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The Electric Dipole Moment of the electron in the decoupling limit of the A2HDM

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We conduct an analysis of model-independent contributions to the Electric Dipole Moment (EDM) of the electron. Our study focuses on contributions that originate from a linearly independent heavy scalar sector. Specifically, we discuss the decoupling limit of the aligned two-Higgs-doublet model (A2HDM). In this framework, Barr-Zee diagrams involving a fermion loop induce logarithmically-enhanced terms proportional to potentially sizable new sources of CP violation. In the decoupling regime these contributions originate from effective dimension-6 operators through the renormalization-group mixing of four-fermion operators into electroweak dipole operators. These logarithmic contributions are absent in more constrained $\mathcal{Z}_2$-symmetric realizations of the 2HDM, where this discrete symmetry dictates the structure of effective operators providing new-physics contributions to the electron EDM and thereby yields a suppression mechanism. During our comparison of the results from the A2HDM with the leading logarithms from SMEFT, we have identified and corrected signs of expressions found in the literature. Finally, we examine how the current experimental bounds on the electron EDM constrain the parameter space of the A2HDM.

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1. Introduction

Phenomena sensitive to Charge-Parity (CP) violation provide a powerful probe of the Standard Model (SM) structure. In searching for New Physics (NP) sources of CP violation, Electric Dipole Moments (EDMs) play a central role, as experimental sensitivities have attained impressive precision, while SM contributions remain largely suppressed.

The presence of additional scalar degrees of freedom enriches the phenomenology of CP violation. New complex phases arising from the scalar sector can modify the Kobayashi-Maskawa picture, leading to several scalar-mediated contributions to EDMs.

The Two-Higgs Doublet Model (2HDM) [1, 2], which extends the SM by introducing an additional scalar doublet, cosntitutes one of the most economical extensions of the SM. The 2HDM suffers in general from tree-level Flavour-Changing Neutral Currents (FCNCs), which are severely constrained by experimental data. A discrete $\mathcal{Z}_2$ symmetry is tipically imposed on the Lagrangian such that each type of right-handed fermion couples to a single scalar doublet, eliminating tree-level FCNCs [3]. Alternatively, the 'Aligned Two-Higgs Doublet Model' (A2HDM) solves this issue by imposing flavour alignment of Yukawa couplings such that the Yukawa matrices of the two scalar doublets are proportional to each other in flavour space [4, 5].

In the A2HDM, the SM is extended with a second complex scalar doublet with hypercharge $Y = 1/2$. After EW symmetry breaking, both doublets acquire complex vacuum expectation values (VEVs). However, it is always possible to rotate the scalar doublets through an SU(2) transformation to the "Higgs-basis" where only the first doublet acquires a non-zero real VEV. In this basis, the scalar doublets read:

$$\Phi_1 = \frac{1}{\sqrt{2}}\begin{pmatrix} \sqrt{2}\,G^+ \\ S_1 + v + i\,G^0 \end{pmatrix}, \qquad \Phi_2 = \frac{1}{\sqrt{2}}\begin{pmatrix} \sqrt{2}\,H^+ \\ S_2 + i\,S_3 \end{pmatrix}, \tag{1}$$

where $v = 246$ GeV is the VEV of the Φ_1 doublet and the components $G^\pm$ and G^0 act as Goldstone bosons. Thus, the scalar sector of this model consists of a pair of charged scalars $H^\pm$, two CP-even scalars S_1, S_2 and a CP-odd scalar (pseudoscalar) S_3. Imposing the gauge invariance from SM, the most general scalar potential reads:

$$\begin{aligned} V =& \mu_1\,\Phi_1^\dagger\Phi_1 + \mu_2\,\Phi_2^\dagger\Phi_2 + \left[\mu_3\,\Phi_1^\dagger\Phi_2 + h.c.\right] + \frac{\lambda_1}{2}\,(\Phi_1^\dagger\Phi_1)^2 + \frac{\lambda_2}{2}\,(\Phi_2^\dagger\Phi_2)^2 + \lambda_3\,(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) \\ &+ \lambda_4\,(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) + \left[\left(\frac{\lambda_5}{2}\,\Phi_1^\dagger\Phi_2 + \lambda_6\,\Phi_1^\dagger\Phi_1 + \lambda_7\,\Phi_2^\dagger\Phi_2\right)(\Phi_1^\dagger\Phi_2) + \text{h.c.}\right], \end{aligned} \tag{2}$$

where $\lambda_{\{5,6,7\}}$ and μ_3 are complex parameters in general. Depending on the parameters of the scalar potential, the neutral scalars S_1, S_2, S_3 mix with each other through an orthogonal matrix $\mathcal{R}$, yielding the mass eigenstates $H_j \in \{H_1, H_2, H_3\}$, where $H_1 = h$ is identified as the lightest state. If some of the parameters in the scalar potential have non-zero complex phases, CP-symmetry is violated, so that the mass eigenstates do not possess a definite CP quantum number.

On the other hand, apart from fermion mass terms, the interaction part of the Yukawa Lagrangian in this model reads:

$$-\mathcal{L}_Y = \sum_{j,f}\left(\frac{y_f^{H_j}}{v}\right) H_j\,\bar{f}M_f\mathcal{P}_R f + \frac{\sqrt{2}H^+}{v}\left[\bar{u}\left\{\varsigma_d V M_d\mathcal{P}_R - \varsigma_u M_u^\dagger V\mathcal{P}_L\right\} d + \varsigma_l\,\bar{\nu}M_l\mathcal{P}_R\,l\right] + \text{h.c.}\,, \tag{3}$$

where $\mathcal{P}_{L,R}$ are chirality projectors, M_f are the diagonal fermion mass matrices, and V is the CKM matrix. The complex coefficients ς_f are called the flavour alignment parameters. The Yukawa couplings of fermions with neutral scalars ($y_f^{H_j}$) are given by:

$$y_{d,l}^{H_j} = \mathcal{R}_{j1} + (\mathcal{R}_{j2} + i\mathcal{R}_{j3})\varsigma_{d,l} \qquad \text{and} \qquad y_u^{H_j} = \mathcal{R}_{j1} + (\mathcal{R}_{j2} - i\mathcal{R}_{j3})\varsigma_u^* \,. \tag{4}$$

The usual 2HDM scenarios can be retrieved from the A2HDM Lagrangian by imposing the following conditions on the alignment parameters:

$$\begin{aligned} &\text{Type I: } \varsigma_u = \varsigma_d = \varsigma_l = \cot\beta, \quad \text{Type II: } \varsigma_u = -\frac{1}{\varsigma_d} = -\frac{1}{\varsigma_l} = \cot\beta\,, \quad \text{Inert: } \varsigma_u = \varsigma_d = \varsigma_l = 0\,, \\ &\text{Type X: } \varsigma_u = \varsigma_d = -\frac{1}{\varsigma_l} = \cot\beta \qquad \text{and} \qquad \text{Type Y: } \varsigma_u = -\frac{1}{\varsigma_d} = \varsigma_l = \cot\beta\,, \end{aligned} \tag{5}$$

along with vanishing $\widetilde{\lambda}_6$ and $\widetilde{\lambda}_7$ terms in the $\mathcal{Z}_2$-symmetric basis.

When studying CP violation in the leptonic sector, operators other than the electron dipole must also be adressed, since four-fermion semileptonic operators generated by the exchange of heavy scalars induce CP-violating electron-nuclon interactions [6]. For the heavy quarks, this leads to lepton-gluonic operators that could be of help in constraining NP couplings. Their discussion is more involved, since both lepton-gluonic operators and operators involving light quarks require dealing with the non-perturbative regime of strong interactions [7, 8].

2. Contributions to the eEDM in the A2HDM

CP violation from the CKM matrix of the SM is known to be insufficient to account for the observed *baryon asymmetry of the universe* [9]. This motivates the study of 2HDM scenarios that can accomodate new sources of CP violation. While in most of these 2HDM scenarios these new CP-violating terms come solely from the scalar potential, the A2HDM allows for CP violation in both the scalar and Yukawa sectors [4]. This CP violation in the Yukawa sector gives rise to new two-loop contributions to the electron EDM (eEDM).

The eEDM d_e can be defined as the coefficient of the following dimension-5 operator in the effective Lagrangian at a very low energy scale ($\mu \ll m_e$):

$$\mathcal{L} \supset -\frac{i}{2} d_e(\mu)\, \bar{\psi}_e\, \sigma^{\mu\nu}\gamma^5\, \psi_e\, F_{\mu\nu} \tag{6}$$

where ψ_e is the electron Dirac spinor and $F^{\mu\nu}$ is the field strength tensor of the photon. In the SM, teh eEDM is dominated by long-distance hadronic effects, giving $d_e^{SM} = 5.8 \times 10^{-40} e\,\text{cm}$ [10], well below the current 90% C.L. experimental bound [11]:

$$|d_e^{exp}| < 4.1 \times 10^{-30}\, e\,\text{cm} \tag{7}$$

It should be noted that the quoted experimental bound on d_e has been extracted from a diatomic molecule, assuming that the eEDM is the only source of CP violation. Diatomic molecules are also sensitive to a CP-odd pseudoscalar-scalar electron-nucleon coupling C_S. It has been estimated that this effective interaction produces a CP-violating effect equivalent in size to $d_e^{SM} \sim 10^{-38}\, e\,\text{cm}$ [12].

In the A2HDM, contributions to the eEDM arise at one loop, but it is well-established that the contribution from two-loop diagrams are the dominant ones [13], since these diagrams contain fewer electron-scalar interaction vertices which generate a mass suppression. The leading two-loop diagrams contributing to the eEDM in the A2HDM can be classified in two classes depending on the number of gauge bosons attached to the electron line: a) Barr-Zee (one gauge boson) and b) Kite (two gauge bosons). Barr-Zee diagrams can can be further divided into three categories depending on the gauge boson connecting the electron line and the internal loop: i) electromagnetic-current mediated (γ), ii) neutral-current mediated (Z) and iii) charged-current mediated ($W^{\pm}$). Each of these can involve three different types of loops: fermion loop, charged Higgs loop and gauge boson loop. We refer to i)-iii) as EM, NC, and CC contributions. The kite diagrams generate two types of contributions from the neutral and charged currents. The combined contribution of these diagrams to the eEDM is given by:

$$d_e = \sum_{L,X} d_{e,L}^{X} + (d_{e,\text{Kite}}^{CC} + d_{e,\text{Kite}}^{NC}), \tag{8}$$

where $L \in \{f,\ H^{\pm},\ W\}$ and $X \in \{\text{CC},\ \text{NC},\ \text{EM}\}$.

In contrast to the Complex 2HDM (C2HDM), which is a special case of the A2HDM, the eEDM in the A2HDM acquires additional contributions from charged-current fermion-loop Barr-Zee diagrams, which are absent in the C2HDM scenario [14]. This contribution is generated by the complex coupling of the $\bar{e}\nu H^{+}$ vertex, which is real in the C2HDM.

2.1 Full model contributions to the eEDM in the decoupling limit of the A2HDM

The scenario where the mass parameter of the second doublet Φ_2 becomes much larger than the VEV of Φ_1, i.e. $\sqrt{\mu_2} \gg v$, is is referred to as the *decoupling limit*. In this regime, the doublet Φ_2 sits at a very high energy scale compared to the SM and decouples. Using the masses of the particles as independent parameters, the condition the decoupling limit can be expressed as $M_{\{H^{\pm},\,H_2,\,H_3\}} \approx M \gg m_h$.

In our work, we computed the leading logarithmic contributions in the decoupling limit of the A2HDM. We found that fermion-loop Barr-Zee diagrams generate a squared logarithm term dependent on ς_u:

$$\left.\frac{d_{e,f}}{e}\right|_{\varsigma_u} = \sum_{X} \left.\frac{d_{e,f}^{X}}{e}\right|_{\varsigma_u} \overset{\text{dec.}}{\approx} m_e \frac{g^2}{(4\pi)^4 v^2}(3+5t_w^2)\,\text{Im}(\varsigma_u^{*}\varsigma_l)\frac{m_t^2}{M^2}\log^2\left(\frac{M^2}{m_t^2}\right), \tag{9}$$

and a single logarithmic term dependent on ς_d:

$$\left.\frac{d_{e,f}}{e}\right|_{\varsigma_d} = \sum_{X} \left.\frac{d_{e,f}^{X}}{e}\right|_{\varsigma_d} \overset{\text{dec.}}{\approx} -m_e \frac{g^2}{(4\pi)^4 v^2}\frac{t_w^2}{2}\,\text{Im}(\varsigma_d^{*}\varsigma_l)\frac{m_b^2}{M^2}\log\left(\frac{M^2}{m_t^2}\right), \tag{10}$$

while the combination of contributions from gauge boson loop Barr-Zee diagrams and charged-current kite diagrams generates another single logarithm [14]:

$$\sum_{X}\frac{d_{e,W}^{X}}{e} + \frac{d_{e,\text{Kite}}^{CC}}{e} \overset{\text{dec.}}{\approx} m_e \frac{g^2}{(4\pi)^4}\frac{3}{4}t_w^2\,\frac{\text{Im}(\lambda_6^{*}\varsigma_l)}{M^2}\log\left(\frac{M^2}{m_W^2}\right). \tag{11}$$

Our findings agree with the recent Ref. [15] which discusses the eEDM in the General 2HDM (G2HDM).

2.2 Model-independent contributions to the eEDM

To have better understanding of the A2HDM in the decoupling limit, we make use of the EFT approach. This approach is only valid when the energy scale of NP (i.e. the masses of the heavy scalars) lies well above the EW scale. In this regime, the effects of these heavy states are captured by effective operators Q_i of dimension higher than 4, associated with Wilson coefficients C_i:

$$\mathcal{L} = \mathcal{L}_{SM} + \sum_i C_i(\mu) Q_i. \tag{12}$$

Following the renormalization group equations (RGEs), these coefficients will run from the NP scale, where they are generated, down to the EW scale, mixing with the coefficients of the EW dipole operators and thus to the eEDM coefficient:

$$d_e = -\sqrt{2}v\Big(c_w \mathrm{Im}(C_{eB}) - s_w \mathrm{Im}(C_{eW})\Big). \tag{13}$$

Within the framework of the SM Effective Field Theory (SMEFT), the leading contribution to the eEDM will come from dimension-6 effective operators. In order to reproduce the leading logarithmic terms obtained in the decoupling limit of the A2HDM, we only need to take into account 4 independent effective non-dipole operators from the well-known Warsaw basis [16].

First, the dimension-6 Yukawa operator Q_{eH} and the four-fermion scalar operator Q_{ledq}:

$$Q_{eH}^{pr} = (H^\dagger H)(\bar{l}_p e_r H), \tag{14}$$
$$Q_{ledq}^{prmn} = (\bar{l}_p^j e_r)(\bar{d}_m q_n^j), \tag{15}$$

both of which mix at two-loop level into the EW dipole operators [17], generating two single logarithm contributions:

$$\frac{d_{e,eH}^{\mathrm{SMEFT}}}{e} = m_e \frac{g^2}{(4\pi)^4}\frac{3}{4}t_w^2 \frac{\mathrm{Im}(\lambda_6^* \varsigma_l)}{M^2} \log\left(\frac{M^2}{m_{EW}^2}\right), \tag{16}$$
$$\frac{d_{e,b}^{\mathrm{SMEFT}}}{e} = -m_e \frac{g^2}{(4\pi)^4 v^2}\frac{t_w^2}{2}\mathrm{Im}(\varsigma_d^* \varsigma_l)\frac{m_b^2}{M^2} \log\left(\frac{M^2}{m_{EW}^2}\right). \tag{17}$$

Then, the four-fermion scalar operator $Q_{lequ}^{(1)}$:

$$Q_{lequ}^{(1),prmn} = (\bar{l}_p^j e_r)\epsilon_{jk}(\bar{q}_m^k u_n), \tag{18}$$

which mixes at one-loop level into the four-fermion tensor operator $Q_{lequ}^{(3)}$:

$$Q_{lequ}^{(3),prmn} = (\bar{l}_p^j \sigma_{\mu\nu} e_r)\epsilon_{jk}(\bar{q}_m^k \sigma^{\mu\nu} u_n), \tag{19}$$

which in turn gets mixed into the EW dipoles generating a squared logarithm contribution to the eEDM:

$$\frac{d_{e,t}^{\mathrm{SMEFT}}}{e} = m_e \frac{g^2}{(4\pi)^4 v^2}(5t_w^2 + 3)\mathrm{Im}(\varsigma_u^* \varsigma_l)\frac{m_t^2}{M^2} \log^2\left(\frac{M^2}{m_{EW}^2}\right). \tag{20}$$

3. Phenomenology

We now investigate the impact of the novel charged-current fermion-loop Barr-Zee contributions on the prediction of the eEDM. With this in mind, we fix the parameters from the scalar potential of the A2HDM to the following benchmark values consistent with Ref. [18]:

$$\begin{aligned} \lambda_3 &= 0.02\,, & \lambda_4 &= 0.04\,, & \lambda_7 &= 0.03\,, & \mathrm{Re}(\lambda_5) &= 0.05\,, \\ \mathrm{Re}(\lambda_6) &= -0.05\,, & \mathrm{Im}(\lambda_6) &= 0.01\,, & \alpha_3 &= \pi/6. \end{aligned} \tag{21}$$

In Fig. 1 (left panel), we show a scatter plot of the value of eEDM for randomly generated values of the mass of the charged scalar, $M_{H^\pm} = M$, and the phases of the flavour alignment parameters, that is, ς_i with $i = u, d, l$. The blue dots are full-model predictions from the A2HDM, while the orange ones are computed by subtracting the new charged-current fermion-loop Barr-Zee contributions from the full result. The black line shows the prediction for the scenario with real ς_i. Finally, we show in gray the upper bound on the modulus of the eEDM from Eq. (7). The current experimental limit on the eEDM can be satisfied with complex ς_i even at a relatively low mass of 400 GeV, which is challenging to achieve with real ς_i parameters, as in the case of the C2HDM.

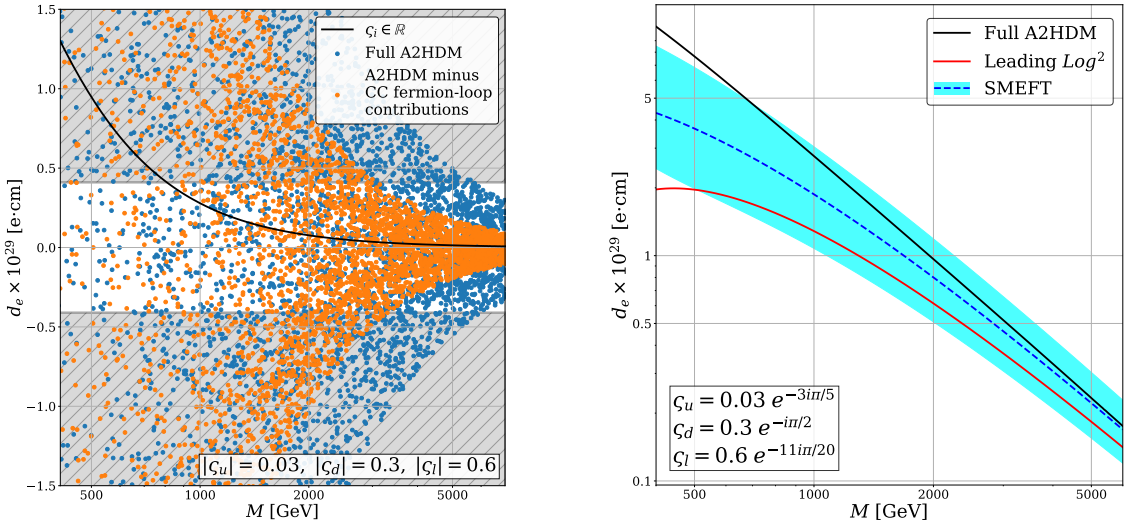

Figure 1: Left: Scatter plot of the eEDM in the A2HDM as a function of M. The blue dots correspond to the full A2HDM, while the expression used to generate the orange dots does not include CC fermion-loop Barr-Zee contributions. The black line assumes real alignment parameters. The gray bands show the upper bound on the modulus of the eEDM. **Right**: Approximations to predictions of the eEDM in the A2HDM as a function of M. The black line is the full two-loop result in the A2HDM. The solid red curve is the leading squared logarithmic approximation, and the dashed blue curve includes the sub-leading logarithms from SMEFT. The shaded blue region is obtained by varying the UV scale from $M/2$ up to $2M$.

In Fig. 1 (right panel), we numerically compare various approximations to the eEDM with respect to the mass of the charged scalar, $M_{H^\pm} = M$. The rest of the parameters are fixed according to the benchmark point in Eq. (21). The black line shows the result of the full two-loop calculation in the A2HDM. The solid red curve shows the leading squared logarithmic approximation. For the values of M displayed in the plot and the mentioned benchmark point, it provides an accurate

prediction for the eEDM. Finally, the dashed blue line shows the SMEFT result, and the shaded band is obtained by varying the Ultra-Violet (UV) scale at which the SMEFT logarithmic terms are evaluated between $M/2$ and $2M$.

4. Conclusions

We present a discussion of the contributions to the eEDM arising from the A2HDM. In contrast to more constrained 2HDM realizations, the A2HDM features additional sources of CP violation in the fermionic sector, giving rise to new contributions to EDMs.

Our study focuses on the contributions from the A2HDM in the decoupling limit. The features observed in this regime can be interpreted using dimension-6 operators of SMEFT. We explicitly demonstrate that the leading logarithmic contributions are captured by a small set of effective dimension-6 operators. Some of these contributions are absent in more extensively studied 2HDMs, where a discrete $\mathcal{Z}_2$ symmetry is applied, constraining the appearance of logarithmically enhanced contributions and providing a suppression mechanism.

Our phenomenological analysis shows that the contributions from charged-current fermion-loop Barr-Zee contributions can dominate over other contributions to the eEDM. Beyond the suppression mechanism induced by a a $\mathcal{Z}_2$ symmetry, accidental cancellation mechanisms may also occur, for instance between fermionic and non-fermionic Barr-Zee diagrams. Therefore, with the same values of the parameters in A2HDM and C2HDM, we can have lower masses of the NP scalars and still satisfy the experimental upper bound in the first case, just by varying the phases of the alignment parameters.

A natural direction for a future work would be to investigate other effects exclusive to the A2HDM, or to more general 2HDM realizations, but that are absent in the more constrained versions of these models.

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