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Right-handed neutrinos and T-violating, P-conserving interactions

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ABSTRACT

We show that experimental probes of the P-conserving, T-violating triple correlation in polarized neutron or nuclear β-decay provide a unique probe of possible T-violation at the TeV scale in the presence of right-handed neutrinos. In contrast to other possible sources of semileptonic T-violation involving only left-handed neutrinos, those involving right-handed neutrinos are relatively unconstrained by present limits on the permanent electric dipole moments of the electron, neutral atoms, and the neutron. On the other hand, LHC results for $pp \to e+$ missing transverse energy imply that an order of magnitude of improvement in D-coefficient sensitivity would be needed for discovery. Finally, we discuss the interplay with the scale of neutrino mass and naturalness considerations.

1. Introduction

The search for time-reversal violation (TV) has long been a subject of considerable experimental and theoretical interest. It is partially motivated by the need for CP-violation beyond that encoded in the Standard Model (SM) Cabibbo–Kobayashi–Maskawa (CKM) matrix to explain the cosmic baryon asymmetry [1]. Assuming CPT is a good symmetry of nature, searches for TV provide a probe of this possible CP-violation. Experimentally, the parity (P)- and T-violating (PVTV) sector is being probed with great sensitivity through electric dipole moment (EDM) searches, with the three most stringent limits having been obtained for the ^{199}Hg atom [2], the electron (extracted from the ThO molecule) [3], and the neutron [4,5]. On the other hand, the P-conserving and T-violating (PCTV) sector (equivalent to C- and CP-violation assuming CPT) has received considerably less attention. Experimental efforts in the sector include measurement of anomalous η-decay channels such as $\eta \to 2\pi^0\gamma$, $3\pi^0\gamma$, 3γ [6] and the D-coefficient in the β-decay of polarized neutrons [7] and ^{19}Ne [8]. These processes are sensitive probes of "new physics" because the Standard Model (SM) contributions are usually small [9,10]. There is a SM final-state interaction that could mimic a non-zero D-coefficient in β-decay at order 10^{-5} for neutron [11] and 10^{-4} for ^{19}Ne [8] but the application of heavy baryon effective field theory allows a precise computation of this contribution (up to 1% accuracy in the case of neutron [12]).

Theoretically, the effect of PCTV physics due to beyond Standard Model (BSM) interactions can be studied in a model-independent way using effective field theory (EFT). In this approach, one has integrated out the BSM heavy degrees of freedom (DOF). In this context, it was observed in Ref. [13] that any EDM limits imply severe bounds on PCTV observables since a PCTV interaction in the presence of P-violating SM radiative corrections will induce an EDM. While special exceptions to this argument may occur [14,15], the question remains as to the prospective impact of, and motivation for, improved probes of flavor-conserving PCTV observables. Recently, the authors of Ref. [16] addressed this question in the EFT context, studying the contribution of the "left-right four fermion" (LR4F) operator to the D-coefficient of the neutron β-decay (defined below). They find that the neutron EDM sets an indirect bound on the D-coefficient that is three orders of magnitude more stringent than its direct experimental bound.

In this paper, we observe that there exists a set of dimension-six four-fermion operators involving right-handed neutrinos that (a) contribute to the D-coefficient and (b) are relatively unconstrained by EDM limits. Because the SM charge changing weak interaction involves purely left-handed leptons, this contribution to neutron decay does not interfere linearly with the SM contribution, resulting in a quadratic, rather than linear, dependence on the operator Wilson coefficients. Nonetheless, present limits on D probe the TeV mass scale. We also show that while a subset of these

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operators generate hadronic EDMs, their effects are suppressed by loop factors as well as Λ_χ/v where $\Lambda_\chi \sim 1$ GeV is the chiral symmetry breaking scale and $v = 246$ GeV is the Higgs vacuum expectation value (VEV). The resulting neutron EDM sensitivity to Λ is also at the TeV scale and does not preclude a non-zero result in a next generation D-coefficient probe.

Interestingly, indirect constraints from T-conserving observables may be more severe. These observables include Large Hadron Collider (LHC) results for the process $pp \to e + X + MET$ (missing transverse energy) and neutrino mass. The latter constraints also rely on naturalness considerations, a somewhat subjective criteria. The former imply that an order of magnitude improvement in D-coefficient sensitivity would be required in order to discover evidence for PCTV right-handed neutrino interactions.

Our analysis leading to these conclusions is organized as follows. In Sec. 2 we introduce the relevant set of dimension-6 operators and discuss the experimental D-coefficient constraint on their Wilson coefficients. We then compare this constraint to those implied by LHC data, hadronic EDMs as well as neutrino mass and naturalness considerations. For comparison, we perform in Sec. 3 a similar analysis of other dimension-6 operators that do not involve right-handed neutrinos. We show that any attempt to evade current EDM constraints and yet keep the size of the D-coefficient experimentally accessible would involve fine tuning at the 10^{-11} level. We conclude in Sec. 4.

2. Dimension-six operators with right-handed neutrinos

The PCTV observable of interest in β-decay involves a triple correlation of the spin of the decaying particle and the momenta of the outgoing leptons that enters the differential β-decay rate. In what follows, we focus on neutron, for which the experimental bound on the PCTV triple correlation is the most stringent. However, the discussion below can be easily generalized to other cases. The differential decay rate for a polarized neutron is given by:

$$\frac{d\Gamma}{dE_e d\Omega_e d\Omega_\nu} = \frac{G_F^2 V_{\rm ud}^2}{(2\pi)^5}(g_V^2 + 3g_A^2)|\vec{p}_e| E_e E_\nu^2 \times \left[1 + a\frac{\vec{p}_e \cdot \vec{p}_\nu}{E_e E_\nu} + \hat{s}\cdot\left(A\frac{\vec{p}_e}{E_e} + B\frac{\vec{p}_\nu}{E_\nu} + D\frac{\vec{p}_e \times \vec{p}_\nu}{E_e E_\nu}\right)\right] \tag{1}$$

where $\hat{s}$ is the unit polarization vector of the neutron; $\vec{p}_e$ and $\vec{p}_\nu$ are the electron and anti-neutrino momenta, respectively, with corresponding energies $E_{e(\nu)}$; G_F is the Fermi constant; and $V_{\rm ud}$ is the first generation element of the Cabibbo–Kobayashi–Maskawa (CKM) matrix. The most stringent experimental limit on the D-coefficient is given by $D = (-0.96 \pm 1.89 \pm 1.01) \times 10^{-4}$ [7] which translates into an upper bound of $|D| < 4 \times 10^{-4}$ at 90% CL [17].

Theoretically, a non-vanishing contribution can be generated by the interference of amplitudes involving a small set of dimension $d = 6$ effective operators. Considering only first generation SM fermions and requiring SM gauge invariance, one finds a limited set of such $d = 6$ TV operators (see Ref. [18] for a complete list of gauge-invariant $d = 6$ operators involving SM fields). As we discuss in Section 3, EDM constraints imply severe bounds on the contribution of these operators to D. Extending the set of fields to include right-handed (RH) neutrinos, one finds an additional set of four-fermion operators that contribute to D at tree-level and that are relatively immune to EDM constraints [19]:

$$\hat{O}_1 = \frac{c_1(\mu)}{\Lambda^2}\bar{L}^i \nu_R \bar{u}_R Q^i + \text{h.c.}$$

$$\hat{O}_2 = \frac{c_2(\mu)}{\Lambda^2}\varepsilon^{ij}\bar{L}^i \nu_R \bar{Q}^j d_R + \text{h.c.}$$

$$\hat{O}_3 = \frac{c_3(\mu)}{\Lambda^2}\varepsilon^{ij}\bar{L}^i \sigma^{\mu\nu}\nu_R \bar{Q}^j \sigma_{\mu\nu} d_R + \text{h.c.} \tag{2}$$

where Λ is the BSM mass scale and μ is the renormalization scale. These operators are analogous to the semi-leptonic four-fermion operators of type $(\bar{L}R)(\bar{L}R)$ and $(\bar{L}R)(\bar{R}L)$ in Ref. [18]. Also notice that the Wilson coefficients $c_1 - c_3$ are functions of the renormalization scale μ, which as to be taken as the hadronic scale when we discuss the bounds of the Wilson coefficients from low-energy experiments.

It is straightforward to compute the contributions of $\hat{O}_{1-3}$ to the D-coefficient. The dominant affect is quadratic in the c_i/Λ^2, as the linear interference term is suppressed by the neutrino mass. Following Ref. [20], we obtain, to leading non-trivial order in $\{c_i\}$,

$$D = -\frac{g_S g_T}{\Lambda^4}\frac{1}{G_F^2 V_{ud}^2(g_V^2 + 3g_A^2)}\text{Im}[(c_1 - c_2)c_3^*]\Big|_{\mu=\mu_h} \tag{3}$$

where g_S and g_T are the nucleon scalar and tensor charges, respectively, and $\mu_h \approx 1$ GeV is the hadronic scale.

Even though one pays a price in BSM sensitivity owing to a quadratic rather than linear dependence on the c_i/Λ^2, the gain achieved by avoiding EDM constraints is considerable (see Section 3). Taking the updated lattice calculation of $g_S = 0.97(12)(6)$ and $g_T = 0.987(51)(20)$ [21], we obtain:

$$\left|\frac{\text{Im}[(c_1 - c_2)c_3^*]}{\Lambda^4}\right|\Big|_{\mu=\mu_h} < 3 \times 10^{-1}\ \text{TeV}^{-4} \tag{4}$$

If we take $\{c_i\} \sim c$ without distinguishing the real and imaginary part, then this inequality implies that existing D-coefficient studies probe BSM T-violating interactions with RH neutrinos with a sensitivity of $(v/\Lambda)^2 c \sim 3 \times 10^{-2}$ at $\mu = \mu_h$. One could estimate the sensitivity to Λ by assuming that $c_i \sim 1$ at $\mu \approx \Lambda$. QCD running in $\overline{\text{MS}}$ scheme gives $c_{1,2}(\Lambda) \approx 0.56 c_{1,2}(\mu_h)$ and $c_3(\Lambda) \approx 1.2 c_3(\mu_h)$ for $\Lambda > m_W$ where m_W is the mass of the W-boson (see, e.g. Ref. [22]). Then, the current bound implies $\Lambda \gtrsim 1$ TeV.

The operators $\hat{O}_{1-3}$ can induce hadronic EDMs at one-loop order, but their contributions also scale quadratically with the c_i/Λ^2. In particular, the combination of $\hat{O}_1$ and $\hat{O}_2$ may induce the CP-odd four-quark operator [23]

$$\frac{C^{(1)}_{quqd}(\mu)}{\Lambda^2}\varepsilon^{ij}\bar{Q}^i u_R \bar{Q}^j d_R + \text{h.c.} \tag{5}$$

via the one-loop graph of Fig. 2a. Contributions from loop momenta $k < \Lambda$ vanish, as seen explicitly in dimensional regularization (DR), because the amplitude involves a quadratically-divergent integral with massless propagators and because it is infrared finite. Non-vanishing contributions result from $k \gtrsim \Lambda$ that are associated with matching onto the *a priori* unknown ultraviolet complete theory that generates the non-vanishing c_i. Estimating these matching contributions using a cut-off regulator [24] yields[1]

$$\frac{C^{(1)}_{quqd}(\Lambda)}{\Lambda^2} \sim \frac{\Lambda^2}{16\pi^2}\frac{c_1^* c_2}{\Lambda^4}\Big|_{\mu=\Lambda} = \frac{c_1^* c_2}{16\pi^2\Lambda^2}\Big|_{\mu=\Lambda} \tag{6}$$

This four-quark operator will in turn induce a neutron EDM d_n. To evaluate this contribution, one must first evolve $C^{(1)}_{quqd}$ from $\mu = \Lambda$ down to $\mu = \mu_h$. In principle, this can only be done if one knows the exact value of Λ. However, since the evolution depends only logarithmically on Λ, it is reasonable to take $\Lambda \sim 1$ TeV as an illustration, giving $C^{(1)}_{quqd}(\mu_h) = 7.2 C^{(1)}_{quqd}(\Lambda)$ [25].

[1] It is possible that in the full theory, a symmetry implies vanishing matching contributions, but we will be more general here.

Next, in order to find the relation between $C^{(1)}_{ququd}(\mu_h)$ and the induced hadronic EDMs one needs to compute corresponding hadronic matrix elements. First-principle calculations of such matrix elements are challenging, and presently only exist for simple systems such as ρ-meson (see, e.g. Ref. [26] and references therein). The results of such calculations are generally consistent with the order-of-magnitude estimation based on naïve dimensional analysis (NDA) [27–29], so here we shall also provide an NDA estimation of d_n:

$$\begin{aligned} d_n &\sim e\frac{\Lambda_\chi}{16\pi^2}\frac{\mathrm{Im}C^{(1)}_{ququd}(\mu_h)}{\Lambda^2} \\ &\approx e\frac{\Lambda_\chi}{16\pi^2}\times 7.2\frac{\mathrm{Im}C^{(1)}_{ququd}(\Lambda)}{\Lambda^2} \\ &\approx 9.4\times 10^{-23}(\frac{v}{\Lambda})^2\mathrm{Im}\{c_1^* c_2\}|_{\mu=\Lambda}\, e\ \mathrm{cm}. \end{aligned} \tag{7}$$

This EDM is suppressed by $1/(16\pi^2)^2$ as well as Λ_χ/v. Given the current upper bound $d_n < 3.0\times 10^{-26} e$ cm at 90% CL [4] we see that the existing neutron EDM limits are probing $(v/\Lambda)^2 c^2 \sim 3\times 10^{-4}$ at $\mu=\Lambda$ which implies $\Lambda \gtrsim 10$ TeV if $c(\Lambda)\sim 1$.

At first glance, the neutron EDM sensitivity to Λ is slightly tighter than that of the D-coefficient. However, since both estimations made in Eq. (6) and (7) allow an error within an order of magnitude, one may reasonably conclude that the sensitivities of d_n and the D-coefficient are comparable. Furthermore, hadronic and atomic EDMs depend only on $c_1^* c_2$ and provide no direct constraint on the contribution from $c_j c_3^*$ $(j=1,2)$ in Eq. (3).

We now consider constraints from T conserving observables. First, we note that LHC studies of the process $pp\to e+X+MET$ place stringent bounds on the operators in (2).[2] Following Ref. [19], one may define two dimensionless quantities:

$$\begin{aligned} \tilde\epsilon_S &= -\frac{c_1-c_2}{2\sqrt{2}G_F V_{ud}\Lambda^2} \\ \tilde\epsilon_T &= \frac{c_3}{2\sqrt{2}G_F V_{ud}\Lambda^2}. \end{aligned} \tag{8}$$

The contribution from $\hat{O}_{1,2,3}$ to the total cross-section σ_{tot} of the $pp\to e+X+MET$ process measured by LHC can be written as $\sigma_{tot}=\sigma_S|\tilde\epsilon_S|^2+\sigma_T|\tilde\epsilon_T|^2$. Therefore LHC is sensitive to $|\tilde\epsilon_S|$ and $|\tilde\epsilon_T|$ while the D-coefficient probes the combination of products $\mathrm{Re}\tilde\epsilon_T\mathrm{Im}\tilde\epsilon_S - \mathrm{Re}\tilde\epsilon_S\mathrm{Im}\tilde\epsilon_T$. The bounds on the $\tilde\epsilon$ parameters obtained in Ref. [19] assume contributions from one operator at a time. However, when comparing with the D-coefficient sensitivity, one must take both $\tilde\epsilon_S$ and $\tilde\epsilon_T$, since the D-coefficient probes products of the two. Recasting the analysis of Ref. [19] is nevertheless straightforward because σ_S and σ_T are known. The constraint equation in Ref. [19] then implies an elliptical bound in the $|\tilde\epsilon_S|-|\tilde\epsilon_T|$ plane. One should also remember that the LHC constraints should be run down to $\mu=\mu_h$ for a fair comparison with the D-coefficient.

Since there are four real parameters in the problem (the Re and Im parts of $\tilde\epsilon_{S,T}$), it us useful to make simplifying assumptions in order to compare the LHC and D-coefficient sensitivities. To that end, we will assume for the moment that $\mathrm{Re}\tilde\epsilon_S=\mathrm{Im}\tilde\epsilon_T=0$ so both the LHC and the neutron D-coefficient results set constraints on $\mathrm{Re}\tilde\epsilon_T$ and $\mathrm{Im}\tilde\epsilon_S$ (see Fig. 1). In this case, one sees that the sensitivity of neutron decay to the D-coefficient has to be improved by roughly a factor of 15 in order to match the sensitivity of the 7-TeV LHC results. Results at $\sqrt{s}=8$ TeV for the same channel at

[2] We thank M. Gonzales-Alonso for pointing out these constraints.

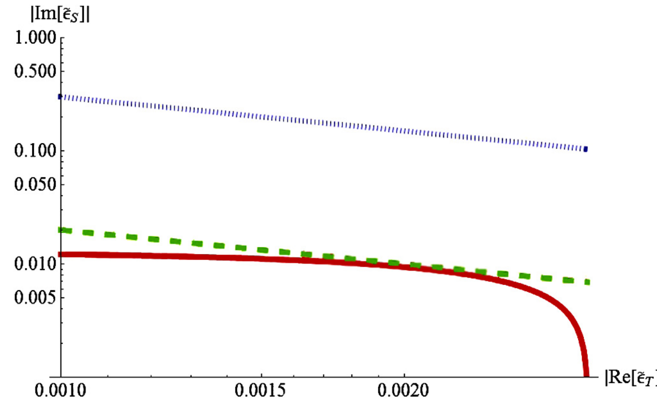

Fig. 1. (Color online) Exclusion plot for $\mathrm{Re}\tilde\epsilon_T$ and $\mathrm{Im}\tilde\epsilon_S$ at $\mu=\mu_h$ from the 7-TeV LHC data (red solid line) as well as the bound from the neutron D-coefficient with the current precision level (blue dotted line) and 15 times of the current precision level (green dashed line) respectively assuming $\mathrm{Re}\tilde\epsilon_S=\mathrm{Im}\tilde\epsilon_T=0$.

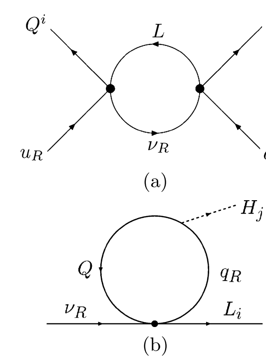

Fig. 2. Leading loop contributions that provide indirect bounds on c_1 and c_2. Figure (a) induces a four-quark operator that generates hadronic EDMs. Figure (b) generates a neutrino mass after electroweak symmetry breaking.

are also available. As there is no significant deviation from SM prediction [30,31], the LHC bound on $|\tilde\epsilon_S|$ and $|\tilde\epsilon_T|$ will be even more stringent than quoted above, although a detailed analysis has yet to be performed.[3]

One may also derive interesting but less direct constraints on $\hat{O}_1$ and $\hat{O}_2$ from the scale of neutrino mass and naturalness considerations. Above the electroweak scale, the leading contribution to m_ν comes from a one-loop diagram with a quark Yukawa insertion, inducing the Yukawa interaction term $\bar{L}\tilde{H}\nu_R$, as shown in Fig. 2b. Again this contribution vanishes in DR so we estimate it using simple dimensional analysis. After electroweak symmetry breaking, one obtains

$$m_\nu \sim \frac{c_i}{\Lambda^2}\frac{\Lambda^2}{16\pi^2}m_q = \frac{c_i m_q}{16\pi^2} \tag{9}$$

where m_q is the light quark mass and $i=1,2$. Taking $m_\nu < 1$ eV and $m_q\approx 5$ MeV we obtain $c_i < 3\times 10^{-5}$. We stress that this bound is not airtight, as the result may vary considerably, depending on the specific symmetry of the underlying BSM scenario. Neutrino mass naturalness bounds also do not constrain the tensor interaction strength c_3. Should a next generate D-coefficient measurement yield a non-vanishing result, the comparison with neutrino mass naturalness considerations would provide interesting input for model-building.

[3] The LHC sensitivity will, of course, improve further with the data obtained from Run II.

3. Operators without right-handed neutrinos

In contrast to the discussion of Section 2, we consider here $d=6$ operators that contain only left-handed (LH) neutrino fields and show that any contributions to the D-coefficient are severely constrained by present EDM limits. In the four-fermion sector, the only operator that gives a tree-level D-coefficient scaling linearly with the BSM coupling strength has the form of $\bar{u}_R\gamma^\mu d_R \bar{e}_L\gamma_\mu \nu_L$ as discussed in Ref. [16]. It is actually derived from a gauge-invariant dim-6 operator:

$$\hat{O}_{Hud} = i\frac{C_{Hud}}{\Lambda^2}(\tilde{H}^\dagger D_\mu H)(\bar{u}_R\gamma^\mu d_R) + \text{h.c.} \tag{10}$$

Below the electroweak scale, exchange of the W-boson contained in the covariant derivative with the left-handed charged weak current leads to both the semi-leptonic four-fermion operator listed above as well as a four-quark operator of the form $\bar{u}_R\gamma^\mu d_R \bar{d}_L\gamma_\mu u_L$. Both operators share the same Wilson coefficient (up to V_{ud}), which is tightly constrained by the four-quark contribution to the neutron EDM.

The three remaining semi-leptonic four-fermion operators that contain T-odd components are the scalar and tensor operators of the type $(\bar{L}R)(\bar{L}R)$ and $(\bar{L}R)(\bar{R}L)$ [18,23]:

$$\begin{aligned}
\hat{O}_{ledq} &= i\frac{\text{Im}C_{ledq}}{\Lambda^2}\bar{L}^i e_R \bar{d}_R Q^i + \text{h.c.}\\
\hat{O}^{(1)}_{lequ} &= i\frac{\text{Im}C^{(1)}_{lequ}}{\Lambda^2}\varepsilon^{ij}\bar{L}^i e_R \bar{Q}^j u_R + \text{h.c.}\\
\hat{O}^{(3)}_{lequ} &= i\frac{\text{Im}C^{(3)}_{lequ}}{\Lambda^2}\varepsilon^{ij}\bar{L}^i \sigma^{\mu\nu} e_R \bar{Q}^j \sigma_{\mu\nu} u_R + \text{h.c.}
\end{aligned} \tag{11}$$

Similar to the operators in Eq. (2), they induce a D-coefficient that scales quadratically with c_i/Λ^2. However, $\hat{O}_{ledq}$, $\hat{O}^{(1)}_{lequ}$, $\hat{O}^{(3)}_{lequ}$ contribute linearly to EDMs of paramagnetic atom and molecules and diamagnetic atoms at tree level as well as hadronic and electron EDMs at the one-loop level. In particular, the ACME limit on the EDM of ThO molecule implies a strong constraint on $\text{Im}(C_{ledq} - C^{(1)}_{ledq})$ (see, e.g. Ref. [32]). The resulting indirect constraints on the associated D-coefficient contributions are severe.

The remaining class of operators that give rise to the D-coefficient at tree-level are dipole-like operators. One may wonder whether EDM constraints to such operators may be avoided with an appropriate choice of Wilson coefficients at low energy. We will show, however, that this is not possible without fine-tuning at the level of many orders of magnitude. To simplify our discussion, let us concentrate on the dipole-like operators in the purely leptonic sector:

$$\begin{aligned}
\hat{O}_{eB} &= i\frac{g'\text{Im}C_{eB}}{\Lambda^2}\bar{L}\sigma^{\mu\nu}He_R B_{\mu\nu} + \text{h.c.}\\
\hat{O}_{eW} &= i\frac{g\text{Im}C_{eW}}{\Lambda^2}\bar{L}\sigma^{\mu\nu}\frac{\tau^i}{2}He_R W^i_{\mu\nu} + \text{h.c.}\\
\hat{O}_{eH^3} &= i\frac{\text{Im}C_{eH^3}}{\Lambda^2}\bar{L}He_R H^\dagger H + \text{h.c.}
\end{aligned} \tag{12}$$

The first-two operators are dipole-like while the third operator is included as well because it mixes with the first two via electroweak renormalization. Only $\hat{O}_{eW}$ contributes to D, as it is the only one containing a W field. After electroweak symmetry-breaking, one finds

$$D = -\frac{4\sqrt{2}g_A^2}{g_V^2 + 3g_A^2}\left(\frac{m_e}{v}\right)\left(\frac{v}{\Lambda}\right)^2 \text{Im}C_{eW} \quad . \tag{13}$$

Note the presence of the m_e/v suppression due to the existence of a derivative in the operator $\hat{O}_{eW}$. The current upper bound on the neutron D-coefficient implies $(v/\Lambda)^2|\text{Im}C_{eW}| < 1\times 10^2$.

The same set of operators also induces an electron EDM, given by

$$d_e = -\frac{\sqrt{2}e}{v}\left(\frac{v}{\Lambda}\right)^2(\text{Im}C_{eB} - \text{Im}C_{eW}) \tag{14}$$

The current upper bound on d_e [3] implies $(v/\Lambda)^2|\text{Im}C_{eB} - \text{Im}C_{eW}| < 7.7\times 10^{-13}$.

At first glance, it seems that one could simply choose $\text{Im}C_{eB} = \text{Im}C_{eW}$ at low energy to avoid the EDM constraint. We want to argue that, however, this choice is highly unnatural because the operators in Eq. (12) mix under electroweak renormalization as

$$\frac{d\Theta}{d\ln\mu} = \begin{pmatrix} \frac{151g'^2-27g^2}{192\pi^2} & -\frac{3gg'}{64\pi^2} & 0\\ -\frac{gg'}{16\pi^2} & \frac{-11g^2+3g'^2}{192\pi^2} & 0\\ -\frac{3g'(g^2-3g'^2)}{16\pi^2} & -\frac{9g(g^2-g'^2)}{32\pi^2} & -\frac{3(9g^2+7g'^2)}{64\pi^2}\end{pmatrix}\Theta \tag{15}$$

where $\Theta = \left(g'\text{Im}C_{eB}\ \ g\text{Im}C_{eW}\ \ \text{Im}C_{eH^3}\right)^T$. Numerically, if we assume that the bounds on the D-coefficient is marginally satisfied at $\mu = m_W$ (i.e. $(v/\Lambda)^2|\text{Im}C_{eW}| = 1\times 10^2$), then after the electroweak renormalization we find that $(v/\Lambda)^2|\text{Im}C_{eB} - \text{Im}C_{eW}| \approx 4.0$ at $\mu = 10$ TeV. However, this number has to be fine-tuned to a precision level of 2×10^{-11}% in order to satisfy the EDM bound at low energy, and therefore it is obviously not natural. The dipole-like operators in the quark sector suffer from the same problem. We conclude that, in the absence of RH neutrinos, EDM constraints imply that the existence of an observable D-coefficient is highly unlikely.

4. Conclusion

If neutrinos are Dirac particles, implying the existence of light ν_R in nature, then present limits on the D-coefficient indicate that the mass scale of any associated PCTV interactions may be quite significant: $\Lambda/c \gtrsim 1$ TeV, where c denotes a $d=6$ operator Wilson coefficient. The corresponding reach associated with limits on the PCTV triple correlation in polarized ^{19}Ne decay are somewhat weaker, but nevertheless quite interesting. The observation of a non-zero effect in a next generation experiment with either the neutron or nuclei is not precluded by constraints from EDM search null results. On the other hand, the LHC results for $pp \to e + X + MET$ present a greater challenge, implying at least an order of magnitude improvement in neutron decay PCTV correlation sensitivity would be needed for discovery of a non-zero D-coefficient. Should such an observation occur, resolving the tension between a non-zero PCTV correlation measurement and neutrino mass naturalness considerations would provide an interesting challenge for model building.

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