Post-ACME2013 CP-violation in Higgs Physics and Electroweak Baryogenesis (Postprint)

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Date: 2017-08-03T00:00:00+00:00

Abstract

We present a class of cancellation mechanisms to suppress the total contributions of Barr-Zee diagrams to the electron electric dipole moment (eEDM). This class of mechanisms is of particular significance after the new eEDM upper limit, which strongly constrains the allowed magnitude of CP-violation in Higgs couplings and hence the feasibility of electroweak baryogenesis (EWBG), was released by the ACME collaboration in 2013. We point out: if both the CP-odd Higgs-photon-photon (Z boson) and the CP-odd Higgs-electron-positron couplings are turned on, a cancellation may occur either between the contributions of a CP-mixing Higgs boson, with the other Higgs bosons being decoupled, or between the contributions of a CP-even and a CP-odd Higgs boson. With the assistance of the cancellation mechanisms, a large CP-phase in Higgs couplings with viable electroweak baryogenesis (EWBG) is still allowed. The reopened parameter regions would be probed by the future neutron, mercury EDM measurements, and direct measurements of Higgs CP-properties at the LHC and future colliders.

Full Text

Preamble

Post-ACME2013 CP-violation in Higgs Physics and Electroweak Baryogenesis

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We present a class of cancellation mechanisms to suppress the total contributions of Barr-Zee diagrams to the electron electric dipole moment (eEDM). These mechanisms are of particular significance after the new eEDM upper limit released by the ACME collaboration in 2013, which strongly constrains the allowed magnitude of CP-violation in Higgs couplings and hence the feasibility of electroweak baryogenesis (EWBG). We point out that if both the CP-odd Higgs-photon-photon (Z boson) and the CP-odd Higgs-electron-positron couplings are turned on, a cancellation may occur either between the contributions of a CP-mixing Higgs boson (with the other Higgs bosons being decoupled) or between the contributions of a CP-even and a CP-odd Higgs boson.

With the assistance of these cancellation mechanisms, a large CP-phase in Higgs couplings with viable electroweak baryogenesis is still allowed. The reopened parameter regions would be probed by future neutron and mercury EDM measurements, as well as direct measurements of Higgs CP-properties at the LHC and future colliders.

Introduction

The baryon asymmetry in the Universe (BAU) nowadays, i.e., $n_B/s \approx (0.7-0.9)\times 10^{-10} \neq 0$, has puzzled people for more than half a century, where s is the entropy density of the Universe. Among various dynamical mechanisms to solve this puzzle, electroweak baryogenesis (EWBG) falls in the most popular class due to its potential testability at the Large Hadron Collider and in other experiments. A generic feature of EWBG is that the CP phases employed to generate the cosmic baryon asymmetry need to enter the couplings between the Higgs sector and particles that either exist in the Standard Model (SM) or are introduced in new physics, regardless of whether the CP-phases are flavor-diagonal, off-diagonal, or flavor-decoupled. Otherwise, these CP-phases are decoupled from the electroweak phase transition (EWPT) and EWBG cannot be implemented. The measurement of Higgs CP-properties therefore provides important information for solving the BAU puzzle.

Motivated by this, the CP-properties of the Higgs boson discovered in 2012 have been extensively studied by both theorists and experimental groups since its discovery, using methods of direct measurements at the LHC. However, given the limited statistics, the sensitivity of the LHC at this stage is still low. On the other hand, fast progress has been made in indirect measurements. Using polar molecules of thorium monoxide (ThO), the ACME collaboration recently reported an upper limit on the eEDM of $|d_e| < 8.7\times 10^{-29} e$cm at 90% confidence level—an order of magnitude stronger than the previous best limit. This limit severely constrains the allowed magnitude of CP-phases in Higgs couplings via Barr-Zee diagrams, causing tension between the observation and the CP-phase required for successful EWBG.

In this letter we point out that these studies have more or less ignored a crucial effect that can dramatically change the conclusions. This is due to the fact that generally both the CP-odd Higgs-photon-photon and the CP-odd Higgs-electron-positron couplings can be or tend to be turned on simultaneously. These two couplings contribute to the eEDM separately and at the same time. If a cancellation exists between their contributions (as we will show below in two different contexts: the type-II two Higgs Doublet Model (2HDM) where the tree-level CP-phase arises from the pure Higgs sector, and the Minimal Supersymmetric Standard Model (MSSM) where the tree-level CP-phase arises from Higgs-superparticle interaction sectors), then even if the magnitudes of the CP-phases in Higgs couplings are large, the current ACME bound can be well-satisfied. In such a case, EWBG can still be successfully implemented.

General Analysis

In an effective Lagrangian for a Higgs sector, the relevant operators are given by:

$$\mathcal{L}_{\text{eff}} = \sum_i \bar{f}(c_i^f + i\tilde{c}_i^f\gamma_5)f\phi_i + \sum_i \left(c_i^\gamma F_{\mu\nu}F^{\mu\nu} + \tilde{c}_i^\gamma F_{\mu\nu}\tilde{F}^{\mu\nu}\right)\frac{\phi_i}{v}$$

where $F_{\mu\nu}$ is the field strength of the photon, with $\tilde{F}^{\mu\nu} \equiv \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}F_{\rho\sigma}$, $V_{\mu\nu}$ is the field strength of the photon and Z boson, with $\tilde{V}^{\mu\nu} \equiv \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}V_{\rho\sigma}$, and θ_i defines the CP-phase of the Yukawa couplings. These operators can be inserted in the Barr-Zee EDM diagrams. Integrating out the internal degrees of freedom, we have $\mathcal{L}_{\text{eff}} = -id_e\bar{e}\sigma_{\mu\nu}\gamma_5 e\partial^\mu A^\nu$, with its contribution to the eEDM given by:

$$d_e = \frac{em_e}{4\pi^3 v^2}\left[\sum_i \tilde{c}_i^e c_i^\gamma \log\frac{\Lambda_i^2}{m_i^2} + \tilde{c}_i^\gamma c_i^e \log\frac{\tilde{\Lambda}_i^2}{m_i^2}\right]$$

Here $v = 246$ GeV is the normalized vacuum expectation value (VEV) of the Higgs fields, and Λ_i ($\tilde{\Lambda}_i$) is the relevant scale for the $h_iF^{\mu\nu}F_{\mu\nu}$ ($h_iF^{\mu\nu}\tilde{F}_{\mu\nu}$) operator. It is clear that the Barr-Zee contributions depend not only on the CP-odd Higgs di-photon coupling $\tilde{c}_i^\gamma$, but also on the CP-even one c_i^γ if the Higgs bosons have a CP-odd coupling with electrons ($\tilde{c}_i^e \neq 0$).

The ACME measurement greatly improved the current bound on the eEDM, leading to:

$$\left|\sum_i \tilde{c}_i^e c_i^\gamma \log\frac{\Lambda_i^2}{m_i^2} + \tilde{c}_i^\gamma c_i^e \log\frac{\tilde{\Lambda}_i^2}{m_i^2}\right| < 0.14$$

This strongly constrains the allowed CP-violation in a single Higgs coupling, e.g., in the case with one (SM-like) Higgs only and meanwhile $\tilde{c}_e = 0$, unless this Higgs boson is decoupled. However, if a cancellation occurs among these interference terms, CP symmetry is allowed to be significantly violated without contradicting the current eEDM bound. Below we will use the type II 2HDM and the MSSM to show two different cancellation mechanisms, both of which are mainly motivated by EWBG: (1) cancellation occurs between the contributions of a CP-mixing Higgs boson, while the other Higgs bosons are decoupled (see the upper diagrams in Fig. 1); and (2) cancellation occurs between the contributions of a CP-even and a CP-odd Higgs boson (see the bottom diagrams in Fig. 1).

Type II 2HDM

As an illustration, we consider type II 2HDM with a softly broken Z_2 symmetry ($\phi_1 \to -\phi_1$ and $\phi_2 \to \phi_2$). Its tree-level Higgs potential is given by:

$$V = m_{11}^2(\phi_1^\dagger\phi_1) + m_{22}^2(\phi_2^\dagger\phi_2) - [m_{12}^2(\phi_1^\dagger\phi_2) + \text{h.c.}] + \frac{\lambda_1}{2}(\phi_1^\dagger\phi_1)^2 + \frac{\lambda_2}{2}(\phi_2^\dagger\phi_2)^2 + \lambda_3(\phi_1^\dagger\phi_1)(\phi_2^\dagger\phi_2) + \lambda_4(\phi_1^\dagger\phi_2)(\phi_2^\dagger\phi_1) + \Big[\frac{\lambda}{2}$$

where m_{12}^2 and λ_5 are complex parameters. Their relative phase $\text{Arg}(\lambda_5 m_{12}^{4*})$ leads to CP violation in the Higgs sector. We take the convention that both Higgs doublets $\phi_{1,2}$ carry a hypercharge of one unit and that the general Higgs VEVs are:

$$\langle\phi_1\rangle = \begin{pmatrix} 0 \\ v_1 \end{pmatrix}, \quad \langle\phi_2\rangle = \begin{pmatrix} 0 \\ v_2 e^{i\xi} \end{pmatrix}$$

with $\sin^2\beta = |v_2|^2/(|v_1|^2 + |v_2|^2)$, $v_1 = v\cos\beta$, $|v_2| = v\sin\beta$. The unitary matrix R, defined to diagonalize the Higgs mass matrix M^2 and satisfy $RM^2R^T = \text{diag}(M_{h_1}^2, M_{h_2}^2, M_{h_3}^2)$ in the mass eigenstate basis (h_1, h_2, h_3), can be easily figured out and is given by:

$$R = \begin{pmatrix} -c_\alpha c_{\alpha_b} & s_\alpha c_{\alpha_b} & s_{\alpha_b} \\ s_\alpha c_{\alpha_c} - c_\alpha s_{\alpha_b} s_{\alpha_c} & c_\alpha c_{\alpha_c} + s_\alpha s_{\alpha_b} s_{\alpha_c} & -c_{\alpha_b} s_{\alpha_c} \\ s_\alpha s_{\alpha_c} + c_\alpha s_{\alpha_b} c_{\alpha_c} & c_\alpha s_{\alpha_c} - s_\alpha s_{\alpha_b} c_{\alpha_c} & c_{\alpha_b} c_{\alpha_c} \end{pmatrix}$$

with $c_i = \cos\alpha_i$, $s_i = \sin\alpha_i$. Here α, α_b, α_c are mixing angles between two CP-even Higgs bosons, between the light CP-even and CP-odd Higgs bosons, and between the heavy CP-even and CP-odd Higgs bosons, respectively. The angular range, beyond which R is repeated, can be chosen as $0 < \alpha \le \pi$, $-\pi < \alpha_b \le \pi$ and $-\pi/2 < \alpha_c \le \pi/2$.

The tree-level h_1 couplings rescaled by the SM values are given by:

$$c_b = c_e = -\frac{\sin\alpha\cos\alpha_b}{\cos\beta}, \quad \tilde{c}_t = -\cot\beta\sin\alpha_b, \quad a_V = \cos\alpha_b\sin(\beta - \alpha)$$

$$\tilde{c}_b = \tilde{c}_e = -\tan\beta\sin\alpha_b$$

where h_1 is SM-like and a_V represents the h_1WW and h_1ZZ couplings. The CP-phase of the top Yukawa coupling θ_t is given by $\tan\theta_t = -\cos\beta\cos\alpha\tan\alpha_b$.

These tree-level effective couplings further contribute to c_γ and $\tilde{c}_\gamma$ at loop level:

$$c_\gamma = \sum_f \frac{Q_f^2 c_f}{2} = \frac{2c_t}{9} - \frac{7a_V}{8}$$

$$\tilde{c}_\gamma = \sum_f \frac{Q_f^2 \tilde{c}_f}{4} = -\frac{\tilde{c}_t}{3}$$

Note that the signs of c_i^γ are derived in the convention of Eq. (3). To achieve the cancellation indicated by Eq. (5), we require $c_{t,W}^\gamma c_e \sim \tilde{c}_i^\gamma \tilde{c}_e/\tilde{c}_t$ or $\tan\beta \sim 1$.

In this setup, the lightest Higgs boson h_1 leads to a leading-order contribution to the eEDM via the Barr-Zee diagrams:

$$d_e^{h_1\gamma\gamma} = \frac{2\alpha G_F m_e}{(4\pi)^3}\left[3f(z_t)\tilde{c}_e c_t + g(z_t)\tilde{c}_t c_e\right]$$

$$d_e^{h_1\gamma Z} = \frac{2\alpha G_F m_e}{(4\pi)^3}\left[3f(z_W) + 5g(z_W)\right] a_V \tilde{c}_e$$

where $z_t = m_t^2/m_{h_1}^2$, $z_W = m_W^2/m_{h_1}^2$, and the loop functions $f(z)$ and $g(z)$ are given in [**?**]. Numerically, we have $f(z_t) = 1.0$ and $g(z_t) = 1.4$. These quantities depend on three free parameters α, α_b and β. For simplicity we work in the alignment limit $\beta = \alpha + \pi/2$, where the free parameters are reduced to β, α_b, with $\tan\theta_t = -\cot\beta\tan\alpha_b$, and the 125 GeV Higgs boson is SM-like if there is no CP-violation. The overall contribution to the eEDM is:

$$d_e^{h_1\gamma\gamma} = \frac{2\alpha G_F m_e}{(4\pi)^3}\sin\alpha_b\cos\alpha_b\left[f'(z_t, \tan\beta) - g'(z_W, \tan\beta)\right]$$

with $f'(z,x) = -8(xf(z) + g(z)/x)/3$ and $g'(z,x) = (3f(z) + 5g(z))x$. The contributions from neutral Higgs with the Z gauge boson and charged Higgs with W gauge boson as propagators are generally smaller, so we neglect them in the calculation.

The fitting results are presented in Fig. 2, where the inclusive LHC data published in March 2013 and the most recent ACME results are applied. In the presence of CP-even and CP-odd Higgs mixing, both the current ACME constraints and the Higgs global fits favor the region with $\tan\beta \sim 1$, where we have $\tan\theta_t \sim -\tan\alpha_b$. The former is easy to understand since a cancellation between $f'(z_t, \tan\beta)$ and $g'(z_W, \tan\beta)$ in Eq. (15) requires $\tan\beta = 1.04$. The latter is because a relatively small $\tan\beta$ can help avoid too large a signal rate of $h \to b\bar{b}$, and hence an over-suppressed $h \to \gamma\gamma$ rate. In this cancellation region, a CP-violation effect with $|\tan\alpha_b|$, $|\tan\theta_t| > 0.1$ is allowed while the most stringent constraints are from the nEDM.

MSSM

Though the MSSM is of type II 2HDM, there is no tree-level CP-violation in the Higgs sector either explicitly or spontaneously, due to a vanishing λ_5 term in the Higgs potential. The CP-phases used for EWBG mostly arise in the tree-level superparticle sectors, such as the chargino, neutralino, squark and slepton sectors. The explicit CP-violation in these sectors can break CP symmetry in the Higgs sector at loop level, leading to CP-even and CP-odd mixing terms in the Higgs squared mass matrix. However, the Higgs CP-mixture caused by this effect is small due to loop suppression. For nonstandard Higgs bosons we notice that the CP-mixture is typically below 10%, consistent with previous studies, even if the CP-phase arises from the stop sector, while for the SM-like Higgs boson, the CP-mixture is suppressed further by an extra $\tan\beta$ factor. So the Higgs eigenstates are approximately CP-eigenstates, with their couplings to electrons being either $|c_e| \gg |\tilde{c}_e|$ (for CP-even Higgs bosons) or $|c_e| \ll |\tilde{c}_e|$ (for CP-odd Higgs bosons). On the other hand, a relatively large $\tan\beta$ is favored in the MSSM, given that the tree-level mass of the SM-like Higgs boson is larger in this case. This leads to $|c_h| \ll |c_{H,A}|$, and hence a small contribution from h to the eEDM (mainly via the $c_e\tilde{c}_\gamma$ term). So in the MSSM with EWBG implemented, the main contributions to the eEDM are made by nonstandard Higgs bosons unless they are highly decoupled.

Among these CP-violating sources, the one arising in the chargino sector is of particular interest because of its high efficiency in generating the BAU via EWBG. The charginos enter the $H\gamma\gamma$ and $A\gamma\gamma$ loops as new mediators, inducing non-trivial contributions to the eEDM via the $c_e\tilde{c}_\gamma$ and $\tilde{c}_e c_\gamma$ terms, respectively. Though these two contributions are comparable in magnitude, their signs are typically different. Given an extra minus sign for the term c_i^γ in Eq. (5), this scenario is strongly constrained by the ACME eEDM bound. However, recall that charged particles like stau leptons also enter the $A\gamma\gamma$ loop. If a non-trivial CP-phase is turned on in the stau sector, new contributions to the eEDM will be introduced via the $\tilde{c}_e c_\gamma$ term. (Note that such a CP-violating coupling will not induce non-trivial contributions to the eEDM via the $H\gamma\gamma$ loop or the $c_e\tilde{c}_\gamma$ term, since stau leptons are scalar particles.) This provides a potential cancellation mechanism such that a CP-phase in the Higgs-chargino couplings large enough for implementing the EWBG mechanism remains allowed.

With CP-violation turned on in the stau sector, the main contributions to the eEDM are then given by:

$$d_e \approx C\tilde{c}_A \sum_{j=1,2} \tilde{c}^{\gamma}_{\tilde{\chi}^{\pm}_j} \log \frac{m^2_{\tilde{\chi}^{\pm}_j}}{m^2_A} + C\tilde{c}_A \sum_{i=1,2} \tilde{c}^{\gamma}_{\tilde{\tau}_i} \log \frac{m^2_{\tilde{\tau}_i}}{m^2_A}$$

with the terms in the first and second lines mediated by A and H, respectively. Here $C = \alpha m_e/(4\pi^3 v^2)$, and:

$$\tilde{c}^{\gamma}_{\tilde{\chi}^{\pm}} = -\sum_{i=1,2} \frac{Q^2_{\tilde{\chi}^{\pm}_i} \tilde{c}_{\tilde{\chi}^{\pm}_i}}{2}, \quad \tilde{c}^{\gamma}_{\tilde{\tau}_i} = -\sum_{i=1,2} \frac{Q^2_{\tilde{\tau}_i} \tilde{c}_{\tilde{\tau}_i}}{4}$$

As for $m_{\tilde{\chi}^{\pm}}$ and $\tilde{c}_{\tilde{\chi}^{\pm}}$ (CP-odd interaction between H and charginos), we refer to the literature for details.

Fig. 3 depicts all low-energy experimental constraints calculated by CPsuperH. Note, however, that the CPsuperH codes used in the analysis were revised by the authors of this letter (for details, see Appendix C), in which we corrected a sign error in the anomalous dimension of the dipole operators, updated the QCD hadron matrix elements, incorporated the mixing effects between the dipole and pseudo-scalar operators during renormalization group running, and added the previously missed contributions to the eEDM from the W-mediated $h\gamma\gamma$ loop in the Barr-Zee diagrams. With these corrections, we notice that the neutron and mercury bounds may become substantially weaker than those given by the original CPsuperH.

As indicated in Fig. 3, the charginos have a negative contribution to the eEDM (black-dashed contours) by coupling with both A and H. The staus, on the other hand, have a positive contribution to the eEDM (blue-dashed contours) by mainly coupling with A. Both contributions are enhanced by $\tan\beta$ because of $\tilde{c}^e_A \propto 1/\cos\beta$. Their dependences on the μ parameter, however, are different. For charginos, the H/A-chargino-chargino couplings $g_P \propto -i(C_R)_{i1}(C_L)^*_{i2}/2 -$ h.c., $i = 1, 2$. With $\mu > M_2$ assumed here, a small μ value will increase the off-diagonal terms in the chargino mixing matrices, $(C_L)_{12}$ and $(C_R)_{21}$, and hence the overall eEDM contribution. As for staus, a larger μ leads to a lighter $\tilde{\tau}_1$ because of a larger mixing term $\propto |\mu \tan\beta - A_\tau|$. Therefore, their contribution to the eEDM increases for larger μ values. Due to their cancellation, there exists a blank region in Fig. 3 where the total eEDM is below the current ACME bound. This region overlaps with the EWBG-favored region that would have been excluded by the ACME bound if only the chargino contribution were taken into account. One benchmark point is presented in Table II.

In the MSSM, a large $\tan\beta$ is favored by the LHC Higgs bounds, and the contributions of the Barr-Zee diagrams to the eEDM are mainly mediated by nonstandard CP-even and CP-odd neutral Higgs bosons (if they are not as heavy as the 10 TeV scale or above). With a cancellation between them, a maximal CP-phase in their couplings with superparticles like charginos is still allowed.

In both cases, successful EWBG is possible, which turns out to be challenging without the assistance of these cancellation mechanisms.

The EWPT, another key element for EWBG, is not addressed in this letter since it is "orthogonal" to the discussions above. It is straightforward to generalize the discussions on eEDM cancellation to the 2HDM plus a singlet, or MSSM extensions with an extra singlet superfield or gauge group, where a strong enough EWPT is not difficult to achieve. Furthermore, in the MSSM, while an extremely light right-handed stop scenario ($\sim$ 110 GeV) has been ruled out by Higgs global fits and direct searches, a moderately light stop with other light scalar particles (like sbottom or stau) may still be viable. Detailed discussions in this regard will be left for future studies.

Acknowledgments

We would like to thank Michael Ramsey-Musolf, Carlos Wagner, and Yue Zhang for useful discussions, and David McKeen for collaboration in the early stages of this project. TL is supported by his start-up fund at the Hong Kong University of Science and Technology. We also acknowledge the hospitality of the Kavli Institute for Theoretical Physics and the Aspen Center for Physics (Simons Foundation), where part of this work was completed.

Discussion and Conclusion

Neutron and Mercury EDMs in the MSSM: CPsuperH vs. Our Analysis

In this letter, we present a class of cancellation mechanisms to suppress the total contributions of Barr-Zee diagrams to the eEDM, which may occur either between the contributions of a CP-mixing Higgs boson (with the other Higgs bosons decoupled) or between the contributions of a CP-even and a CP-odd Higgs boson.

As an illustration, we study two scenarios: the type-II 2HDM where the tree-level CP-phase arises from the Higgs sector, and the MSSM where the tree-level CP-phase arises from Higgs-superparticle interaction sectors. In the 2HDM, $\tan\beta \sim 1$ is favored by the LHC Higgs bounds, and the contributions of the Barr-Zee diagrams to the eEDM are mainly mediated by a CP-mixed SM-like Higgs boson. With a cancellation between them, a CP-phase as large as $\mathcal{O}(0.1-1)$ is still allowed for the top Yukawa coupling, induced by Higgs CP mixing.

In the MSSM, a large $\tan\beta$ is favored by the LHC Higgs bounds, and the contributions of the Barr-Zee diagrams to the eEDM are mainly mediated by nonstandard CP-even and CP-odd neutral Higgs bosons (if they are not as heavy as the 10 TeV scale or above). With a cancellation between them, a maximal CP-phase in their couplings with superparticles like charginos is still allowed. In both cases, successful EWBG is possible, which turns out to be challenging without the assistance of these cancellation mechanisms.

In the Appendix, we introduce the modifications and corrections made to CPsuperH in our analysis. We then compare the constraints on neutron and mercury EDMs given by the original CPsuperH and the modified version. To begin, we briefly introduce the theoretical methods used for calculating neutron and mercury EDMs, basically following refs [?, ?].

(1) Renormalization Group Running of Wilson Coefficients The EDM of a fermion (d_f^E; electron EDM is an exception denoted by d_e in this letter), the chromo-EDM of a quark (d_q^C), and the Weinberg operator coupling (d^G) are defined by the Lagrangian:

$$\mathcal{L}_{\rm EDM} = -\frac{i}{2} d_f^E F_{\mu\nu} \bar{f} \sigma^{\mu\nu} \gamma_5 f - \frac{i}{2} d_q^C g_s G^a_{\mu\nu} \bar{q} \sigma^{\mu\nu} \gamma_5 T^a q + \frac{1}{6} d^G g_s f^{abc} \epsilon^{\mu\nu\lambda\sigma} G^a_{\mu\nu} G^b_{\lambda\rho} G_\sigma{}^c$$

These can be generated through CP-violating neutral Higgs-boson mixing in the t-channel and CP-violating Yukawa threshold corrections. The corresponding CP-odd coefficients are given by:

$$\delta_f = -\frac{2eQ_f m_f}{\Lambda^2} c_f \tilde{c}_f, \quad \tilde{\delta}_q = -\frac{2g_s m_q}{\Lambda^2} c_q \tilde{c}_q, \quad C_{\tilde{G}} = -\frac{3g_s}{\Lambda^2} \sum_q c_q \tilde{c}_q$$

where m_f and Λ denote the fermion masses and the MSSM scale, respectively. The Lagrangian can be rewritten as:

$$\mathcal{L}_{\rm EDM} = i \sum_f \frac{m_f Q_f e}{\Lambda^2} \delta_f F_{\mu\nu} \bar{f} \sigma^{\mu\nu} \gamma_5 f + i \sum_q \frac{m_q g_s}{\Lambda^2} \tilde{\delta}_q G^a_{\mu\nu} \bar{q} \sigma^{\mu\nu} \gamma_5 T^a q - \frac{3g_s}{\Lambda^2} C_{\tilde{G}} f^{abc} \epsilon^{\mu\nu\lambda\sigma} G^a_{\mu\nu} G^b_{\lambda\rho} G_\sigma{}^c$$

with $g_{q(q')} = m_{q(q')}/v$ and $v = 2M_W/g$. The last CP-odd four-fermion operator $\mathcal{O}^4_{q'q} = \bar{q}'_\alpha \sigma_{\mu\nu} q'_\beta \bar{q}_\beta i\sigma^{\mu\nu}\gamma_5 q_\alpha$, on the other hand, is generated from operator mixing effects of $C^1_{q'q}$ which follow Eq. (27) below.

To calculate the Wilson coefficients δ_q, $\tilde{\delta}_q$, and $C_{\tilde{G}}$ at a GeV scale, we need to evolve them from the MSSM scale Λ down to the GeV scale based on the Renormalization Group Equations (RGE) [?, ?]:

$$\frac{d}{d\ln\mu} \begin{pmatrix} \delta_q \\ \tilde{\delta}_q \\ C_{\tilde{G}} \\ C^1_{q'q} \\ C^1_{qq'} \\ C^4_{q'q} \end{pmatrix} = \gamma \begin{pmatrix} \delta_q \\ \tilde{\delta}_q \\ C_{\tilde{G}} \\ C^1_{q'q} \\ C^1_{qq'} \\ C^4_{q'q} \end{pmatrix}$$

The one-loop anomalous dimension matrix is given by the standard results in the literature.

During this process, flavor-conserving CP-odd four-fermion operators may lead to non-trivial corrections to the Wilson coefficients of the CEDM and Weinberg operators via mixing. A more complete Lagrangian for calculating neutron and mercury EDMs should include these operators.

Besides the RG running and operator mixing effects, there are two additional important bottom quark threshold effects. First, the Weinberg operator receives a shift from the bottom quark CEDM after the bottom quark is integrated out at m_b:

$$\Delta C_{\tilde{G}}(m_b) = \frac{\alpha_S(m_b)}{4\pi}\tilde{\delta}_b(m_b)$$

where $\tilde{\delta}_b(m_b) = \tilde{\delta}_b^0(m_b) + \Delta\tilde{\delta}_b(m_b)$ is the b-quark CEDM at the m_b scale, with $\tilde{\delta}_b^0(m_b)$ being the direct CEDM from two-loop Barr-Zee graphs. The $\Delta\tilde{\delta}_b(m_b)$ is the bottom quark CEDM correction from RG running and operator mixing. Keeping the leading logarithmic terms that contribute to $\Delta\tilde{\delta}_b(m_b)$ at the matching scale $\mu = m_b$, we have:

$$\Delta\tilde{\delta}_b(m_b) \approx \frac{8\pi^2}{g_s^2} C_{bq}^1(m_H) \log\frac{m_H^2}{m_b^2}$$

which could be figured out from the RGE and arises from integrating out the b quark at one-loop level.

The second important correction from integrating out the bottom quark is the shift to CEDMs of other quarks:

$$\Delta\tilde{\delta}_q(m_b) \approx \frac{\pi}{g_s^2}\left(\tilde{C}_{bq}^1(m_H) + \tilde{C}_{qb}^1(m_H)\right)\left(\frac{m_b}{m_q}\right)^2$$

which are actually two-loop effects where the operator $\tilde{C}_{bq}^1$ first mixes into $\mathcal{O}_{qq'}^4$ and then mixes into $\Delta\tilde{\delta}_q(m_b)$. Numerically, the suppression from one additional loop is compensated by the m_b/m_q enhancement.

(2) Neutron and Mercury EDMs The neutron EDM is calculated by:

$$d_n = e\zeta_n^u\delta_u + e\zeta_n^d\delta_d + e\tilde{\zeta}_n^u\tilde{\delta}_u + e\tilde{\zeta}_n^d\tilde{\delta}_d + \beta_n^G C_{\tilde{G}}$$

In our analysis, we updated hadronic matrix elements with $\zeta_n^u = 0.82\times10^{-8}$, $\zeta_n^d = 1.63\times10^{-8}$, $\tilde{\zeta}_n^u = -3.3\times10^{-8}$, $\tilde{\zeta}_n^d = 0.82\times10^{-8}$, and $\beta_n^G = 2\times10^{-20}$ e cm.

Compared to the electron EDM d_e and CP-odd electron-nucleon interactions, the nuclear Schiff moment has a larger contribution to the mercury EDM. The Schiff moment is generated by long-range, pion-exchange mediated P- and T-violating nucleon-nucleon interactions:

$$\mathcal{L}_{\pi NN} = \bar{N}\left[\bar{g}_\pi^{(0)}\vec{\tau}\cdot\vec{\pi} + \bar{g}_\pi^{(1)}\pi^0 + \bar{g}_\pi^{(2)}(2\tau_3\pi^0 - \vec{\tau}\cdot\vec{\pi})\right]N$$

In a general context, the isoscalar and isovector couplings $\bar{g}_\pi^{(0)}$ and $\bar{g}_\pi^{(1)}$ dominate over the isotensor coupling $\bar{g}_\pi^{(2)}$, so the mercury EDM is approximately given by:

$$d_{\rm Hg} \approx \kappa_S\left(2m_N g_A a_0 \bar{g}_\pi^{(0)} + a_1\bar{g}_\pi^{(1)}\right)$$

with $g_A \approx 1.26$, $F_\pi = 186$ MeV. To calculate the $\bar{g}_\pi^{(i)}$ couplings, we need the Wilson coefficients $\{C_{\tilde{G}}, \tilde{\delta}_u, \tilde{\delta}_d\}$ at the hadron scale. In our analysis, we updated hadronic matrix elements with $\tilde{\eta}^{(0)} = -2\times10^{-7}$, $\tilde{\eta}^{(1)} = -4\times10^{-7}$, $\gamma_{\tilde{G}}^{(0)} \approx \gamma_{\tilde{G}}^{(1)} = 2\times10^{-6}$; updated nuclear matrix elements with $a_0 = 0.01$ e fm³, $a_1 = \pm0.02$ e fm³; and assumed a new atomic sensitivity coefficient $\kappa_S = -2.8\times10^{-4}$ fm ². We emphasize again that both the sign correction of the anomalous dimension coefficient γ_e and the mixing effect of the RGE operators have been incorporated in calculating $\{\delta_u, \delta_d, \tilde{\delta}_u, \tilde{\delta}_d, C_{\tilde{G}}\}$ at the hadron scale in our analysis, which may cause an important change in the bounds of both neutron and mercury EDMs compared to the original CPsuperH.

(3) The CPsuperH vs. Our Analysis In CPsuperH, we made the following corrections and modifications:

1. **Sign error in anomalous-dimension coefficient**: A sign error in γ_e was fixed, which corresponds to the (1,1) component of the anomalous dimension matrix. CPsuperH follows Ref. [**?**] where $\gamma_e = 8/3$ is defined. As pointed out in Refs. [**?**, **?**], however, γ_e should be $-8/3$ rather than $8/3$ in that notation.
2. **Mixing effects**: The mixing effects between the Weinberg gluonic operator, quark color dipole operator, quark electric dipole operator, and CP-odd four-fermion operators have been incorporated.
3. **Barr-Zee diagrams with Higgs-photon-photon(Z) loop mediated by W boson**: These contributions have been added.

These corrections and modifications to CPsuperH may significantly change the neutron and mercury EDM bounds for CP-violation in the MSSM. As a comparison, the neutron and mercury EDM bounds after and before the revision of CPsuperH are shown in Fig. 3 and Fig. 4 [Figure 4: see original paper], respectively, with a chiral model assumed for hadrons and the same parameter values in both plots. The comparison indicates that both the neutron and mercury EDM bounds are overestimated by the original CPsuperH for the assumed parameter values.

Note: Figure translations are in progress. See original paper for figures.