Anomalous B meson mixing and baryogenesis
in a two Higgs doublet model with top-charm flavor violation

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There exist experimental hints from the $B$ sector for CP violation beyond the Standard Model (SM) CKM paradigm. An anomalous dimuon asymmetry was reported by the D0 collaboration, while tension exists between $B \to \tau\nu$ and $S_{\tau K}$. These measurements, disfavoring the SM at the $\sim 3\sigma$ level, can be explained by new physics in both $B^0_d$-$\bar{B}^0_d$ and $B^0_s$-$\bar{B}^0_s$ mixing, arising from (1) new bosonic degrees of freedom at or near the electroweak scale, and (2) new, large CP-violating phases. These two new physics ingredients are precisely what is required for electroweak baryogenesis to work in an extension of the SM. We show that a simple two Higgs doublet model with top-charm flavor violation can explain the $B$ anomalies and the baryon asymmetry of the Universe. Moreover, the presence of a large relative phase in the top-charm Yukawa coupling, favored by $B^0_d$-$\bar{B}^0_d$ mixing, weakens constraints from $\epsilon_K$ and $b \to s\gamma$, allowing for a light charged Higgs mass of $\mathcal{O}(100 \, \text{GeV})$.

I. INTRODUCTION

Precision tests of CP violation have shown a remarkable consistency with the Standard Model (SM), where all CP-violating observables are governed uniquely by the single phase in the Cabibbo-Kobayashi-Maskawa (CKM) matrix \textsuperscript{[1]}. Yet the search continues. Many well-motivated extensions of the SM, such as supersymmetry, contain new sources of CP violation at the electroweak scale. Furthermore, new CP violation beyond the CKM phase is likely required to explain the origin of the baryon asymmetry of the Universe.

Recent analyses have suggested that the CKM paradigm may be in trouble. First, the D0 collaboration has measured the like-sign dimuon asymmetry, arising from CP violation in the mixing and decays of $B^0_d,s$ mesons, in excess over SM prediction at the $3.2\sigma$ level \textsuperscript{[2]}. Second, there is tension at the $\sim 3\sigma$ level between the branching ratio for $B^+ \to \tau^+\nu$ and the CP asymmetry $S_{\tau K}$ in $B^0_d \to J/\psi K$ \textsuperscript{[3, 4]}. Additionally, CDF and D0 have measured the CP asymmetry $S_{\psi\phi}$ in $B^0_s \to J/\psi\phi$. While their earlier results (each with $2.8 \, \text{fb}^{-1}$ data) showed a $\sim 2\sigma$ deviation from the SM \textsuperscript{[3, 4]}, this discrepancy has been reduced in their updated analyses with more data ($5.2$ and $6.1 \, \text{fb}^{-1}$, respectively) \textsuperscript{[5]}. Although further experimental study is required, taken at face value, these anomalies suggest CP violation from new physics (NP) in the mixing and/or decay amplitudes of $B^0_d,s$. Recently, the CKMFitter group has performed a global fit to all flavor observables, allowing for arbitrary new physics in $B^0_d,s$-$\bar{B}^0_d,s$ mixing amplitudes \textsuperscript{[6]}. They conclude that the SM is disfavored at $3.4\sigma$, while the data seem to favor NP with large CP-violating phases relative to the SM in both $B^0_d$ and $B^0_s$ mixing. At the level of effective theory, this NP takes the form

$$\mathcal{L}_{\text{NP}} \sim \frac{c_d}{\Lambda^2} (bd)^2 + \frac{c_s}{\Lambda^2} (bs)^2 + \text{h.c.}$$

These operators can arise from new bosonic degrees of freedom at or near the weak scale, with new large CP-violating phases \textsuperscript{[6, 7]}. It is suggestive that the same NP ingredients, new weak-scale bosons and new CP violation, can also lead to successful electroweak baryogenesis (EWBG). EWBG, in which the baryon asymmetry is generated during the electroweak phase transition \textsuperscript{[13, 15]}, is particularly attractive since two out of three Sakharov conditions \textsuperscript{[16]} can be tested experimentally. First, a departure from thermal equilibrium is provided by a strong first-order phase transition, proceeding by bubble nucleation. While this does not occur in the SM \textsuperscript{[17]}, additional weak-scale bosonic degrees of freedom can induce the required phase transition; these new bosons can be searched for at colliders. Second, there must exist new CP violation beyond the SM \textsuperscript{[18]}.

If we wish to connect Eq. \textsuperscript{[11]} to EWBG, it is better to generate these operators at one-loop, rather than tree-level. Constraints on the mass differences $\Delta M_{d,s}$ in the $B^0_d,s$ systems require that $\Lambda^2/|c_d| \gtrsim (500 \, \text{TeV})^2$ and $\Lambda^2/|c_s| \gtrsim (100 \, \text{TeV})^2$ \textsuperscript{[21]}. For tree-level exchange, it seems unlikely that all three Sakharov conditions can be met at once. Sufficient baryon number generation typically requires couplings $\gtrsim \mathcal{O}(10^{-1})$, such that $c_{d,s} \gtrsim \mathcal{O}(10^{-2})$, while a viable phase transition requires $\Lambda \lesssim 1 \, \text{TeV}$. Therefore, EWBG requires $\Lambda^2/|c_{d,s}| \lesssim (10 \, \text{TeV})^2$, at odds with $\Delta M_{d,s}$ constraints. However, if the operators in Eq. \textsuperscript{[11]} arise at one-loop order, $c_{d,s}$ will have an...
additional $1/(4\pi^2)$ loop suppression, allowing for both large couplings and lighter scale $\Lambda$, without conflicting with $\Delta M_{ls}$ constraints.

In this work, we propose that a simple two Higgs doublet model (2HDM) can account for both anomalous CP violation in $B^0_{d,s} \to \bar{B}^0_{d,s}$ mixing and EWBG. Previous works have studied CP violation in $B^0_{d,s} \to \bar{B}^0_{d,s}$ mixing within a 2HDM [3,12]. Our setup, described in Sec. III is different: we assume the NP Higgs doublet $(H^+, H^0+iA^0)$ mediates top-charm flavor violation. In this case, the NP $B^0_{d,s} \to \bar{B}^0_{d,s}$ mixing amplitudes $(M^d_{12})_{NP}$ are generated at one-loop order through charge current interactions mediated by $H^+$ (similar to Ref. [12]), rather than through tree-level exchange [8,11]. In Sec. III we compute $(M^d_{12})_{NP}$ in our model. We find:

- The best fit values to both $M^d_{12}$ and $M^s_{12}$, from Ref. [8], can be explained in terms of a single NP phase $\theta_{tc}$ (defined below).

- For large values of $\theta_{tc}$ preferred by $B^0_{d,s} \to \bar{B}^0_{d,s}$ mixing observables, constraints from $\epsilon_K$ and $b \to s\gamma$ are weakened and $H^\pm$ can be light ($m_{H^\pm} \sim 100$ GeV).

In Sec. IV we discuss in detail EWBG in our 2HDM model. We focus on the CP violation aspects of EWBG, computing the baryon asymmetry in terms of the underlying parameters of our model by solving a system of coupled Boltzmann equations. We find:

$\mathcal{L}_{yuk} \supset \bar{u}_L (y_U H_1 + \bar{y}_U H_2) Q_L + \text{h.c.}$

where the left-handed quark doublet is $Q_L \equiv (u_L, V d_L)$. The SU(2)$_L$ contraction is $H_i Q_L = H^+_i (V d_L) - H^0_i u_L$. The 3 $\times$ 3 Yukawa matrices $y_U$ and $\bar{y}_U$ couple right-handed $u$-type quarks $u_R \equiv (u, c, t)L$ and left-handed quarks $u_L \equiv (u, c, t)L$ and $d_L \equiv (d, s, b)L$. Working in the mass eigenstate basis, the matrix

$y_U = \text{diag}(y_u, y_c, y_t) = \text{diag}(m_u, m_c, m_t)/v$

is a diagonal matrix of SM Yukawas, and $V$ is the CKM matrix. Analogous Yukawa couplings arise for down quarks and charged leptons:

$\mathcal{L}_{yuk} \supset -\tilde{d}_R(y_D H^+_1 + \bar{y}_D H^+_2) Q_L - \tilde{e}_R(y_L H^+_1 + \bar{y}_L H^+_2) L_L + \text{h.c.}$

where $y_D = \text{diag}(y_d, y_s, y_b)$ and $y_L = \text{diag}(y_e, y_\mu, y_\tau)$ are the SM Yukawas.

The NP Yukawa matrices $\tilde{y}_U, D, L$ can be arbitrary. However, the absence of anomalously large flavor-violating processes provides strong motivation for an organizing principle. In this work, we assume that flavor violation arises predominantly in the top sector. Specifically, we take

$\tilde{y}_U = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & \tilde{y}_{tc} & \tilde{y}_{tc} \end{pmatrix}, \quad \tilde{y}_{D,L} = 0.$

That is, we consider a hierarchical structure where the $t_Rt_L$ and $t_Rc_L$ couplings are dominant (with $|\tilde{y}_{tc}| \gg |\tilde{y}_{tc}|$), while others are suppressed. The zeros in Eq. (6) are meant to indicate these subleading couplings that for simplicity we neglect in our analysis. In our setup, flavor violation in meson observables arises at one-loop order through $H^\pm$ charge current interactions, discussed in the next section.

III. FLAVOR CONSTRAINTS

Mixing and CP violation in the $B^0_q \to \bar{B}^0_q$ system ($q = d, s$) is governed by the off-diagonal matrix element $M^q_{12} = i \Gamma_{12}^q$ in the Hamiltonian [24,25], with $M^q_{12}$ ($\Gamma_{12}^q$) associated with the (anti-)Hermitian part. Only the relative phase $\phi_q \equiv \arg(-M^q_{12}/\Gamma_{12}^q)$ is physical. The relevant observables are the mass and width differences between the two eigenstates

$\Delta M_q = 2|M^q_{12}|, \quad \Delta \Gamma_q = 2|\Gamma_{12}^q| \cos \phi_q,$

and the wrong sign semileptonic asymmetry

$a_{sL}^q = \frac{\Gamma(B^0_q \to \mu^+ X) - \Gamma(B^0_q \to \mu^- X)}{\Gamma(B^0_q \to \mu^+ X) + \Gamma(B^0_q \to \mu^- X)} \approx \frac{|\Gamma_{12}^q|}{|M_{12}^q|} \sin \phi_q.$

The dimuon asymmetry measured by D0 arises from wrong sign semileptonic decays of both $B^0_d$ and $B^0_s$ mesons and is given by $A_{sL} \approx 0.5 a_{sL} + 0.5 a_{sT}^d$ [2].
FIG. 1: New physics $B_0^0$-$B_d^0$ and $B_s^0$-$B_{d,s}^0$ mixing amplitudes $(M_{12}^d)$$_{NP}$ arising from box graphs with $H^\pm$ exchange.

In the SM, the mixing amplitude $M_{12}^d$ arises from box graphs, while the $f^d_{12}$ comes from tree-level decays. Therefore, it is plausible that NP effects enter predominantly through mixing. Deviations in $M_{12}^d$ from the SM can be parametrized by

$$M_{12}^d = (M_{12}^d)_{SM} + (M_{12}^d)_{NP} \equiv (M_{12}^d)_{SM} \Delta q . \quad (9)$$

The consistency of $\Delta M_{d,s}$ with SM predictions constrains $|\Delta q/d,s| \approx 1$, at the O(20%) level [8], while the dimuon asymmetry measurement disagrees with SM prediction at 3.2$\sigma$ and requires O(1) NP phases $\phi_1^d \equiv \arg(\Delta q) [2]$. Phases $\phi_4^d$ also enter into CP asymmetries due to interference between $B_0^0$-$B_d^0$ decay amplitudes without mixing: e.g., the asymmetry for $B_0^0 \to J/\psi K^0_S$ is $S_{\phi_4^d} = \sin(2\beta + \phi_4^d)$, with CKM angle $\beta \equiv \arg( - V_{td}V_{ts}^\dagger V_{tb}V_{tb}^\dagger )$. As emphasized in Ref. [4], the presence of non-zero $\phi_4^d$ can alleviate tension between $S_{\phi_4^d}$ and Br($B^+ \to \tau^+\nu$), which is sensitive to $\beta$ but not $\phi_4^d$.

To quantify these tensions, the CKMfitter group performed a global fit allowing for arbitrary $\Delta q/d,s$ (dubbed “Scenario I”), finding that the SM point ($\Delta q/d,s = 1$) is disfavored at 3.4$\sigma$. Moreover, their best fit point favors NP CP-violating phases in both $B_0^0$ and $B_d^0$ mixing: $\phi_4^d \equiv (-12^{+3.3}_{-2.9})^\circ$ and $\phi_6^d = (-129^{+12}_{-9.4})^\circ \cup (-51.6^{+14.1}_{-9.4})^\circ$.

In our model, NP effects enter $B_{d,s}^0$ observables predominantly through mixing, via box diagrams shown in Fig. 1. We find

$$\Delta q = 1 + c_{eq} f_1(x_H,x_t)/S_0(x_t) + c_{eq}^2 f_2(x_H,x_t)/S_0(x_t) , \quad (10)$$

where

$$c_{ij} \equiv \frac{(\bar y_U V_{ti})(\bar y_U V_{ti})}{4\sqrt{2}G_F m_W^2 V_{tb}V_{tb}^\dagger} . \quad (11)$$

The $\bar t_R d_{i,L}^H H^+$ charge current couplings are $(\bar y_U V_{ti}) = \bar y_{tt}^i V_{ti} + \bar y_{tt}^i V_{ci}$, for $i = d,s,b$. The NP loop functions are

$$F_1(x_H,x_t) = \frac{x_t x_H (x_H - 4) \log x_H}{(x_H - 1) (x_H - x_t)^2} - \frac{x_t (x_t - 4)}{x_t (x_H x_t - 2 x_t x_H + 4 x_H - 3 x_t^2) \log x_t} \frac{(x_t - 1)^2 (x_H - x_t)^2}{(x_H - x_t)^2} \quad (12)$$

$$F_2(x_H,x_t) = \frac{x_t^2 - x_t^2 - 2 x_t x_H \log(x_H/x_t)}{(x_H - x_t)^3} \quad (13)$$

where $x_{t,H} \equiv m_{t,H}^2/m_W^2$, and $S_0(x_t) \approx 2.35$ is the SM loop function (e.g., see [25]).

$B_{d,s}^0$-$B_{d,s}^0$ mixing from box graphs in a 2HDM have been computed previously [27]. Here, a novel feature arises from the NP CP-violating phase associated with $\bar y_{tt}$. We can write $(\bar y_U V_{ti})$ as

$$(\bar y_U V_{ti})_b \simeq \bar y_{tt} V_{ti} \quad (14)$$

$\bar y_U^i V_{ti} = \bar y_{tt} V_{ti} (1 + (\bar y_{tt} V_{ti}) e^{i\vartheta_{tt}})$

where $\vartheta_{tt} \equiv \arg(\bar y_{tt} V_{ti}^* V_{ti}^\dagger)$. In the limit $|\bar y_{tt}| \gg |\bar y_{tt}|$, we neglect the term $\bar y_{tt} V_{ti}$ for $i = b$; however, $\bar y_{tt}$ is non-negligible for $i = d,s$ because the $\bar y_{tt}$ terms are Cabibbo suppressed.

The NP phase that enters $(M_{12}^d)_{NP}$ is $\vartheta_{tt}$, while for $(M_{12}^d)_{NP}$ it is $(\vartheta_{tt} + \beta)$, due to the different CKM structures of $(\bar y_U V_{ti})$ and $(\bar y_U V_{ti})$. The best fit values for $\phi_{d,s}^4$ are quite different numerically, but due to this extra $e^{i\beta}$, we can explain both $\phi_{d,s}^4$ in terms of the single NP phase $\vartheta_{tt}$. For $\bar y_{tt} = 0$, our model gives $\phi_{d,s}^4 = 0$, since $(M_{12}^d)_{NP}$ would have the same complex phase $(V_{tb}V_{tb}^\dagger)_{SM}$.

Our results for $B_{d,s}^0$-$B_{d,s}^0$ mixing are shown in Fig. 2. Here, we map best fit regions for $\Delta q/d,s$ from Ref. [8] into the parameter space of our model. We fix $|\bar y_{tt}|$ and $m_{H^\pm}$ and evaluate the preferred regions for $|\bar y_{tt}|$ and $\vartheta_{tt}$ consistent with $B_{d,s}^0$-$B_{d,s}^0$ mixing constraints. (As discussed below, EWBG favors $|\bar y_{tt}| \sim 1$ and $m_{H^\pm} \lesssim 500$ GeV.) The blue (red) contours correspond to the best fit regions at 1$\sigma$ (inner) and 2$\sigma$ (outer), for $\Delta q/d,s$. Since $\Delta q/d,s$ are quadratic functions of $|\bar y_{tt}| e^{i\vartheta_{tt}}$, the best fit regions for $\Delta q/d,s$ each map into two best fit regions in $|\bar y_{tt}|$, $\vartheta_{tt}$ parameter space.

We also implement constraints on our model from $b \to s\gamma$ and $b \to sK$. The branching ratios for $b \to s\gamma$, as measured experimentally [29] and evaluated theoretically in the SM

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1 Ref. [8] did not include in their fit updated CDF and D0 results for $S_{\phi_4^d}$ [8], which showed improved consistency with the SM over previous results favoring non-zero $\phi_4^d$ [8].

2 We neglect running between the scales $m_t, m_W$, and $m^2_F$, integrating out these degrees of freedom at a common electroweak scale. Moreover, we have neglected a NP QCD correction factor $\eta(x_H,x_t)/\eta_B$ arising at next-to-leading order [25].
In Fig. 2, the white (light grey) region corresponds to constraint is due to model: in quadrature, we take the following constraint on our model at NLO following Refs. [30, 31], except that contributions have been computed at NLO only. At NNLO in the SM [32], we work at NLO since 2HDM contributions have been computed at NLO only.

The remaining SM input parameters in Eq. (14) are defined and tabulated in Ref. [8]. Assuming a theoretical error bar as in Ref. [33], we take the following constraint on our model

\[ |\epsilon_K|_{\text{SM} + \text{NP}} = (2.23 \pm 0.30) \times 10^{-3} \]  

(18)

It appears that since \( |\epsilon_K|_{\text{SM}} < |\epsilon_K|_{\text{exp}} \), this constraint would favor a small, positive contribution from NP. However, \( |\epsilon_K|_{\text{SM}} \) itself is shifted to a central value \( |\epsilon_K|_{\text{SM}} = 2.40 \times 10^{-3} \) because the best fit CKM parameters in the presence of NP in \( B_{d,s} \to B_{d,s} \) mixing (given in Table 11 of Ref. [8]) are different than in a SM-only fit. As a result, Eq. (18) favors a small, negative contribution from NP. In Fig. 2, the parameter region within the dashed dark (light) green contours is consistent with \( \epsilon_K \) constraint in Eq. (18) at 1σ (2σ).

Here, we make several important points.

- Despite the fact that \( \phi_\Delta^s \) and \( \phi_\Delta^c \) are quite different numerically, there exists regions of parameter space where both NP in \( B_0^0 - B_0^0 \) and \( B_d^0 - B_d^0 \) can be explained by a single phase \( \theta_{1c} \). The 1σ best fit regions for \( \Delta_{d,s} \) overlap within the parameter space of our model (neglecting correlations between \( \Delta_d \) and \( \Delta_s \)).

- The \( \Delta_s \) region that overlaps with the \( \Delta_d \) region in Fig. 2 corresponds to the \( \phi_\Delta^c = (-51.6_{-9.4}^{+ 14.1})^\circ \) solution. Therefore, our model predicts \( \Delta \Gamma_s > 0 \).

- Although \( b \to s\gamma \) and \( \epsilon_K \) constrain a large parametric region of our model, these two observables are consistent with observation in regions favored by \( \bar{B} \) mixing observables.

- A large phase \( \theta_{1c} \) can weaken \( b \to s\gamma \) and \( \epsilon_K \) constraints, and a light charged Higgs \( (m_{H^\pm} \sim 100 \text{ GeV}) \) is not excluded.

- The values of \( (|\bar{y}_t|, m_{H^\pm}) \) shown in Fig. 2 are consistent with \( \bar{R}_b = \text{BR}[Z \to b\bar{b}] / \text{BR}[Z \to \text{hadrons}] \) at 95% CL [12].

\[ |\epsilon_K|_{\text{SM} + \text{NP}} = (2.23 \pm 0.30) \times 10^{-3} \]
Although we chose only two illustrative values \((\bar{y}_{ul}, m_{H^\pm}) = (0.8, 100 \text{ GeV})\) and \((1.2, 350 \text{ GeV})\) in Fig. 2, there exists a consistency region between all these observables for parameters \(\bar{y}_{ul} \sim 1\), \(\bar{y}_{uc} \sim 0.05 - 0.1\), and \(\theta_t \sim 3\pi/4\), for \(100 < m_{H^\pm} < 500 \text{ GeV}\). As we discuss below, EWBG favors \(\bar{y}_{tc} \sim 1\) and \(m_{H^\pm} \lesssim 500 \text{ GeV}\).

**IV. ELECTROWEAK BARYOGENESIS**

Given a NP model, viable EWBG requires: (1) the electroweak phase transition must be strongly first order to prevent washout of baryon number, and (2) CP violation must be sufficient to account for the observed baryon-to-entropy ratio \(Y_B^{\text{obs}} \approx 9 \times 10^{-11}\). EWBG in a 2HDM has been studied many times previously. Most recently, Ref. [36] showed that a strong first order phase transition can occur in a type-II 2HDM for \(m_{H^0} \lesssim 200 \text{ GeV}\) and \(300 \lesssim m_{H^\pm} \lesssim 500 \text{ GeV}\). Although our 2HDM is not exactly the same as in Ref. [36], we assume that a strong first order transition does occur. (The phase transition can also be further strengthened or modified by the presence of scalar gauge singlets or non-renormalizable operators.)

We now study baryon number generation during the phase transition. The dynamical Higgs fields during the transition gives rise to a spacetime dependent mass matrix \(M(x)\) for, e.g., u-type quarks:

\[
\mathcal{L}_{\text{mass}} = -\bar{u}_R M u_L + \text{h.c.}, \quad M = y_U v_1(T) + \bar{y}_U v_2(T)
\]

where \(v_1, v_2(T) \equiv \langle H_{1,2}^0 \rangle_{T \neq 0}\) are the vevs at finite temperature \(T \approx 100 \text{ GeV}\). At zero temperature, when \(v_1(T), v_2(T) \to v, 0\), we recover the usual \(T = 0\) masses. However, if \(v_2(T) \neq 0\), then CP-violating quark charge density can arise from \(\bar{y}_U\), as we show below. Left-handed quark charge, in turn, leads to baryon number production through weak sphalerons. In previous studies, CP asymmetries were generated by a spacetime-dependent Higgs vev phase, arising from CP violation in the Higgs sector. Here, we assume that the Higgs potential is CP-conserving, such that \(v_1, v_2(T)\) do not have spacetime-dependent phases and can be taken to be real.

Is it plausible that \(v_2(T) \neq 0\) during the phase transition? Following Ref. [10], the most general potential for \(H_{1,2}\) can be written

\[
V = \lambda_1 H_1^4 - v^2 H_1^2 + m_{H_1}^2 H_1^2 + \lambda_2 H_2^4 + \lambda_3 (H_1^2 H_2^2) + \lambda_4 H_1^2 H_2^2 + \lambda_5 H_1^2 H_2^2 \quad \text{for} \quad \lambda_5 > 0
\]

Our basis choice that \(\langle H_2^0 \rangle_{T=0} = 0\) requires that no terms linear in \(H_2\) survive when \(H_1^0 \to v\). The same statement does not hold at \(T \neq 0\) due to thermal corrections to \(V\). First, since we expect \(v_1(T) \neq v\), terms linear in \(H_2\) appear proportional to \(\lambda_5\). Second, top quark loops generate a contribution to the potential \((y_u y_t T^2 H_1^2 H_2^2/4 + \text{h.c.})\), given here in the high \(T\) limit, also linear in \(H_2\). A proper treatment of this issue requires a numerical evaluation of the bubble wall solutions of the finite \(T\) Higgs potential, which is beyond the scope of this project. Here, we treat \(\tan \beta(T) \equiv v_2(T)/v_1(T)\) as a free parameter, and we work in the \(\beta(T) \sim 1\) limit. Intuitively, we expect \(\beta(T)\) to be suppressed in the limit \(m_{H_2} \gg T^2\), since the vev will be confined along the \(\langle H_2^0 \rangle = 0\) valley.

The charge transport dynamics of EWBG are governed by a system of Boltzmann equations of the form

\[
\dot{N}_a = S^a_{\alpha} + D_a \nabla^2 n_a + \sum_{b} \Gamma_{ab} n_b
\]

Here, \(n_a\) is the charge density for species \(a\). The CP-violating source \(S^a\) generates non-zero \(n_a\) within the expanding bubble wall, at the boundary between broken and unbroken phases, due to the spacetime-varying vevs \(v_1, v_2\). The diffusion constant \(D_a\) describes how \(n_a\) is transported ahead of the wall into the unbroken phase, where weak sphalerons are active. The remaining terms describe inelastic interactions that convert \(n_a\) into charge density of other species \(b\), with rate \(\Gamma_{ab}\). Our setup of the Boltzmann equations follows standard methods, described in detail in Ref. [40].

Following Ref. [39], we assume a planar bubble wall geometry, with velocity \(v_w \ll 1\) and coordinate \(z\) normal to the wall. The \(z > 0\) (\(z < 0\)) region corresponds to the (un)broken phase. We look for steady state solutions in the rest frame of the wall that only depend on \(z\). Therefore, we replace \(\dot{n}_a = v_w n_a' + \nabla^2 n_a \to n_a''\), where prime denotes \(\partial/\partial z\). We adopt kink bubble wall profiles

\[
v(T)/T = \xi [1 + \tanh(z/L_w)]/(2\sqrt{2}),
\]

\[
\beta(T) = \Delta \beta [1 + \tanh(z/L_w)]/2,
\]

where \(v(T) \equiv v_1(T)^2 + v_2(T)^2\). We take \(\xi = 1.5\), wall width \(L_w = 5/T\), and \(T = 100 \text{ GeV}\). Ref. [39] found viable first-order phase transitions with \(1 < \xi < 2.5\) and \(2 < L_w T < 15\), depending on the Higgs parameters. For definiteness, we take \(m_{H_2} = 400 \text{ GeV}\); however, our analysis does not account for the crucially important \(m_{H_2}\)-dependence of the bubble profiles.

Specializing to our 2HDM, the complete set of Boltzmann equations is

\[
v_w n_a' = D_q n_a'' + \delta_{qa} (S^a + \gamma_q Q_y + \gamma_m Q_m) + 2 \frac{\Gamma_{ss} Q_{ss}}{Q_{ss}}
\]

\[
v_w n_a' = D_q n_a'' - \delta_{qa} (S^a + \gamma_q Q_y + \gamma_m Q_m) + \Gamma_{ss} Q_{ss}
\]

\[
v_w n_H' = D_H n_H'' + \gamma_y Q_y - \Gamma_b Q_b
\]

\[4\text{ although the usual } \tan \beta \text{ is not physical at } T = 0, \text{ the angle } \beta(T) \text{ between the } T = 0 \text{ and } T \neq 0 \text{ vev directions is physical.}\]
with linear combinations of charge densities

$$Q_y = \frac{n_{u_3}}{k_{u_3}} - \frac{n_{q_3}}{k_{q_3}} - \frac{n_H}{k_H}, \quad Q_m = \frac{n_{u_3}}{k_{u_3}} - \frac{n_{q_3}}{k_{q_3}} \tag{24}$$

$$Q_{ss} = \sum_{a=1}^{3} \left( \frac{2n_{u_3}}{k_{u_3}} - \frac{n_{u_a}}{k_{u_a}} - \frac{n_{u_s}}{k_{u_s}} \right), \quad Q_h = \frac{n_H}{k_H}. \tag{25}$$

The relevant densities are the ath generation left(right)-handed quark charges $n_{qa}$ ($n_{ua}$, $n_{ua}$), and the Higgs charge density $n_H \equiv n_{H_1} + n_{H_2}$ (we treat $H_{1,2}$ as mass eigenstates in the unbroken phase). We assume that (Cabibbo unsuppressed) gauge interactions are in equilibrium, as are Higgs interactions that chemically equilibrate $H_{1,2}$ (provided by $\lambda_{3,4,5}$ quartic couplings in $V$). Lepton densities do not get sourced and can be neglected.

The $k$-factors are defined by $n_a = T^2 k_a \mu_a / 6$, with chemical potential $\mu_a$.

In the Eqs. (23), we take these transport coefficients as input:

$$S^{\text{CP}}_t \approx 0.1 \times N_c |y_t B_{\text{t}}| \sin \theta_{tt} v(T)^2 v_w \beta_y \Theta(T) T \tag{26}$$

$$\Gamma_m \approx 0.1 \times N_c |y_t v_1(T) + \tilde{y}_t v_2(T)|^2 T^{-1} \tag{27}$$

$$\Gamma_y \approx \frac{27 \zeta_3^2}{2 \pi^2} \alpha_s q^2 T + 9 |y_t|^2 T \left( \frac{m_H}{2 \pi T} \right)^{5/2} e^{-m_H^2/T} \tag{28}$$

$$\Gamma_{ss} \approx 14 \alpha_s^2 T, \quad D_q \approx 6 / T, \quad D_H \approx 100 / T. \tag{29}$$

We compute the CP-violating source $S^{\text{CP}}_t$ and relaxation rate $\Gamma_m$, arising for $t_{L,R}$ only, following the vev-insertion formalism [41, 42] (explicit forms can be found in [43]). The sole source of CP violation here is the phase $\theta_{tt} \equiv \arg (\tilde{y}_t)$, which is not the same phase that enters into $D^{dd}_L, D^{dd}_R$ mixing. The dimensionless numerical factors (0.1), obtained following Ref. [42], arise from integrals over $t_{L,R}$ quasi-particle momenta, taking as input are the thermal masses (tabulated in [43]) and thermal widths ($\gamma_{tt} \approx 0.15 g^2 T$ [46]). The top Yukawa rate $\Gamma_y$ comes from processes $H_{1,2} \leftrightarrow t \bar{t}$ and $H_2 \leftrightarrow t_{R,L}$ [43, 44]. The strong sphaleron rate $\Gamma_{ss}$ plays a crucial role in EWBG, in the 2HDM [49], discussed below, and $D_q, D_H$ are the Higgs diffusion constants [50]. The relaxation rate $\Gamma_h$ is due to Higgs charge non-conservation when the vev is zero. For simplicity, we set $\Gamma_h = \Gamma_m$ [51]; we find deviations from this estimate lead to $\lesssim O(1)$ variations in our computed $Y_B$. We have omitted from Eq. (25) additional Yukawa interactions induced by $y_{tc}$ (e.g., $H_2 \leftrightarrow t_{R,L}$) because we find they have negligible impact on $Y_B$. Moreover, CP-violating sources from $y_{tc}$ do not arise at leading order in vev-insertions. Therefore, $y_{tc}$ plays no role in our

The reparametrization invariant phase is $\theta_{tt} \equiv \arg (\tilde{y}_t)$, and $y_t, y_{tc}$ are real and positive.

5 Although there exist more sophisticated treatments, the reliability of quantitative EWBG computations remains an open question (see discussion in [43]).

6 The reparametrization invariant phase is $\theta_{tt} \equiv \arg (\tilde{y}_t)$, but we have adopted a convention where $v_{1,2}(T)$ and $y_t$ are real and positive.

EWBG setup (this conclusion may not hold beyond the vev-insertion formalism).

Thus far, we have neglected baryon number violation; this is reasonable since the weak sphaleron rate $\Gamma_{ws} \approx 120 a_T^2 T$ [51] is slow and out of equilibrium. Therefore, we solve for the total left-handed charge $n_L \equiv \sum_a n_{qa}$ from Eqs. (23), neglecting $\Gamma_{ws}$, and then treat $n_L$ as a source for baryon density $n_B$, according to

$$v_w n_B' - D_q n_B'' = -(3 \Gamma_{ws} n_L + R n_B) h, \tag{30}$$

with the relaxation rate $R = (15/4) \Gamma_{ws}$ [52]. The sphaleron profile $h(z)$ governs how $\Gamma_{ws}$ turns off in the broken phase [53]. Since the energy of the $T = 0$ sphaleron is $E_{\text{sph}} \approx 4 M_W / \alpha_w$, we take [54]

$$h(z) = \exp \left( - E_{\text{sph}}(T) / v \right), \quad E_{\text{sph}}(T) = E_{\text{sph}} v(T) / v. \tag{31}$$

Effectively, this cuts off the weak sphaleron rate for relatively small values of the vev: $v(T, z) / T \gtrsim g_2 / (8 \pi)$.

In Fig. 3, we show the spatial charge densities resulting from a numerical solution to Eqs. (23) for an example choice of parameters giving $Y_B \approx 9 \times 10^{-11}$. In general, the individual charge densities have long diffusion tails in the unbroken phase ($z < 0$). However, $n_L$ is strongly localized near the bubble wall ($z = 0$), due to strong sphalerons, thereby suppressing $n_B$. This effect can be understood as follows: at the level of Eqs. (23), $B$ is conserved, implying $\sum_a (n_{qa} + n_{ua} + n_{da}) = 0$; additionally, strong sphalerons relax the linear combination of densities

$$Q_{ss} \approx (1 / N_c) \sum_a (n_{qa} - n_{ua} - n_{da}) \tag{32}$$

to zero. These considerations imply that $n_L \approx 0$ if strong sphalerons are in equilibrium. In Fig. 3 we see that strong sphalerons are equilibrated and $n_L$ vanishes for $z \lesssim -10 L_w$. Since $n_L$ is non-zero only near the wall, it is important to treat the weak sphaleron profile accurately in this region, rather than with a simple step function. Nevertheless, despite this suppression, EWBG can account for $Y_B^{\text{obs}}$. (We also note the significant Higgs charge
We proposed a simple 2HDM that can account for these $B$ meson anomalies and the baryon asymmetry. An interesting feature of our setup is a top-charm flavor-violating Yukawa coupling of the new physics Higgs doublet. The large relative phase of this coupling can explain both the dimuon asymmetry and tension in $\text{BR}(B \to \tau \nu)$. Although top-charm flavor violation can give potentially large contributions to $b \to s \gamma$ and $\epsilon_K$ (i.e., less CKM-suppressed than SM contributions), these bounds are weakened in precisely the same region of parameter space consistent with $B_{d,s} \to B_{d,s}$ observables.

We also discussed electroweak baryogenesis. We showed that, provided a strong first-order electroweak phase transition occurs, our model can easily explain the observed baryon asymmetry of the Universe. CP violation during the phase transition is provided by the relative phase in the flavor-diagonal $t_L$-$t_R$ Yukawa coupling $\tilde{y}_L$ to the new Higgs, and the relevant phase is not related to the top-charm CP phase entering flavor observables. However, flavor observables and baryogenesis both require $|\tilde{y}_{1L}| \sim 1$. Additionally, baryon generation is dependent on a parameter $\Delta \beta$ related to the shift in the ratio of Higgs vevs across the bubble wall. We expect $\Delta \beta$ to be suppressed in the limit $m_{H^\pm} \gg m_W$. However, we showed that the charged Higgs state $H^\pm$ can be light ($m_{H^\pm} \sim 100 \, \text{GeV}$) without conflicting with flavor observables due to the large top-charm phase in our model (as opposed to the limit $m_{H^\pm} > 315 \, \text{GeV}$ from $b \to s \gamma$ in a type-II 2HDM [30, 31]).

It would be interesting to explore the consequences of our model for Higgs- and top-related CP-violating and flavor-violating observables measurable in colliders, and also for rare decays such as $K \to \pi \nu \bar{\nu}$. Additionally, a more robust analysis of EWBG requires an analysis of the finite temperature effective potential in a Type-III 2HDM, addressing the phase transition strength and bubble wall profiles.

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