tilde{V}_{1/2}^L \equiv \tilde{V}$.

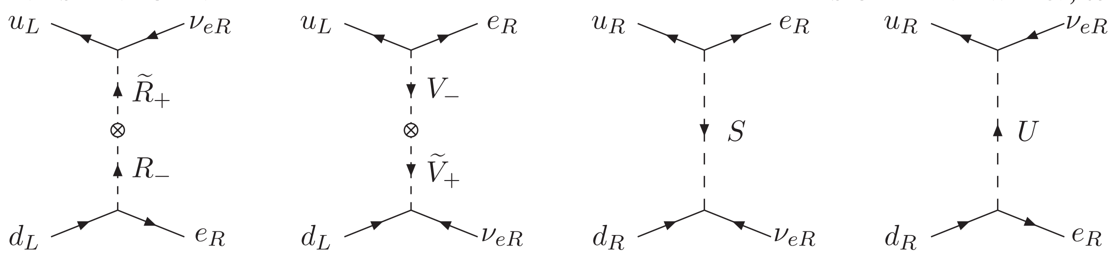

FIG. 5. $D_t = k\,\mathrm{Im}(a^V_{RL} a^{V*}_{RR})$ is generated by LQ couplings involving right-handed neutrinos, with a^V_{RL} from R-, $\tilde{R}$- or V-, $\tilde{V}$-exchange (mixing denoted $\otimes$), and a^V_{RR} from S- or U-exchange.

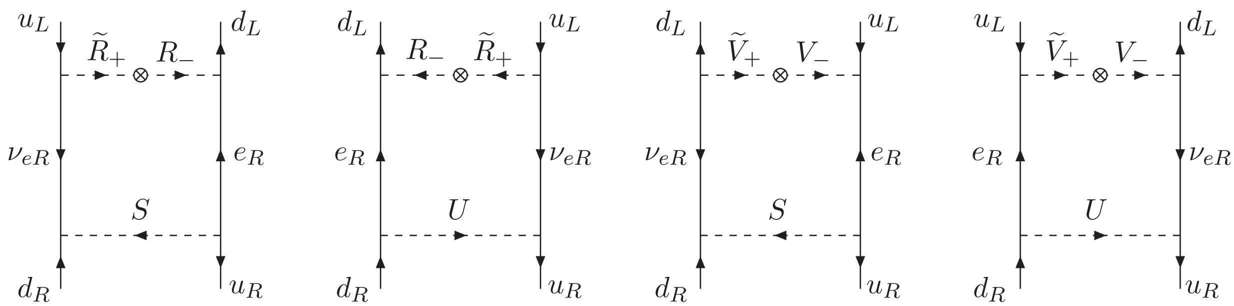

FIG. 6. LQ contributions to D_t generate radiatively CP-odd operator $\mathcal{O}_{LR}$ contributing to EDMs.

$$\mathcal{R}_{1,2}\text{-exchange: } a^V_{RL} = -\frac{\tilde{h}_R h^*_R \sin 2\theta_R e^{-i\phi_R}}{8\sqrt{2}G_F} \times \left(\frac{1}{m^2_{\mathcal{R}_1}} - \frac{1}{m^2_{\mathcal{R}_2}}\right), \tag{34a}$$

$$\mathcal{V}_{1,2}\text{-exchange: } a^V_{RL} = \frac{\tilde{g}_R g^*_R \sin 2\theta_V e^{-i\phi_V}}{4\sqrt{2}G_F} \times \left(\frac{1}{m^2_{\mathcal{V}_1}} - \frac{1}{m^2_{\mathcal{V}_2}}\right). \tag{34b}$$

In order to generate a^V_{RR}, we introduce two additional LQ states S and U, with quantum numbers

$$\text{scalar LQ}: S \sim (\bar{3}, 1, 1/3), \qquad \text{vector LQ}: U \sim (3, 1, 2/3) \tag{35}$$

and quark-lepton couplings

$$\mathcal{L}_{\text{int}} = (g_S \bar{u}^c_R e_R + g'_S \bar{d}^c_R \nu_{eR})S + (h_U \bar{d}_R \gamma^\mu e_R + h'_U \bar{u}_R \gamma^\mu \nu_{eR})U_\mu + \text{h.c.} \tag{36}$$

Through tree-level exchange, these states generate

$$S\text{-exchange: } a^V_{RR} = \frac{g'_U g^*_U}{4\sqrt{2}G_F m^2_S}, \qquad U\text{-exchange: } a^V_{RR} = -\frac{h'_U h^*_U}{2\sqrt{2}G_F m^2_U}. \tag{37}$$

There are four possible contributions to $D_t = \kappa\,\mathrm{Im}(a^V_{RL} a^{V*}_{RR})$ depending on which of the combinations

$$(\mathcal{R}_1, \mathcal{R}_2, S), \qquad (\mathcal{R}_1, \mathcal{R}_2, U), \qquad (\mathcal{V}_1, \mathcal{V}_2, S), \qquad (\mathcal{V}_1, \mathcal{V}_2, U) \tag{38}$$

we consider contributing to a^V_{RL} and a^V_{RR}.

Next, we consider each of these combinations separately and compute the resulting EDM induced by the CP-odd four-quark operator in Eq. (7). There are four possible contributions, shown in Fig. 6, and they all give nearly identical results:

$$|k_{LR}| = \frac{\sqrt{2}G_F m^2_{LQ}}{(4\pi)^2} |\mathrm{Im}(a^V_{RL} a^{V*}_{RR})| \hat{f}(m^2_1, m^2_2, m^2_3). \tag{39}$$

The loop function is

$$\hat{f}(m^2_1, m^2_2, m^2_3) \equiv \frac{2m^2_1 m^2_2 m^2_3 (m^2_1 \log(m^2_2/m^2_3) + m^2_2 \log(m^2_3/m^2_1) + m^2_3 \log(m^2_1/m^2_2))}{m^2_{LQ}(m^2_2 - m^2_1)(m^2_3 - m^2_2)(m^2_1 - m^2_3)}, \tag{40}$$

where, for each case in Eq. (38), $m^2_{1,2,3}$ corresponds to the masses of the three states, with m^2_{LQ} being the smallest of the three. Defined in this way, we have $\hat{f} \geq 1$, with equality if all states are degenerate.

Assuming that one *CP*-violating phase is dominant in D_t, the resulting EDMs arising from that phase are

$$|d_n| > 9 \times 10^{-22}\, e\,\mathrm{cm} \times |D_t/\kappa| \left(\frac{m_{LQ}}{300\ \mathrm{GeV}}\right)^2, \quad (41a)$$

$$|d_{\mathrm{Hg}}| > 7 \times 10^{-26}\, e\,\mathrm{cm} \times |D_t/\kappa| \left(\frac{m_{LQ}}{300\ \mathrm{GeV}}\right)^2, \quad (41b)$$

$$|d_D| > 4 \times 10^{-22}\, e\,\mathrm{cm} \times |D_t/\kappa| \left(\frac{m_{LQ}}{300\ \mathrm{GeV}}\right)^2. \quad (41c)$$

Comparing Eqs. (26) and (39), we find that D_t from LQ scenarios involving right-handed neutrinos is less constrained by EDMs by a factor 4 compared those involving left-handed neutrinos (for fixed m_{LQ}). For $m_{LQ} > 300$ GeV, the neutron EDM bound implies $|D_t/\kappa| \lesssim 3 \times 10^{-5}$.

Constraints on scalar and vector LQ masses from pair production at hadron colliders are the same as in Eqs. (30) and (31). However, in the limit $h_U \ll h'_U$, the vector U decays primarily via $U \to j\nu$ and is subject to the relatively weaker mass bound $m_U > 144$ GeV [84]. Significantly stronger bounds are provided by the H1 Collaboration for $\beta_e(U) \approx 0$ [85], which depend on the U-e-d coupling h_U:

$$m_U \gtrsim \begin{cases} 250\ \mathrm{GeV} & (h_U = 0.03) \\ 300\ \mathrm{GeV} & (h_U = 0.06) \\ 1\ \mathrm{TeV} & (h_U = 0.3). \end{cases} \quad (42)$$

Although suppressing h_U weakens the bound on m_U, the contribution to $D_t(\propto h_U/m_U^2)$ is also suppressed. Assuming $h_U \gtrsim \mathcal{O}(0.06)$ (to avoid too much additional suppression in D_t) we take $m_U \gtrsim 300$ GeV.[9] In addition, the weak charge of cesium is given by (using results from Ref. [87])

$$Q_w^{\mathrm{LQ}}(\mathrm{Cs}) \simeq 5.7\ \mathrm{TeV}^2 \frac{|g_S|^2}{m_S^2} - 24.2\ \mathrm{TeV}^2 \frac{|h_R|^2}{m_R^2} + 12.8\ \mathrm{TeV}^2 \frac{|h_U|^2}{m_U^2} + 24.2\ \mathrm{TeV}^2 \frac{|g_R|^2}{m_V^2}. \quad (43)$$

Considering each state individually, we have

$$\begin{aligned} m_S &> 3.0|g_S|\mathrm{TeV}, \qquad m_R > 5.7|h_R|\mathrm{TeV}, \\ m_V &> 6.1|g_R|\mathrm{TeV}, \qquad m_U > 4.4|h_U|\mathrm{TeV}, \end{aligned} \quad (44)$$

although clearly these bounds are weakened in the presence of cancellations.

[9] It seems plausible that the best trade-off between small m_U and small h_U occurs for $m_U \sim 300$ GeV, corresponding to the center-of-mass energy $\sqrt{s} = 319$ GeV at HERA. For $m_U < \sqrt{s}$, on-shell LQ production dominates, allowing for relatively stronger constraints on h_U; for $m_U > \sqrt{s}$, only off-shell production is allowed, and the constraints are weaker [85]. A more precise analysis is beyond the scope of this work.

V. CONCLUSIONS

The emiT Collaboration has measured $D_n = (-1.0 \pm 2.1) \times 10^{-4}$ [15], consistent with the SM prediction dominated by $\mathcal{O}(10^{-5})$ final state effects. Here, we studied several new physics scenarios beyond the SM and showed that the current neutron EDM measurement $|d_n| < 2.9 \times 10^{-26}\, e\,\mathrm{cm}$ provided in all cases stronger bounds on D.

(i) $|D_t/\kappa| < 3 \times 10^{-7}$ in left-right symmetric models, exotic fermion models, and the *R*-parity violating MSSM. EDM bounds on this class of models, given in Eq. (16), can be understood in an otherwise model-independent operator framework through a coupling of the W boson to the right-handed quark charge current $\bar{u}_R \gamma^\mu d_R$.

(ii) $|D_t/\kappa| < 3 \times 10^{-5}$ (7×10^{-6}) in leptoquark models with (without) light right-handed neutrinos. Moreover, EDM constraints will become more severe if collider bounds on leptoquark masses are improved, as shown in Eqs. (29) and (41).

We recall that $\kappa \approx 0.87$ (for the neutron) is defined in Eq. (6), and D_t denotes the contribution to D from fundamental T violation (as opposed to final state effects). Analogous constraints from the mercury EDM bound are weaker by an order of magnitude (with large uncertainties), although the situation may change with future improvements in the nuclear computations. A future constraint on the deuteron EDM of $|d_D| \lesssim 10^{-28}\, e\,\mathrm{cm}$ would improve all aforementioned bounds on D_t by 2 orders of magnitude. These bounds can in principle be evaded by fine-tuned cancellations with other *CP*-odd phases contributing to EDMs, but not to D_t.

Even though D is not as sensitive as EDMs to *CP* violation beyond the SM, clearly it worthwhile to push D measurements to greater sensitivities. Since any single EDM measurement has little model discriminating power, it is desirable to consider as many observables as possible—especially if a nonzero EDM were measured. In this case, D could play an important role in untangling the nature of *CP* violation and potentially shedding light on the origin of matter in the Universe.

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