Cancellation mechanism for the electron electric dipole moment connected with baryon asymmetry of the Universe

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We elucidate a cancellation mechanism for the electric dipole moment of the electron in the general two Higgs doublet model. The impressive improvement by ACME Collaboration in 2018 suggests the presence of a new electron Yukawa coupling that brings in exquisite cancellations among dangerous diagrams, broadening the solution space for electroweak baryogenesis driven by an extra top Yukawa coupling. The cancellation mechanism calls for the new Yukawa couplings to have hierarchical structures that echo the observed pattern of the Standard Model Yukawa couplings.

Introduction.— It is remarkable that the Cabibbo-Kobayashi-Maskawa (CKM) framework is able to explain all laboratory-based measurements of charge-parity, or CP, violation (CPV). But it is well known that the CPV phase arising from the CKM matrix is by far insufficient in generating the baryon asymmetry of the Universe (BAU), hence some new CPV phase(s) must exist to address this cosmological problem. Thus, in many well-motivated models beyond the Standard Model (SM), the existence of such beyond CKM phases is often a common theme. Detecting the effect of such new CPV phases would provide a powerful probe of new energy thresholds above the electroweak (EW) scale.

Owing to its high testability, EW baryogenesis (EWBG) is of primary importance and broad interest. However, data from the Large Hadron Collider (LHC), such as the measurement of Higgs boson properties, have diminished or even completely eliminated the EWBG parameter space in most models. Complementary to collider probes, extreme low-energy searches such as electric dipole moment (EDM) of the electron, neutron etc., have put further stress on models. In particular, with the new upper bound on the electron EDM (eEDM), \( |d_e| < 1.1 \times 10^{-29} \) e cm at 90\% confidence level (C.L.), given by the ACME collaboration in 2018 (ACME18), many EWBG scenarios are now in jeopardy. Although calculations of CPV sources still have significant uncertainties, hence the amount of BAU might go upward by more refined analyses, the impact of the ACME18 bound is nevertheless overwhelming. In other words, if EWBG is the true mechanism behind BAU, the unprecedented ACME18 result may indicate some undisclosed mechanism that renders \( d_e \) small.

In a previous paper, we have explored the general 2HDM (g2HDM), i.e. without the ad hoc discrete \( Z_2 \) symmetry, where an additional \( 3 \times 3 \) Yukawa coupling matrix for each type of charged fermion should be CP and flavor violating. It was shown that the extra top Yukawa couplings, naturally \( O(1) \) in magnitude, can provide sufficient CPV needed for BAU. The specific scenario exemplified in Ref. is now excluded by ACME18, but one should explore more generic parameter space to see how one can survive the ACME18 bound.

In this Letter we find a built-in cancellation mechanism among the diagrams of Fig. that can evade the ACME18 bound, and support EWBG via top transport in g2HDM. The new bound suggests the existence of a new electron Yukawa coupling that, in conjunction with the extra top Yukawa coupling, can render eEDM sufficiently small. The cancellation mechanism works only when the hierarchical structure of the new Yukawa couplings is close to those of the SM Yukawa couplings and with particular pattern of CPV phases, which may reflect an underlying flavor structure in g2HDM.

\( g2HDM, \) EWBG, and ThO EDM.— The g2HDM extends SM by adding one extra Higgs doublet, but without imposing a \( Z_2 \) symmetry. With concerns of flavor changing neutral Higgs (FCNH) couplings controlled by fermion mass and mixing hierarchies plus alignment (i.e. close proximity of \( h(125) \) to SM Higgs boson), the phenomenological consequences of g2HDM is much

\[ \text{FIG. 1. Two-loop Barr-Zee diagrams contributing to the electron EDM, where } \phi \text{ denotes neutral and charged Higgs bosons, and } V \text{ denotes vector bosons } \gamma, Z \text{ and } W. \]
The Yukawa interactions in the mass eigenbasis are

$$-\mathcal{L}_Y = \bar{f}_i y^f_i V_R f \phi + \bar{f}_i \left[ V\rho^i_{\alpha} R - \rho^{i\dagger}_\alpha V L \right] f_i H^+ + \text{H.c.}$$

where \( f = u, d, e, f = u, v, f_i = d, e, L, R = (1 \mp \gamma_5)/2 \), \( \phi = h, H, A \) are the neutral scalars and \( H^+ \) the charged scalar, \( V \) is the CKM matrix for quarks and unit matrix for leptons. In Eq. (1), \( \rho^f_i \) are \( 3 \times 3 \) Yukawa matrices which are new sources of CPV and flavor violation, and \( y^f_i \) are related \( 3 \times 3 \) matrices with elements

\[
\begin{align*}
  y_{hi}^f &= \frac{\lambda^f_i}{\sqrt{2}} s_{ij} \gamma^i + \frac{\rho^f_i}{\sqrt{2}} c_i, \\
  y_{hi}^f &= \frac{\lambda^f_i}{\sqrt{2}} s_{ij} c_i - \frac{\rho^f_i}{\sqrt{2}} s_i, \\
  y_{hi}^f &= -i \frac{\rho^f_i}{\sqrt{2}} f_i, \quad y_{hi}^f = i \frac{\rho^f_i}{\sqrt{2}} 2
\end{align*}
\]

where \( \lambda^f_i = \sqrt{2} m_i^f / v \) (\( v = 246 \text{ GeV} \)), \( s_i = \sin \gamma \) and \( c_i = \cos \gamma \), and alignment implies \( \cos^2 \gamma \) is quite small. We will comment on later the further mixing between \( h, H, A \) induced by CPV phases of \( \rho^f_i \) at one-loop level.

As far as EWBG is concerned, not all complex phases are relevant. As found in Ref. [6], \( |\rho_{tt}| \gtrsim 0.01 \) with moderate CPV phase can generate sufficient BAU, while \( O(1) \rho_{tt} \) with maximal phase can also play a role in case \( |\rho_{tt}| \lesssim 0.01 \). Even though the \( \rho_{tt} \) mechanism is more efficient, the parameter space is severely constrained by ACME18. In the \( \rho_{tt} \) mechanism, on the other hand, in exchange for less efficient baryogenesis, it does not induce dangerous eEDMs by itself. The two mechanisms are therefore complementary. In this work, we focus exclusively on the \( \rho_{tt} \) case, and parameterize \( \rho_{tt} = |\rho_{ij}| e^{i\delta_{ij}} \).

The effective EDM for thorium monoxide (ThO) is given by [7,8]

$$d_{\text{ThO}} = d_e + \alpha_{\text{ThO}} C_2,$$

where the dimension-5 operator \( -\frac{2}{3} d_e (e\sigma^{\mu\nu} \gamma_5 e) F_{\mu\nu} \) (\( F_{\mu\nu} \) is the electromagnetic field strength tensor) induces the first term, while the second term arises from non-\( Z_2 \)-independent electron-nucleon interaction described by \( \frac{G_F}{\sqrt{2}} C_2 (\bar{N} N) (\bar{e} i \gamma_5 e) \), where \( G_F \) is the Fermi constant. ACME18 gives [2] \( d_{\text{ThO}} = (4.3 \pm 4.0) \times 10^{-30} \text{ e cm, with} \) the stated bound on \( d_e \) obtained by assuming \( C_S = 0 \). With the estimate [3] of \( \alpha_{\text{ThO}} = 1.5 \times 10^{-20} \), as we will see below, \( C_S \) cannot be completely neglected in our case, so we shall use \( d_{\text{ThO}} \) of ACME18 to explore the model.

In g2HDM, the dominant contributions to \( d_e \) come from the Barr-Zee diagrams [4], as depicted in Fig. 1 which we decompose into three pieces, depending on the particles attached to the electron line. That is,

$$d_e = d_e^{\phi\gamma} + d_e^{Z\gamma} + d_e^{W\gamma},$$

where \( \phi \) can be the neutral \( h, H, A \) bosons or the \( H^+ \) boson. CPV is violated at the lower and/or upper vertices of the \( \phi \) line. It is known that \( d_e^{\phi\gamma} \) gives the dominant contribution among the three pieces, hence the cancellation must occur in this sector. We note, however, that although \( d_e^{Z\gamma} \) and \( d_e^{W\gamma} \) are subleading, they are not always smaller than the ACME18 bound.

We further decompose each \( d_e^{\phi\gamma} \) in Eq. (4) into three types of diagrams, consisting of fermions, \( W \) and \( H^+ \) loops for \( d_e^{\phi\gamma} \) and \( d_e^{Z\gamma} \), and \( f_i/f_1, W/\phi \) and \( H^+/\phi \) loops for \( d_e^{W\gamma} \). These are denoted as \( (d_e^{\phi\gamma})_i \), \( i = f, W, H^+ \) for \( V = \gamma, Z, \) and \( (d_e^{W\gamma})_i, i = f_1/f_2, W/\phi, H^+/\phi \).

We comment on later the further mixing between \( h, H, A \) induced by CPV phases of \( \rho^f_i \) at one-loop level.

If \( \rho_{tt} \) is the only element that has nonzero CPV phase and other \( \rho \) elements are zero, one would have \( C_S = 0 \), and \( d_e \) hence \( d_{\text{ThO}} \) is solely induced by \( (d_e^{\phi\gamma})_t \), which is the left diagram of Fig. 2. We find

\[
\frac{(d_e^{\phi\gamma})_t}{e} = \frac{\alpha_{\text{em}} s_{2\gamma}}{12 \sqrt{2} \pi^2 v m_t} \frac{m_t}{m_e} \text{Im} \rho_{tt} \Delta g,
\]

\[
= -6.6 \times 10^{-29} \left( \frac{2\gamma}{0.2} \right) \left( \frac{\text{Im} \rho_{tt}}{-0.1} \right) \left( \frac{\Delta g}{0.94} \right),
\]

where \( e \) is the positron charge, \( \alpha_{\text{em}} = e^2/4\pi \) and \( \Delta g = (m_t^2/m_h^2) - (m_\gamma^2/m_h^2) \), and the loop function \( g \) is defined in Ref. [4].

In the second line of Eq. (7), we take one of the benchmark points considered in Ref. [5], i.e. \( c_\gamma = 0.1, m_b = 125 \text{ GeV and } m_H = m_A = 500 \text{ GeV, which is now excluded by ACME18. This could be circumvented by making Im} \rho_{tt} \) and/or \( c_\gamma \) small. For instance, \( |(d_e^{\phi\gamma})_t| \) would become smaller than the ACME18 bound if \( |\text{Im} \rho_{tt}| \lesssim 0.01 \), without changing the value of \( c_\gamma \). However, this would no longer be the \( \rho_{tt} \)-driven EWBG scenario [5]. For smaller \( c_\gamma \), the dependence of BAU on \( c_\gamma \) has not been studied yet in g2HDM. However, since \( c_\gamma \to 0 \) corresponds to the SM-like limit, the variation of the VEV ratio \( \Delta \beta \) during electroweak phase transition would be suppressed with decreasing \( c_\gamma \).

We conclude that the \( \rho_{tt} \)-driven EWBG case as stated above [5] is unlikely to survive the ACME18 bound.

Cancellation mechanism for ThO EDM. — In Ref. [5] we set \( \rho_{ee} = 0 \) for simplicity, but there is no symmetry or mechanism to make it zero exactly. Once complex \( \rho_{ee} \) comes in, \( (d_e^{\phi\gamma})_W \) as shown in the right diagram of Fig. 2...
can be comparable or even bigger than \((d^e_\gamma)^t\), which is analogous to \(h\) decay to diphoton. To elucidate our cancellation mechanism, we decompose \((d^e_\gamma)^t\) into two parts

\[
(d^e_\gamma)^t = (d^e_\gamma)_{\text{mix}} + (d^e_\gamma)_{\text{extr}},
\]

where the first term arises from the mixing between SM and extra Yukawa couplings, while the second term is purely from extra Yukawa couplings. For the top-loop contribution, one has

\[
\frac{(d^e_\gamma)_{\text{mix}}}{e} = \frac{\alpha_{em} s_{2\gamma}}{12\sqrt{2}\pi^3 v} \left[ \text{Im} \rho_{ee} \Delta f + \frac{m_e}{m_t} \text{Im} \rho_{HT} \Delta g \right],
\]

where \(\Delta f = X(\tau_h) - X(\tau_{1/2})\), with \(X = f, g\) defined in Refs. [4] are monotonically increasing loop functions, so \(\Delta X > 0\) for \(m_h < m_H\). For \((d^e_\gamma)_{\text{extr}}\), we take the approximation of \(\tau \ll 1\) and \(m_H \simeq m_A\).

For the \(W\)-loop contribution, on the other hand, there is no extra Yukawa coupling in the \(\phi WW\) vertex, so the \((d^e_\gamma)_{W}\) is solely given by \((d^e_\gamma)_{\text{mix}}\), which is

\[
\frac{(d^e_\gamma)_{\text{mix}}}{e} = -\frac{\alpha_{em} s_{2\gamma}}{64\sqrt{2}\pi^3 v} \text{Im} \rho_{ee} \Delta J^\gamma_W,
\]

where \(\Delta J^\gamma_W = J^\gamma_W(m_i) - J^\gamma_W(m_H)\), with \(J^\gamma_W\) defined in Ref. [10], which is a monotonically decreasing function, hence \(\Delta J^\gamma_W > 0\) for \(m_h < m_H\).

We consider the cancellation \((d^e_\gamma)_{\text{mix}} + (d^e_\gamma)_{W} = 0\), under the condition that \((d^e_\gamma)_{\text{extr}} \neq 0\). The case of having \((d^e_\gamma)_{\text{extr}} = 0\) is discussed later. From Eqs. \([9], [10] and [11]\), these two conditions lead, respectively, to

\[
\text{Im} \rho_{ee} \rho_{HT} = c \times \frac{\lambda}{\lambda_t}, \quad \text{Re} \rho_{ee} \rho_{HT} = \frac{\text{Im} \rho_{ee}}{\text{Im} \rho_{HT}}.
\]

where \(c = (16/3) \Delta g / (\Delta J^\gamma_W - 16/3 \Delta f)\). For instance, \(c \approx 0.71\) for \(m_h = 125\) GeV and \(m_H = 500\) GeV. Combining the two conditions in Eq. \([12]\), one gets \(\rho_{ee}/\rho_{HT} = c \times \lambda_t/\lambda\), with correlated phase between \(\rho_{HT}\) and \(\rho_{ee}\). Note that \(c\) is not sensitive to the exotic Higgs spectrum that is consistent with first order electroweak phase transition, hence does not change drastically in the parameter range for EWBG.

With the above cancellation, \(d^{\phi Z}_{\gamma}\), \(d^{\phi W}_{\gamma}\) and \(C_S\) become potentially important. We estimate \([11]\) \(C_S\) as

\[
C_S = -2e^2 \left[ 6.3 (C_{ue} + C_{du}) + C_{ue} \frac{41 \text{ MeV}}{m_s} \right. \\
+ C_{cc} \frac{79 \text{ MeV}}{m_c} + 0.062 \left( \frac{C_{be}}{m_b} + \frac{C_{ce}}{m_t} \right) \left( C_{ce} \right)_{_t},
\]

where \(C_{ce} = \text{Im} (\rho_{ee} \rho_{HT})\), which emerges after integrating out all neutral Higgs bosons. The quark mass suppressions are cancelled by corresponding Yukawa couplings in \(C_q\), so all quark flavors are generically relevant. Note that for \(s_t \approx 1\) and \(m_H \simeq m_A\), \(C_{qe}\) for \(u\) and \(d\)-type quarks are cast in the form of \(C_{uc} \approx \text{Im}(\rho_{ec} \rho_{HT})/(2m_n^2)\) and \(C_{de} \approx \text{Im}(\rho_{ec} \rho_{HT})/(2m_n^2)\), respectively, which implies that \(C_{qe} \approx 0\) if \((d^e_\gamma)_{\text{extr}} \neq 0\).

Before turning to numerical results, we comment on CPV effects at one-loop level, where \(h, H\) can mix with \(A\) through \(\Im \rho_{HT}\) and \(\Im \rho_{ee}\), hence are no longer CP eigenstates. The mass eigenstates are obtained by \((H_1, H_2, H_3)^T = O(h, H, A)^T\), where \(O\) is an orthogonal matrix that diagonalizes the Higgs mass squared matrix \(M_{\phi N}^2\), i.e. \(O^T M_{\phi N}^2 O = \text{diag}(m_H^2, m_h^2, m_A^2)\). The dominant contributions to the CP-mixing entries are \((M_{\phi N}^2)_{11} = -3\lambda_t \text{Im} \rho_{HT} m_f^2/4\pi^2\) and \((M_{\phi N}^2)_{13} = -3 \text{Re} \rho_{HT} m_f^2/4\pi^2\). For \(\phi_H = -90^\circ\), one finds that \(\theta_{23} \simeq \tan^{-1} \left[ 2(M_{\phi N}^2)_{13}/(m_{H_2}^2 - m_A^2) \right] / 2 \approx 9.6 \times 10^{-3}\) for \(m_H = 500\) GeV, and the effects are small enough to be ignored. For \(\phi_H \neq -90^\circ\), on the other hand, \(M_{\phi N}^2)_{23}\) being loop-induced, the 2-3 mixing angle would be \(\theta_{23} \simeq \tan^{-1} \left[ 2(M_{\phi N}^2)_{23}/(m_{H_2}^2 - m_A^2) \right] / 2 \approx 45^\circ\) if \(m_H \approx m_A\), and \(H\) and \(A\) cannot be identified as CP eigenstates at all. But even for this case, \(d_s\) would not be much affected because of the orthogonality of the matrix \(O\). For example, we estimate the relevant part for \((d^e_\gamma)^t\) as \(\sum f = O_{23} O_{31} f(m_i^2/m_{H_2}^2) \simeq O_{21} O_{31} f(m_i^2/m_{H_1}^2) + O_{22} O_{32} f(m_i^2/m_{H_2}^2) \ll 1\), where \(m_{H_2} \approx m_{H_3}\) and \(\sum O_{23} O_{31} = 0\) have been used. We conclude that the one-loop CPV effects are rather minor.

**Numerical results.** We choose \(\rho_{HT}\) to be consistent with successful EWBG, and parameterize the other diagonal \(\rho_{fi}\) elements as \(\text{Re} \rho_{fi} = a_f / \lambda / \lambda_t \text{Re} \rho_{HT}\) and \(\text{Im} \rho_{fi} = b_f / \lambda_t \text{Im} \rho_{HT}\), where \(a_f\) and \(b_f\) are real parameters such that \(|a_f| = |b_f| = r_f\). From the argument given above, the cancellation mechanism would be at work if \(a_e < 0\) and \(b_t > 0\). In what follows, we consider a flavor-blind scaling of \(a_f = -r, b_f = r\).

To see the cancellation behavior, we first investigate the magnitude of \(d_{T:\phi}\). In Fig. [3] (left) we plot \(|d_{T:\phi}|\) (black, solid) and its compositions \(|d_{T}\|\) (red, solid), \(|\alpha_{T:\phi} C_S\|\) (blue, solid), \(|d^{\phi Z}_\gamma|\) (red, dotted), \(|d^{\phi W}_\gamma|\) (red, dot-dashed) as functions of \(r\), where we set \(\text{Re} \rho_{HT} = \text{Im} \rho_{HT} = -0.1\) for illustration. The AMCE18 [2] and previous [12] (ACME14) bounds are shown as the gray and brown shaded regions as marked. The absence of \(\rho_{ee}\) would correspond to the case of \(r \approx 0\), with \(d_s \simeq (d^e_\gamma)^t\) estimated in Eq. [4]. This specific point [3] is excluded by ACME18. The situation changes considerably, however, for \(r \neq 0\).

As can be seen, strong cancellation occurs in \(d^{\phi Z}_\gamma\) around \(r \approx 0.75\). This is owing to the presence of \((d^e_\gamma)^t\), and \(d^{\phi W}_\gamma\) becomes dominant, followed by \(d^{\phi Z}_\gamma\), shifting the cancellation point in \(d_s\) upward. However, the dip
Taking Re$\rho$ the room for value of $d\phi$ in for EWBG. The gray shaded region (larger $|B|$ excluded by $\phi$ considered have $\phi$ approaches the cancellation point at $r$ with $\rho$ contours correspond to these contours are allowed, while to the right of the black solid contour, $Y_B/Y_B^{ob} = 1$, is allowed, while the gray shaded region is excluded by $B_s$-$\bar{B}_s$ mixing. Other input parameters are the same as in the left plot.

in $d_{ThO}$ moves downward due to the $C_S$ contribution. In any case, $d_{ThO}$ can be suppressed by two orders of magnitude owing to the cancellation mechanism.

We display, in Fig. 3 [right], the $2\sigma$ allowed region of $d_{ThO}$ in the ($|\rho_{tt}|$, $\phi_{tt}$) plane, taking $r = 1.0$ (blue, solid), 0.9 (red, dashed), 0.8 (magenta, dotted) and 0.75 (navy blue, dot-dashed), respectively. The region to the left of these contours are allowed, while to the right of the black contours correspond to $Y_B/Y_B^{obs} = 8.59 \times 10^{-11}$ for EWBG. The gray shaded region (larger $|\rho_{tt}|$ values) is excluded by $B_s$-$\bar{B}_s$ mixing [14]. Note that in Ref. [2], we considered $\phi_{tt} < 0$ for BAU positive. However, one can have $\phi_{tt} > 0$ by flipping the sign of $\Delta \beta$. Since the central value of $d_{ThO}$ is positive, the allowed region is asymmetric in $\phi_{tt}$. For $r = 1.0$ and 0.9, only $\phi_{tt} < 0$ is consistent with $\rho_{tt}$-driven EWBG, but $\phi_{tt} > 0$ becomes possible as $r$ approaches the cancellation point at $r \sim 0.75$, enlarging the room for $\rho_{tt}$-driven EWBG.

Let us comment on the case where $(d^\phi_{ThO})^{extr} \neq 0$. Taking $Re\rho_{ee} \simeq 0$ for illustration, we find

\begin{equation}
\frac{Im\rho_{ee}}{Im\rho_{tt}} \simeq \frac{(16/3)\Delta g}{\Delta f_{W}^f - (16/3)\Delta f + \epsilon} \simeq c' \times \frac{\lambda_e}{\lambda^f},
\end{equation}

where $s_{2\gamma}, \lambda^f = -(16/3)Re\rho_{tt}[f(\tau_{fA}) + g(\tau_{fA})]$. Thus, the coefficient $c$ in Eq. (12) can be altered by $\epsilon$, where $|c'|$ can become much larger than one when $\epsilon$ makes the denominator small. But then $d_{ThO}$ gets too large due to sizable $Im\rho_{ee}$, becoming inconsistent with ACME18. We find $|c'| \gtrsim 0.3$ for experimentally allowed $Re\rho_{tt}$. This indicates that the cancellation mechanism still suggests the $\rho$ matrices follow the SM Yukawa coupling hierarchy. It is also worth mentioning that, despite the small parameter space, further cancellation in $d_{ThO}$ can occur if we take flavor-dependent $a_f$ and $b_f$ such that $|a_f|, |b_f| < 1$. In this case, $\rho_{ob}$ could play an elevated role.

Before closing, we note that the ACME14 bound was confirmed by an independent experiment using the polar molecule $^{180}$Hf$^{19}$F$^+$ [13]. Given the significance of the ACME18 result, it should be similarly crosschecked, preferably using different methods. It is quite interesting that, while the largest diagonal extra Yukawa coupling, $\rho_{uu}$ is responsible for BAU, it works in concert with the smallest diagonal extra Yukawa coupling, $\rho_{ee}$ to generate an eEDM that might be revealed soon by very low energy, ultra-precision probes. We look forward to updates on electron EDM that may further probe the parameter space of $\rho_{tt}$-driven EWBG.

**Conclusion.**— In the scenario where an extra Yukawa coupling $\rho_{tt}$ drives EWBG, we demonstrate that the ACME18 result suggests the presence of a new electron Yukawa coupling, bringing in an exquisite cancellation mechanism for eEDM measured in ThO, which broadens the parameter space. This cancellation can be at work only when the hierarchical structure of the new Yukawa couplings is similar to those of the SM Yukawa couplings, which may reflect some underlying flavor structure in the general 2HDM. Alternatively, EWBG may be due to the weaker mechanism from flavor changing $\rho_{ee}$ coupling that evade the eEDM bound.

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