Improved bounds on heavy quark electric dipole moments

H. Gisbert\textsuperscript{1}\thanks{Electronic address: Hector.Gisbert@ific.uv.es} and J. Ruiz Vidal\textsuperscript{1}\thanks{Electronic address: Joan.Ruiz@ific.uv.es}

\textsuperscript{1}IFIC, Universitat de Val\`encia-CSIC, Valencia, Spain

(Dated: May 8, 2019)

New bounds on the electric dipole moment (EDM) of charm and bottom quarks are derived using the stringent limits on their chromo-EDMs. The new limits, $|d_c| < 1.5 \times 10^{-21}$ cm and $|d_b| < 1.2 \times 10^{-20}$ cm, improve the previous ones by about three orders of magnitude. These indirect bounds can have important implications for models of new physics.

PACS numbers: 13.40.Em, 14.65.Dw, 14.20.Lq, 31.30.jn 11.10.Hi

Searches for electric dipole moments (EDMs) are currently setting stringent constraints on models of new physics (NP) with additional $CP$-violation sources \cite{1}. Since the standard model predictions are well below the current experimental accuracy, any signal of a non-zero EDM would be a clear sign of NP. Moreover, the persisting B-anomalies suggest a non-trivial flavor structure in NP models, which can enhance the heavy quark EDMs \cite{6,7}. Due to their very small lifetime, direct EDM searches on heavy-flavoured hadrons represent an experimental challenge and only indirect limits on heavy quark dipole couplings have been obtained to date. However, this situation may change with the new proposals to search for the EDM of charmed and bottom baryons at the LHC \cite{9,11}. In this Letter, a new approach for setting indirect bounds on quark EDM couplings is presented. By exploiting the mixing of operators under the renormalization group and using constraints on the chromo-EDM of charm and bottom quarks \cite{12,13}, we extract new bounds on their corresponding EDMs that improve the current ones by several orders of magnitude.

For that purpose, let us consider the following flavour-conserving $CP$-violating effective Lagrangian

$$\mathcal{L}_{\text{eff}} = \sum_{i=1}^{2} \sum_{q} C_{i}^{q}(\mu) O_{i}^{q}(\mu) + C_{3}(\mu) O_{3}(\mu) , \quad (1)$$

where the index $q$ runs over the relevant flavours at the chosen renormalization scale. The effective operators are defined as

$$O_{1}^{q} = -\frac{i}{2} e Q_q m_q \bar{q}^{\alpha} \sigma^{\mu\nu} \gamma_5 q^\alpha F_{\mu\nu} ,$$

$$O_{2}^{q} = -\frac{i}{2} g_s m_q \bar{q}^{\alpha} \sigma^{\mu\nu} T_{a} \gamma_5 q^\alpha G_{\mu\nu}^{a} ,$$

$$O_{3} = -\frac{1}{6} g_s f_{abc} e^{\mu\nu\lambda}\kappa G_{\mu\rho}^{a} C_{\nu}^{b} C_{\lambda}^{c} , \quad (2)$$

where $Q_q$ and $m_q$ are the quark charge and quark mass, respectively. The quark EDM, chromo-EDM, and the usually defined coefficient $\omega(\mu)$ of the Weinberg operator are related to the Wilson coefficients by

$$d_q(\mu) = e Q_q m_q(\mu) C_{1}(\mu) ,$$

$$\bar{d}_q(\mu) = m_q(\mu) C_{2}^{\pm}(\mu) , \quad (3)$$

$$\omega(\mu) = -\frac{1}{2} g_s(\mu) C_{3}(\mu) .$$

When a heavy quark is integrated out, its chromo-EDM gives a finite contribution to the Weinberg operator \cite{13,15}, which is strongly constrained from the limits on the neutron EDM. This allows to bound the quark chromo-EDMs to be \cite{12,13},

$$|\bar{d}_c(m_c)| < 1.0 \times 10^{-22} \text{ cm} ,$$

$$|\bar{d}_b(m_b)| < 1.1 \times 10^{-21} \text{ cm} . \quad (4)$$

Attempts to constraint heavy quark EDMs have followed different strategies: flavor-mixing contributions into light quark EDMs \cite{12,16,17}, $b \rightarrow s\gamma$ transitions \cite{12}, mixing into the electron EDM via light-by-light scattering diagrams \cite{17} and tree-level contributions to the $e^+e^- \rightarrow q\bar{q}$ total cross section \cite{18,19}. All of these approaches yield results within the same order of magnitude, the most restrictive ones being \cite{12,19},

$$|d_c(m_c)| < 4.4 \times 10^{-17} \text{ cm} ,$$

$$|d_b(m_b)| < 2.0 \times 10^{-17} \text{ cm} . \quad (5)$$

In this work we follow a new strategy that relates the EDM and chromo-EDM operators in order to find new limits on $\bar{d}_q$ from the already available strong bounds on $d_q$. This relation is done in a model-independent way using the renormalization group equations, which mix the effective operators when the energy scale is changed. The relevant diagrams include photon loops which have been neglected in previous works due to its small size compared with pure QCD corrections. Nevertheless, they represent the first non-zero contribution to the mixing we are interested in.
The evolution of the Wilson coefficients is given by

\[ \frac{d}{d \ln \mu} \overline{C}(\mu) = \tilde{\gamma}^T \overline{C}(\mu) , \tag{6} \]

where \( \overline{C} \equiv (C_7, C_9, C_{10}) \) and \( \tilde{\gamma} \) is the anomalous dimension matrix. This matrix can be expanded in powers of the QCD and QED coupling constants, \( \alpha_s \) and \( \alpha_e \), respectively,

\[ \tilde{\gamma} = \frac{\alpha_s}{4 \pi} \gamma_s(0) + \left( \frac{\alpha_s}{4 \pi} \right)^2 \gamma_s(1) + \frac{\alpha_e}{4 \pi} \gamma_e(0) + \cdots , \tag{7} \]

where \( \gamma_s(0) \) and \( \gamma_s(1) \) represent the one- and two-loop QCD corrections, while \( \gamma_e(0) \) encodes the one-loop QED correction \( \mathcal{O}(\alpha_e^2) \). At \( \mathcal{O}(\alpha_s^2) \), the quark EDM does not mix into the chromo-EDM and the first contribution only appears at \( \mathcal{O}(\alpha_e) \) from photon-loop diagrams as shown in Figure 1. Applying the standard techniques for the computation of anomalous dimensions \( \mathcal{O}(\alpha_s^3) \) we obtain the matrix element \( (\gamma_e)(0)_{12} = 8 \), in agreement with a previous calculation \( \mathcal{O}(\alpha_s^4) \).

Solving Eq. (6) by adding this contribution, the evolution of the charm and bottom chromo-EDMs read

\[ \tilde{d}_c(m_c) = -0.04 \frac{d_c(M_{\text{NP}})}{e} + 0.74 \tilde{d}_c(M_{\text{NP}}) , \tag{8} \]

\[ \tilde{d}_b(m_b) = 0.08 \frac{d_b(M_{\text{NP}})}{e} + 0.88 \tilde{d}_b(M_{\text{NP}}) , \tag{9} \]

where we have taken \( M_{\text{NP}} \sim 1 \text{ TeV} \) as the scale of NP. In this result, we have neglected the mixing of the Weinberg operator into the chromo-EDM due to the very strong bounds on \( \omega \) from constraints on the neutron EDM \( \mathcal{O}(\alpha_e^2) \). The mixing of \( \tilde{d}_b \) into itself, described by the second piece of Eqs. (8) and (9), has leading contributions from pure QCD corrections, then corrections of \( \mathcal{O}(\alpha_e) \) can be safely neglected.

Using the bounds on the chromo-EDMs at the low scales quoted in Eq. (4), the parameter space on the \( \tilde{d}_c-d_e \) plane is constrained as shown in Figure 2. Strong fine-tuned cancellations between the two pieces of Eqs. (8) and (9) result in an allowed region extending along a straight line which is unlikely to be realised in NP models.

Hence, we assume constructive interference between the EDM and chromo-EDM contributions at the NP scale to extract bounds on \( d_{e}(M_{\text{NP}}) \). Then, using the evolution of the EDM operator to bring these bounds down to the quark mass scale, the new bounds on the charm and bottom quark EDMs are

\[ |d_c(m_c)| < 1.5 \times 10^{-21} \text{ cm} , \]

\[ |d_b(m_b)| < 1.2 \times 10^{-20} \text{ cm} , \tag{10} \]

which improve the previous ones quoted in Eq. (5) by three and four orders of magnitude, respectively. This approach does not improve the current bounds on the top quark EDM \( \mathcal{O}(\alpha_e^2) \) given that the limit on its chromo-EDM is of similar size \( \mathcal{O}(\alpha_e^3) \). The theoretical uncertainty of this result is dominated by the contribution of the Weinberg operator to the neutron EDM, since it determines the size of the chromo-EDM bounds. Note also that higher values of the NP scale yield less conservative results, e.g.

\[ m_{\text{NP}} > \left( \frac{|d_{e}(m_{\text{NP}})|}{|d_{e}(M_{\text{NP}})|} \right)^{1/2} \text{ TeV} , \]

where \( m_{\text{NP}} \) is the scale of NP.
a 30% stronger bounds for $M_{NP} = 10$ TeV.

The new constraints for the charm and bottom quark EDMs are in tension with the predictions of different theories beyond the standard model [29–32] and will provide valuable input for future phenomenological analysis of NP models.

We want to express our gratitude to W. Dekens, A. Delhom, M. Cerda-Sevilla, F. Martínez-Vidal, N. Neri and A. Pich for useful discussions. This work has been supported in part by the Spanish State Research Agency and ERDF funds from the EU Commission [Grants FPA2017-84445-P, FPA2017-85140-C3-3-P, FPA2015-68318-R and FPA2014-53631-C2-1-P], by Generalitat Valenciana [Grant Prometeo/2017-053 and Grant Prometeo/2014-049], and by the Spanish Centro de Excelencia Severo Ochoa Programme [Grant SEV-2014-0398]. The work of H.G. is supported by a FPI doctoral contract [BES-2015-073138], funded by the Spanish Ministry of Economy, Industry and Competitiveness.

[1] T. Chupp, P. Fierlinger, M. Ramsey-Musolf and J. Singh, Rev. Mod. Phys. 91 (2019) no.1, 015001.
[2] N. Yamanaka, B. K. Sahoo, N. Yoshinaga, T. Sato, K. Asahi and B. P. Das, Eur. Phys. J. A 53 (2017) no.3, 54.
[3] M. Pospelov and A. Ritz, Annals Phys. 318 (2005) 119.
[4] J. Engel, M. J. Ramsey-Musolf and U. van Kolck, Prog. Part. Nucl. Phys. 71 (2013) 21.
[5] W. Dekens, J. de Vries, J. Bsaisou, W. Bernreuther, C. Hanhart, U. G. Meißner, A. Nogga and A. Wirzba, JHEP 1407 (2014) 069.
[6] D. Buttazzo, A. Grelio, G. Isidori and D. Marzocca, JHEP 1711 (2017) 044.
[7] W. Dekens, J. de Vries, M. Jung and K. K. Vos, JHEP 1901 (2019) 069.
[8] F. J. Botella, L. M. Garcia Martin, D. Marangotto, F. Martinez Vidal, A. Merli, N. Neri, A. Oyanguren and J. Ruiz Vidal, Eur. Phys. J. C 77 (2017) no.3, 181.
[9] E. Bagli et al., Eur. Phys. J. C 77 (2017) no.12, 828.
[10] A. S. Fomin et al., JHEP 1708 (2017) 120.
[11] V. G. Baryshevsky, Phys. Lett. B 757 (2016) 426.
[12] F. Sala, JHEP 1403 (2014) 061.
[13] D. Chang, W. Y. Keung, C. S. Li and T. C. Yuan, Phys. Lett. B 241 (1990) 589.
[14] E. Braaten, C. S. Li and T. C. Yuan, Phys. Rev. Lett. 64 (1990) 1709.
[15] G. Boyd, A. K. Gupta, S. P. Trivedi and M. B. Wise, Phys. Lett. B 241 (1990) 584.
[16] A. Cordero-Cid, J. M. Hernandez, G. Tavares-Velasco and J. J. Toscano, J. Phys. G 35 (2008) 025004.
[17] A. G. Grozin, I. B. Khriplovich and A. S. Rudenko, Nucl. Phys. B 821 (2009) 285.
[18] R. Escribano and E. Masso, Nucl. Phys. B 429 (1994) 19.
[19] A. E. Blinov and A. S. Rudenko, Nucl. Phys. Proc. Suppl. 189 (2009) 257.
[20] S. Weinberg, Phys. Rev. Lett. 63 (1989) 2333.
[21] F. Wilczek and A. Zee, Phys. Rev. D 15 (1977) 2660.
[22] G. Degrassi, E. Franco, S. Marchetti and L. Silvestrini, JHEP 0511 (2005) 044.
[23] E. E. Jenkins, A. V. Manohar and P. Stoffer, JHEP 1801 (2018) 084.
[24] G. Buchalla, A. J. Buras and M. E. Lautenbacher, Rev. Mod. Phys. 68 (1996) 1125.
[25] A. J. Buras, hep-ph/9806471.
[26] C. A. Baker et al., Phys. Rev. Lett. 97 (2006) 131801.
[27] V. Cirigliano, W. Dekens, J. de Vries and E. Mereghetti, Phys. Rev. D 94 (2016) no.1, 016002.