PHYSICAL REVIEW D **84,** 034013 (2011)

Anomalous B-meson mixing and baryogenesis

Sean Tulin[1] and Peter Winslow[1,2]

[1]*Theory Group, TRIUMF, 4004 Wesbrook Mall, Vancouver, British Columbia, V6T 2A3, Canada*

[2]*Department of Physics and Astronomy, University of British Columbia, 6224 Agricultural Road, Vancouver, British Columbia, V6T 1Z1, Canada*

(Received 27 May 2011; published 9 August 2011)

There exist experimental hints from the B sector for CP violation beyond the standard model (SM) paradigm. An anomalous dimuon asymmetry was reported by the D0 Collaboration, while tension exists between $B \to \tau\nu$ and $S_{\psi K}$. These measurements, disfavoring the SM at the $\sim 3\sigma$ level, can be explained by new physics in both $B_d^0 - \bar{B}_d^0$ and $B_s^0 - \bar{B}_s^0$ 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 $B_{d,s}^0$-$\bar{B}_{d,s}^0$ mixing, weakens constraints from ϵ_K and $b \to s\gamma$, allowing for a light charged Higgs mass of $\mathcal{O}(100\ \text{GeV})$.

DOI: 10.1103/PhysRevD.84.034013 PACS numbers: 14.40.Nd, 12.60.Fr

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 [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 $B_{d,s}^0$ mesons, in excess over SM prediction at the 3.2σ level [2]. Second, there is tension at the $\sim 3\sigma$ level between the branching ratio for $B^+ \to \tau^+\nu$ and the CP asymmetry $S_{\psi K}$ in $B_d^0 \to J/\psi K$ [3,4]. Additionally, CDF and D0 have measured the CP asymmetry $S_{\psi\phi}$ in $B_s^0 \to J/\psi\phi$. While their earlier results (each with 2.8 fb^{-1} data) showed a $\sim 2\sigma$ deviation from the SM [5], this discrepancy has been reduced in their updated analyses with more data (5.2 and 6.1 fb^{-1}, respectively) [6].

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_d^0 and B_s^0 mesons [7]. Recently, the CKMfitter group has performed a global fit to all flavor observables, allowing for arbitrary new physics in $B_{d,s}^0$-$\bar{B}_{d,s}^0$ mixing amplitudes [8]. They conclude that the SM is disfavored at 3.4σ, while the data seem to favor NP with large CP-violating phases relative to the SM in *both* B_d^0 and B_s^0 mixing. At the level of effective theory, this NP takes the form

$$\mathcal{L}_{\mathrm{NP}} \sim \frac{c_d}{\Lambda^2}(\bar{b}d)^2 + \frac{c_s}{\Lambda^2}(\bar{b}s)^2 + \text{H.c.} \tag{1}$$

These operators can arise from new bosonic degrees of freedom at or near the weak scale, with new large CP-violating phases [9–12].

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 [13–15], is particularly attractive since two out of three Sakharov conditions [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 [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 [18]. This CP violation must involve particles with large couplings to the Higgs boson, since it is the interactions of those particles with the dynamical Higgs background field that leads to baryon production. Precision tests, such as electric dipole moment searches [19] and flavor observables, can probe directly CP violation relevant for EWBG. (The third condition, baryon number violation, is provided in the SM by weak sphalerons [20]; however, it is difficult to probe experimentally, since the sphaleron rate is highly suppressed at temperatures below the weak scale.)

If we wish to connect Eq. (1) 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 $B_{d,s}^0$ systems require that $\Lambda^2/|c_d| \gtrsim (500\ \text{TeV})^2$ and $\Lambda^2/|c_s| \gtrsim (100\ \text{TeV})^2$ [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 $\Lambda \lesssim 1$ 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. (1) 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 Λ, without conflicting with $\Delta M_{d,s}$ constraints.

In this work, we propose that a simple two Higgs doublet model (2HDM) can account for both anomalous CP violation in $B^0_{d,s}$-$\bar{B}^0_{d,s}$ mixing and EWBG. Previous works have studied CP violation in $B^0_{d,s}$-$\bar{B}^0_{d,s}$ mixing within a 2HDM [9–12]. Our setup, described in Sec. II, is different: we assume the NP Higgs doublet (H^+, $H^0 + iA^0$) mediates top-charm flavor violation. In this case, the NP $B^0_{d,s}$-$\bar{B}^0_{d,s}$ mixing amplitudes $(M_{12}^{d,s})_{\rm NP}$ are generated at one-loop order through charge current interactions mediated by H^+ (similar to Ref. [12]), rather than through tree-level exchange [9–11]. In Sec. III, we compute $(M_{12}^{d,s})_{\rm NP}$ in our model. We find the following:

(i) The best fit values to *both* M_{12}^d and M_{12}^s, from Ref. [8], can be explained in terms of a single NP phase ϑ_{tc} (defined below).

(ii) For large values of ϑ_{tc} preferred by $B^0_{d,s}$-$\bar{B}^0_{d,s}$ mixing observables, constraints from ϵ_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 that the parameter region favored by flavor observables (specifically, a large $\bar{t}_R t_L H^0$ coupling) can easily account for the observed baryon asymmetry. However, the relevant CP-violating phase is unrelated to the phase ϑ_{tc} entering flavor observables. In Sec. V, we summarize our conclusions.

II. MODEL

In a general (type-III) two Higgs doublet model [22,23], where both Higgs fields couple to each SM fermion, one can perform a field redefinition such that only one Higgs field acquires a real, positive vacuum expectation value (vev) [24]. We denote the two Higgs doublets by

$$H_1 = \begin{pmatrix} G^+ \\ v + \frac{h^0 + iG^0}{\sqrt{2}} \end{pmatrix}, \qquad H_2 = \begin{pmatrix} H^+ \\ \frac{H^0 + iA^0}{\sqrt{2}} \end{pmatrix}, \tag{2}$$

where h^0, H^0 (A^0) are the neutral (pseudo)scalars, $H^\pm$ is a charged scalar, and $G^{\pm,0}$ are the Goldstone modes eaten by the electroweak gauge bosons. The vev is $v \approx 174$ GeV. In general, the physical neutral states can be admixtures of h^0, H^0, A^0, depending on the details of the Higgs potential. We neglect mixing in our analysis; in this case, H_1 is exactly an SM Higgs doublet.

The most general Yukawa interaction for u-type quarks is

$$\mathcal{L}_{\rm yuk} \supset \bar{u}_R(y_U H_1 + \tilde{y}_U H_2) Q_L + \text{H.c.}, \tag{3}$$

where the left-handed quark doublet is $Q_L \equiv (u_L, V d_L)$. The SU(2)$_L$ contraction is $H_i Q_L \equiv H_i^+(V d_L) - H_i^0 u_L$. The 3×3 Yukawa matrices y_U and $\tilde{y}_U$ couple right-handed u-type quarks $u_R \equiv (u, c, t)_R$ to 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

$$y_U = \text{diag}(y_u, y_c, y_t) = \text{diag}(m_u, m_c, m_t)/v \tag{4}$$

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

$$\begin{aligned}\mathcal{L}_{\rm yuk} \supset & -\bar{d}_R(y_D H_1^\dagger + \tilde{y}_D H_2^\dagger) Q_L \\ & - \bar{e}_R(y_L H_1^\dagger + \tilde{y}_L H_2^\dagger) L_L + \text{H.c.},\end{aligned} \tag{5}$$

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 Yukawa couplings.

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

$$\tilde{y}_U = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & \tilde{y}_{tc} & \tilde{y}_{tt} \end{pmatrix}, \qquad \tilde{y}_{D,L} = 0. \tag{6}$$

That is, we consider a hierarchical structure, in the spirit of Ref. [23], where the t_R-t_L and t_R-c_L couplings are dominant (with $|\tilde{y}_{tt}| \gg |\tilde{y}_{tc}|$), 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-$\bar{B}^0_q$ system ($q = d, s$) is governed by the off-diagonal matrix element $M_{12}^q - \frac{i}{2}\Gamma_{12}^q$ in the Hamiltonian [25,26], with M_{12}^q (Γ_{12}^q) associated with the (anti-)Hermitian part. Only the relative phase $\phi_q \equiv \arg(-M_{12}^q/\Gamma_{12}^q)$ is physical. The relevant observables are the mass and width differences between the two eigenstates

$$\Delta M_q = 2|M_{12}^q|, \qquad \Delta\Gamma_q = 2|\Gamma_{12}^q|\cos\phi_q, \tag{7}$$

and the wrong sign semileptonic asymmetry

$$a_{\rm sl}^q \equiv \frac{\Gamma(\bar{B}_q^0 \to \mu^+ X) - \Gamma(B_q^0 \to \mu^- X)}{\Gamma(\bar{B}_q^0 \to \mu^+ X) + \Gamma(B_q^0 \to \mu^- X)} = \frac{|\Gamma_{12}^q|}{|M_{12}^q|} \sin\phi_q. \tag{8}$$

The dimuon asymmetry measured by D0 arises from wrong sign semileptonic decays of both B_d^0 and B_s^0 mesons and is given by $A_{\rm sl}^b \approx 0.5 a_{\rm sl}^d + 0.5 a_{\rm sl}^s$ [2].

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

$$M_{12}^q = (M_{12}^q)_{\rm SM} + (M_{12}^q)_{\rm NP} \equiv (M_{12}^q)_{\rm SM} \Delta_q. \tag{9}$$

The consistency of $\Delta M_{d,s}$ with SM predictions constrains $|\Delta_{d,s}| \approx 1$, at the $\mathcal{O}(20\%)$ level [8], while the dimuon asymmetry measurement disagrees with SM prediction at 3.2σ and requires $\mathcal{O}(1)$ NP phases $\phi_q^\Delta \equiv \arg(\Delta_q)$ [2]. Phases ϕ_q^Δ also enter into CP asymmetries due to interference between $B_{d,s}^0$ decay amplitudes with and without mixing: e.g., the asymmetry for $B_d^0 \to J/\psi K_S^0$ is $S_{\phi K_S} = \sin(2\beta + \phi_d^\Delta)$, with CKM angle $\beta \equiv \arg(-V_{cd}V_{cb}^* V_{td}^* V_{tb})$. As emphasized in Ref. [4], the presence of nonzero ϕ_d^Δ can alleviate tension between $S_{\phi K_S}$ and $\mathrm{Br}(B^+ \to \tau^+ \nu)$, which is sensitive to β but not ϕ_d^Δ.

To quantify these tensions, the CKMfitter group performed a global fit allowing for arbitrary $\Delta_{d,s}$ (dubbed "Scenario I"), finding that the SM point ($\Delta_d = \Delta_s = 1$) is disfavored at 3.4σ [8]. Moreover, their best fit point favors NP CP-violating phases in both B_d^0 and B_s^0 mixing: $\phi_d^\Delta = (-12^{+3.3}_{-3.4})^\circ$ and $\phi_s^\Delta = (-129^{+12}_{-12})^\circ \cup (-51.6^{+14.1}_{-9.4})^\circ$.[1]

In our model, NP effects enter $B_{d,s}^0$ observables predominantly through mixing, via box diagrams shown in Fig. 1. We find[2]

$$\Delta_q = 1 + c_{bq} F_1(x_H, x_t)/S_0(x_t) + c_{bq}^2 F_2(x_H, x_t)/S_0(x_t), \tag{10}$$

where

$$c_{ij} \equiv \frac{(\tilde{y}_U V)_{ti} (\tilde{y}_U V)_{tj}^*}{4\sqrt{2} G_F m_W^2 V_{ti} V_{tj}^*}. \tag{11}$$

The $\bar{t}_R d_L^i H^+$ charge current couplings are $(\tilde{y}_U V)_{ti} = \tilde{y}_{tt} V_{ti} + \tilde{y}_{tc} V_{ci}$, for $i = d, s, b$. The NP loop functions are

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

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

where $x_{t,H} \equiv m_{t,H^\pm}^2 / m_W^2$, and $S_0(x_t) \approx 2.35$ is the SM loop function (e.g., see [26]).

$B_{d,s}^0$-$\bar{B}_{d,s}^0$ mixing from box graphs in a 2HDM have been computed previously [28]. Here, a novel feature arises from the NP CP-violating phase associated with $\tilde{y}_{tc}$ [29]. We can write $(\tilde{y}_U V)_{ti}$ as

$$\begin{aligned} (\tilde{y}_U V)_{tb} &\simeq \tilde{y}_{tt} V_{tb}, \\ (\tilde{y}_U V)_{ts} &= \tilde{y}_{tt} V_{ts} \left(1 + \left|\frac{\tilde{y}_{tc} V_{cs}}{\tilde{y}_{tt} V_{ts}}\right| e^{i\vartheta_{tc}}\right), \\ (\tilde{y}_U V)_{td} &= \tilde{y}_{tt} V_{td} \left(1 + \left|\frac{\tilde{y}_{tc} V_{cd}}{\tilde{y}_{tt} V_{td}}\right| e^{i(\vartheta_{tc} + \beta)}\right), \end{aligned} \tag{14}$$

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

The NP phase that enters $(M_{12}^s)_{\rm NP}$ is ϑ_{tc}, while for $(M_{12}^d)_{\rm NP}$ it is $(\vartheta_{tc} + \beta)$, due to the different CKM structures of $(\tilde{y}_U V)_{ts}$ and $(\tilde{y}_U V)_{td}$. The best fit values for $\phi_{d,s}^\Delta$ are quite different numerically, but due to this extra $e^{i\beta}$, we can explain both $\phi_{d,s}^\Delta$ in terms of the single NP phase ϑ_{tc}. [For $\tilde{y}_{tc} = 0$, our model gives $\phi_{d,s}^\Delta = 0$, since $(M_{12}^q)_{\rm NP}$ would have the same complex phase $(V_{tb} V_{tq}^*)^2$ as $(M_{12}^q)_{\rm SM}$.]

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

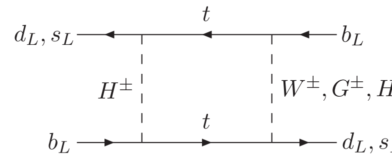

FIG. 1. New physics B_d^0-$\bar{B}_d^0$ and B_s^0-$\bar{B}_s^0$ mixing amplitudes $(M_{12}^{d,s})_{\rm NP}$ arising from box graphs with $H^\pm$ exchange.

[1]Reference [8] did not include in their fit updated CDF and D0 results for $S_{\phi\psi}$ [6], which showed improved consistency with the SM over previous results favoring nonzero ϕ_s^Δ [5].

[2]We neglect running between the scales m_t, m_W, and $m_H^\pm$, 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 [27].

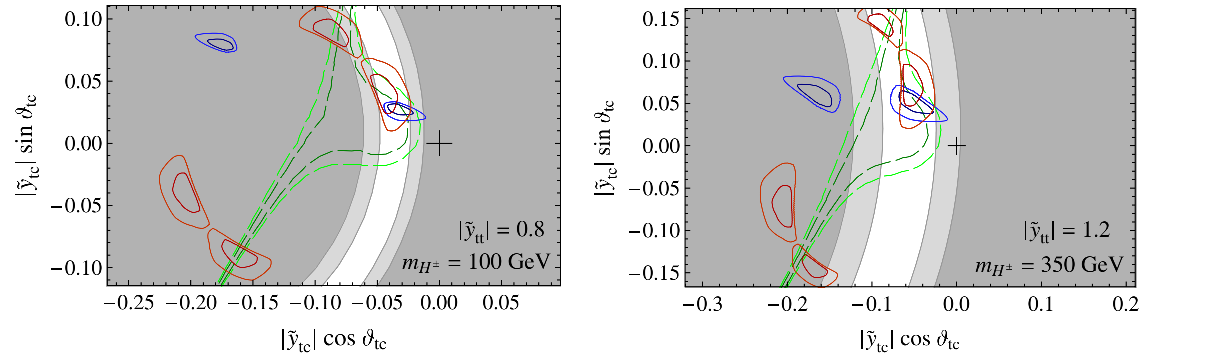

FIG. 2 (color online). Top-charm flavor violation parameter space ($|\tilde{y}_{tt}|$, ϑ_{tc}) consistent with flavor observables, for two choices of $|\tilde{y}_{tt}|$, $m_{H\pm}$ 68% and 95% C.L. regions for Δ_d (Δ_s) from Ref. [8] shown by blue (red) contours. Region within dark (light) dashed green contours is consistent with ϵ_K at 68% (95%) C.L. Light (dark) grey region is excluded at 68% (95%) C.L. from $BR(\bar{B} \to X_s\gamma)$.

We also implement constraints on our model from $b \to s\gamma$ and ϵ_K. The branching ratios for $b \to s\gamma$, as measured experimentally [30] and evaluated theoretically in the SM at next-to-leading order (NLO) [31], are given by[3]

$$\mathrm{BR}[\bar{B} \to X_s\gamma]^{\rm exp}_{E_\gamma > 1.6\,\mathrm{GeV}} = (3.55 \pm 0.24 \pm 0.09) \times 10^{-4},$$

$$\mathrm{BR}[\bar{B} \to X_s\gamma]^{\rm SM}_{E_\gamma > 1.6\,\mathrm{GeV}} = (3.60 \pm 0.30) \times 10^{-4}. \qquad (15)$$

We evaluate SM + NP contributions to $\mathrm{BR}[\bar{B} \to X_s\gamma]$ in our model at NLO following Refs. [31,33], except that we take as inputs the best fit CKM parameters given in Table 11 of Ref. [8]. Adding all errors in Eqs. (15) in quadrature, we take the following constraint on our model:

$$\mathrm{BR}\,[\bar{B} \to X_s\gamma]^{\rm SM+NP}_{E_\gamma > 1.6\,\mathrm{GeV}} = (3.55 \pm 0.39) \times 10^{-4}. \quad (16)$$

In Fig. 2, the white (light grey) region corresponds to $|\tilde{y}_{tc}|$, ϑ_{tc} parameter space consistent with Eq. (16) at less than 1σ (2σ), while the dark grey region is excluded at 2σ.

NP contributions to K^0-$\bar{K}^0$ mixing arise in our model through box graphs analogous to Fig. 1. The strongest constraint is due to ϵ_K. In the SM, $|\epsilon_K|_{\rm SM} = (1.90 \pm 0.26) \times 10^{-3}$ [34], while experimentally $|\epsilon_K|_{\rm exp} = (2.228 \pm 0.011) \times 10^{-3}$ [35]. The SM + NP value of ϵ_K is

$$\begin{aligned} |\epsilon_K|_{\rm SM+NP} = \kappa_\epsilon C_\epsilon \hat{B}_K\, \mathrm{Im}[&(V_{ts}V_{td}^*)^2 \eta_2 (S_0(x_t) \\ &+ c_{sd} F_1(x_H, x_t) + c_{sd}^2 F_2(x_H, x_t)) \\ &+ (V_{cs}V_{cd}^*)^2 \eta_1 S_0(x_c) \\ &+ 2(V_{cs}V_{cd}^* V_{ts}V_{td}^*) \eta_3 S_0(x_c, x_t)], \end{aligned} \qquad (17)$$

where NP enters through the coefficients c_{sd} defined in Eq. (11). (We neglect NP NLO corrections to η_2.) The remaining SM input parameters in Eq. (17) are defined and tabulated in Ref. [8]. Assuming a theoretical error bar as in Ref. [34], we take the following constraint on our model:

$$|\epsilon_K|_{\rm SM+NP} = (2.23 \pm 0.30) \times 10^{-3}. \qquad (18)$$

It appears that since $|\epsilon_K|_{\rm SM} < |\epsilon_K|_{\rm exp}$, this constraint would favor a small, positive contribution from NP. However, $|\epsilon_K|_{\rm SM}$ itself is shifted to a central value $|\epsilon_K|_{\rm SM} = 2.40 \times 10^{-3}$ because the best fit CKM parameters in the presence of NP in $B^0_{d,s}$-$\bar{B}^0_{d,s}$ mixing (given in Table 11 of Ref. [8]) are different than in an 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 the ϵ_K constraint in Eq. (18) at 1σ (2σ).

Here, we make several important points.

(i) Despite the fact that ϕ_d^Δ and ϕ_s^Δ are quite different numerically, there exists regions of parameter space where *both* NP in B_d^0-$\bar{B}_d^0$ and B_s^0-$\bar{B}_s^0$ can be explained by a single phase ϑ_{tc}. The 1σ best fit regions for $\Delta_{d,s}$ overlap within the parameter space of our model (neglecting correlations between Δ_d and Δ_s).

(ii) The Δ_s region that overlaps with the Δ_d region in Fig. 2 corresponds to the $\phi_s^\Delta = (-51.6^{+14.1}_{-9.4})°$ solution. Therefore, our model predicts $\Delta\Gamma_s > 0$.

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

(iv) A large phase ϑ_{tc} can weaken $b \to s\gamma$ and ϵ_K constraints, and a light charged Higgs boson ($m_{H^\pm} \sim 100$ GeV) is not excluded.

(v) The values of ($|\tilde{y}_{tt}|$, $m_{H^\pm}$) shown in Fig. 2 are consistent with $R_b \equiv \mathrm{BR}[Z \to b\bar{b}]/\mathrm{BR}[Z \to \mathrm{hadrons}]$ at 95% C.L. [12].

Although we chose only two illustrative values ($|\tilde{y}_{tt}|$, $m_{H^\pm}$) = (0.8, 100 GeV) and (1.2, 350 GeV) in

[3]In the observed value, the first error is experimental, while the second is a theoretical error associated with a photon shape function used to extrapolate the branching ratio to different photon energies E_γ. Also, although $\mathrm{BR}[\bar{B} \to X_s\gamma]$ has been computed at NNLO in the SM [32], we work at NLO since 2HDM contributions have been computed at NLO only.

Fig. 2, there exists a consistency region between all these observables for parameters $|\tilde{y}_{tt}| \sim 1$, $|\tilde{y}_{tc}| \sim 0.05$–$0.1$, and $\vartheta_{tc} \sim 3\pi/4$, for $100 < m_{H^\pm} < 500$ GeV. As we discuss below, EWBG favors $|\tilde{y}_{tt}| \sim 1$ and $m_{H^\pm} \lesssim 500$ GeV.

IV. ELECTROWEAK BARYOGENESIS

Given an 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 $Y_B^{\rm obs} \approx 9 \times 10^{-11}$. EWBG in a 2HDM has been studied many times previously [36]. Most recently, Ref. [37] showed that a strong first-order phase transition can occur in a type-II 2HDM for $m_{h^0} \lesssim 200$ GeV and $300 \lesssim m_{H^0} \lesssim 500$ GeV. Although our 2HDM is not exactly the same as in Ref. [37], 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 [38] or nonrenormalizable operators [39].)

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:

$$\mathcal{L}_{\rm mass} = -\bar{u}_R M u_L + \text{H.c.},\qquad M = y_U v_1(T) + \tilde{y}_U v_2(T), \tag{19}$$

where $v_{1,2}(T) \equiv \langle H^0_{1,2}\rangle_{T\neq 0}$ are the vevs at finite temperature $T \approx 100$ 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 $\tilde{y}_U$, 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 [36,37]. Here, we assume that the Higgs potential is CP-conserving, such that $v_{1,2}(T)$ 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 [10], the most general potential for $H_{1,2}$ can be written

$$\begin{aligned} V = {}& \lambda(H_1^\dagger H_1 - v^2)^2 + m_{H_2}^2 H_2^\dagger H_2 + \lambda_1 H_1^\dagger H_1 H_2^\dagger H_2 \\ & + \lambda_2 H_1^\dagger H_2 H_2^\dagger H_1 + [\lambda_3 (H_1^\dagger H_2)^2 + \lambda_4 H_1^\dagger H_2 H_2^\dagger H_2 \\ & + \lambda_5 H_2^\dagger H_1 (H_1^\dagger H_1 - v^2) + \text{H.c.}] + \lambda_6 (H_2^\dagger H_2)^2 \end{aligned} \tag{20}$$

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 λ_5. Second, top quark loops generate a contribution to the potential $(y_t \tilde{y}_{tt} T^2 H_1^\dagger H_2/4 + \text{H.c.})$, 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,[4] and we work in the $\beta(T) \ll 1$ limit. Intuitively, we expect $\beta(T)$ to be suppressed in the limit $m_{H_2}^2 \gg T^2$, 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 $\dot{n}_a = S_a^{\not{CP}} + D_a \nabla^2 n_a + \sum_b \Gamma_{ab} n_b$ [40]. Here n_a is the charge density for species a. The CP-violating source $S_a^{\not{CP}}$ generates nonzero n_a within the expanding bubble wall, at the boundary between broken and unbroken phases, due to the spacetime-varying vevs $v_{1,2}(T)$. 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 Γ_{ab}. Our setup of the Boltzmann equations follows standard methods, described in detail in Ref. [41].

Following Ref. [40], 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 $\dot{n}_a \to v_w n_a'$ and $\nabla^2 n_a \to n_a''$, where prime denotes $\partial/\partial z$. We adopt kink bubble wall profiles

$$v(T)/T = \xi[1 + \tanh(z/L_w)]/(2\sqrt{2}), \tag{21}$$

$$\beta(T) = \Delta\beta[1 + \tanh(z/L_w)]/2, \tag{22}$$

where $v(T)^2 \equiv v_1(T)^2 + v_2(T)^2$. We take $\xi = 1.5$, wall width $L_w = 5/T$, and $T = 100$ GeV. Reference [37] 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 $m_{H_2} = 400$ 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

$$\begin{aligned} v_w n'_{q_a} &= D_q n''_{q_a} + \delta_{3a}(S_t^{\not{CP}} + \Gamma_y Q_y + \Gamma_m Q_m) - 2\Gamma_{ss} Q_{ss}, \\ v_w n'_{u_a} &= D_q n''_{u_a} - \delta_{3a}(S_t^{\not{CP}} + \Gamma_y Q_y + \Gamma_m Q_m) + \Gamma_{ss} Q_{ss}, \\ v_w n'_{d_a} &= D_q n''_{d_a} + \Gamma_{ss} Q_{ss}, \\ v_w n'_H &= D_H n''_H + \Gamma_y Q_y - \Gamma_h Q_h, \end{aligned} \tag{23}$$

with linear combinations of charge densities

[4] Although the usual $\tan\beta$ is not physical at $T = 0$, the angle $\beta(T)$ between the $T = 0$ and $T \neq 0$ vev directions *is* physical.

$$Q_y \equiv \frac{n_{u_3}}{k_{u_3}} - \frac{n_{q_3}}{k_{q_3}} - \frac{n_H}{k_H}, \qquad Q_m \equiv \frac{n_{u_3}}{k_{u_3}} - \frac{n_{q_3}}{k_{q_3}},$$
$$Q_{ss} \equiv \sum_{a=1}^{3}\left(\frac{2n_{q_a}}{k_{q_a}} - \frac{n_{u_a}}{k_{u_a}} - \frac{n_{u_a}}{k_{u_a}}\right), \qquad Q_h = \frac{n_H}{k_H}. \tag{24}$$

The relevant densities are the ath generation left(right)-handed quark charges n_{q_a} (n_{u_a}, n_{d_a}), 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 μ_a.

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

$$S_t^{\not{CP}} \approx 0.1 \times N_c |y_t \tilde{y}_{tt}| \sin\theta_{tt} v(T)^2 v_w \beta(T)' T, \tag{26}$$

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

$$\Gamma_y \approx \frac{27\zeta_3^2}{2\pi^2}\alpha_s y_t^2 T + 9|\tilde{y}_{tt}|^2 T\left(\frac{m_{H_2}}{2\pi T}\right)^{5/2} e^{-m_{H_2}/T}, \tag{28}$$

$$\Gamma_{ss} \approx 14\alpha_s^4 T, \qquad D_q \approx 6/T, \qquad D_H \approx 100/T. \tag{29}$$

We compute the CP-violating source $S_t^{\not{CP}}$ and relaxation rate Γ_m, arising for $t_{L,R}$ only, following the vev-insertion formalism [42,43] (explicit formulas can be found in [44]).[5] The sole source of CP violation here is the phase $\theta_{tt} \equiv \arg(\tilde{y}_{tt})$, which is *not* the same phase that enters into $B^0_{d,s}$-$\bar{B}^0_{d,s}$ mixing.[6] The dimensionless numerical factors (0.1), obtained following Ref. [43], arise from integrals over $t_{L,R}$ quasiparticle momenta, taking as input the thermal masses (tabulated in [46]) and thermal widths ($\gamma_{t_{L,R}} \approx 0.15 g_s^2 T$ [47]). The top Yukawa rate Γ_y comes from processes $H_1 t_L \leftrightarrow t_R g$ and $H_2 \leftrightarrow t_R \bar{t}_L$ [46,48]. The strong sphaleron rate Γ_{ss} [49] plays a crucial role in EWBG in the 2HDM [50], discussed below, and $D_{q,H}$ are the quark and Higgs diffusion constants [51]. The relaxation rate Γ_h is due to Higgs charge nonconservation when the vev is nonzero. For simplicity, we set $\Gamma_h = \Gamma_m$ [40]; we find deviations from this estimate lead to $\lesssim \mathcal{O}(1)$ variations in our computed Y_B. We have omitted from Eq. (23) additional Yukawa interactions induced by y_{tc} (e.g., $H_2 \leftrightarrow t_R \bar{c}_L$) 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 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 $\Gamma_{ws} \approx 120\alpha_w^5 T$ [52] is slow and out of equilibrium. Therefore, we solve for the total left-handed charge $n_L \equiv \sum_a n_{q_a}$ from Eqs. (23), neglecting Γ_{ws}, and then treat n_L as a source for baryon density n_B, according to

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

with the relaxation rate $\mathcal{R} = (15/4)\Gamma_{ws}$ [53]. The sphaleron profile $h(z)$ governs how Γ_{ws} turns off in the broken phase [54]. Since the energy of the $T = 0$ sphaleron is $E_{\rm sph} \approx 4M_W/\alpha_w$, we take [55]

$$h(z) = \exp(-E_{\rm sph}(T)/T), \qquad E_{\rm sph}(T) = E_{\rm sph} v(T)/v. \tag{31}$$

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 into 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 [50]. This effect can be understood as follows: at the level of Eqs. (23), B is conserved, implying $\sum_a (n_{q_a} + n_{u_a} + n_{d_a}) = 0$; additionally, strong sphalerons relax the linear combination of densities

$$Q_{ss} \approx (1/N_c)\sum_a (n_{q_a} - n_{u_a} - n_{d_a}) \tag{32}$$

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 $z \lesssim -10L_w$. Since n_L is nonzero 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

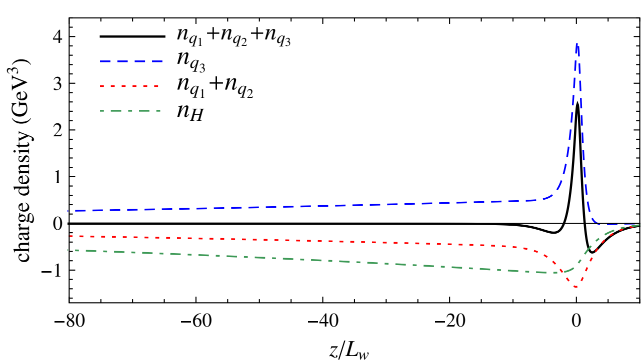

FIG. 3 (color online). Charge densities in the unbroken phase ($z < 0$) for $|\tilde{y}_{tt}| = 1$, $\theta_{tt} = 0.18$, $v_w = 0.05$, $\Delta\beta = 10^{-2}$, giving $Y_B \approx 9 \times 10^{-11}$.

[5] Although there exist more sophisticated treatments, the reliability of quantitative EWBG computations remains an open question (see discussion in [45]).

[6] The reparametrization invariant phase is $\theta_{tt} \equiv \arg(\tilde{y}_{tt} y_t^* v_1^* v_2)$, but we have adopted a convention where $v_{1,2}(T)$ and y_t are real and positive.

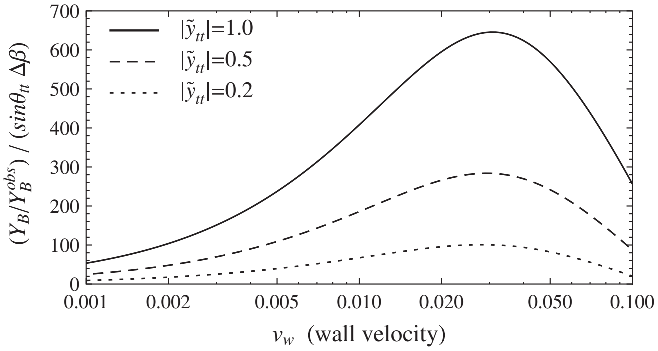

FIG. 4. Computed baryon asymmetry Y_B, normalized with respect to $Y_B^{\rm obs} \sim 9 \times 10^{-11}$ and $(\sin\theta_{tt}\Delta\beta)$, as a function of v_w and $|\tilde{y}_{tt}|$. Vertical axis shows $(\sin\theta_{tt}\Delta\beta)^{-1}$, required for viable EWBG.

for $Y_B^{\rm obs}$. (We also note the significant Higgs charge n_H in the broken phase. Although we neglect lepton Yukawas here, it is possible that n_H could be efficiently transferred into left-handed lepton charge via $\tilde{y}_L$, thereby driving EWBG without suffering from strong sphaleron suppression, analogous to Ref. [46].)

In Fig. 4, we show how large Y_B can be in our model. The most important parameters are $\Delta\beta$, $\tilde{y}_{tt}$, and v_w (we find Y_B is not strongly sensitive to L_w or ξ). The vertical axis shows the (inverse) value of $\Delta\beta \times \sin\theta_{tt}$ required for successful EWBG ($Y_B = Y_B^{\rm obs}$), for different values of $|\tilde{y}_{tt}|$ and v_w. Our main conclusion is that our model can easily account for the baryon asymmetry of the Universe—even if $\Delta\beta$ is as small as $10^{-3} - 10^{-2}$, provided the NP Yukawa coupling has magnitude $|\tilde{y}_{tt}| \gtrsim 0.2$, with $\mathcal{O}(1)$ phase. Moreover, $|\tilde{y}_{tt}| \sim 1$ is preferred by consistency with flavor observables.

V. CONCLUSIONS

The dimuon asymmetry reported by D0 [2] and the branching ratio BR$(B \to \tau\nu)$ [3,4] seem to disfavor the CKM paradigm of CP violation in the SM at the $\sim 3\sigma$ level. Although more experimental scrutiny is required, taken at face value, these anomalies can be accounted for by new physics in both B_d^0-$\bar{B}_d^0$ and B_s^0-$\bar{B}_s^0$ mixing [8]. Such new physics would involve new weak-scale bosonic degrees of freedom and new large CP-violating phases. These two ingredients are precisely what is required for viable electroweak baryogenesis in extensions of the SM.

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 BR$(B \to \tau\nu)$. Although top-charm flavor violation can give potentially large contributions to $b \to s\gamma$ and ϵ_K (i.e., less CKM-suppressed than SM contributions), these bounds are weakened in precisely the same region of parameter space consistent with $B_{d,s}^0$-$\bar{B}_{d,s}^0$ observables.

We also discussed electroweak baryogenesis. We showed that, provided a strong first-order eletroweak 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 $\tilde{y}_{tt}$ to the new Higgs boson, and the relevant phase is not related to the top-charm CP phase entering flavor observables. However, flavor observables and baryogenesis both require $|\tilde{y}_{tt}| \sim 1$. 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 ($m_{H^\pm} \sim 100$ GeV) without conflicting with flavor observables due to the large top-charm phase in our model (as opposed to the limit $m_{H^\pm} > 315$ GeV from $b \to s\gamma$ in a type-II 2HDM [31,35]).

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.

ACKNOWLEDGMENTS

We thank the authors of Ref. [8] for sharing with us numerical data from their global fit analysis. We are also indebted to K. Blum, V. Cirigliano, J. Cline, Y. Hochberg, D. Morrissey, Y. Nir, M. Pospelov, and M. Trott for helpful discussions and suggestions. S. T. is supported by NSERC of Canada.

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