

Mesogenesis through the ephemeral dark decay of B mesons

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Mesogenesis provides a path for generating the baryon asymmetry of the Universe, using only the CP violation furnished by the Standard Model in the decay of B mesons. While this is an intriguing possibility, it is largely constrained by the data on B meson branching fractions into baryons and missing energy carried into the dark sector. We point out that it is possible to make this branching fraction dominant only in the early Universe, through an ultralight scalar coupled to the dark sector and the Standard Model leptons. A scenario is examined where the thermal density of muons in the early Universe temporarily lowers the mass of a dark fermion, allowing for efficient B meson decays. This “dark” decay channel is shut off later when the muon number density falls, making the scenario compatible with flavor data. Our model can be consistent with the LHC constraints on color-charged heavy bosons required to implement mesogenesis; such states may be discovered in the future runs as their masses cannot be far above the current bounds. We also outline other possible signals, which can arise in future displaced vertex searches, long range force searches, and observations of neutron star binary mergers.

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I. INTRODUCTION

The mechanism for generation of the observed baryon asymmetry of the Universe (BAU) [1]

$$Y_B \equiv (n_B - n_{\bar{B}})/s = (8.718 \pm 0.004) \times 10^{-11}, \quad (1)$$

remains an open question in particle physics and cosmology. In the above, n_B and $n_{\bar{B}}$ are the number densities of baryons and antibaryons, respectively; s is the entropy density. Quite generally, it is expected that whatever mechanism generated the BAU must satisfy the Sakharov conditions [2], one of which is the presence of CP violation. While the Standard Model (SM) contains CP violating interactions, it is widely believed that it is by far not at sufficient levels [3–6]. This circumstance has been one of the main motivations to postulate and look for new CP violating physics.

A counter example to the above conclusion was proposed in Ref. [7], where the generation of the BAU was sourced by

the CP violation present in the SM, specifically in certain B meson decays in the early Universe.¹ The model in Ref. [7] is a mesogenesis model [17–20] modified to overcome exotic B -meson branching fraction constraints as well as LHC constraints on color-triplet scalars in order to be able to utilize the measured CP violation in the B meson sector. In mesogenesis, a color-triplet scalar particle mediates interactions between b -quarks and the dark sector. (The new interactions are not necessarily baryon number violating since the dark sector can also carry baryon number. However, baryon number in the visible sector can still be *violated* while being balanced by an equal and opposite asymmetry in the dark sector.) In Ref. [7], it is proposed that this mediator go through a mass change in the early Universe, starting with $M_\gamma = O(100 \text{ GeV})$ before the generation of the BAU and ending with an $O(\text{TeV})$ mass, currently. These different masses allow for a large branching fraction for the B -violating B -meson decays in the early Universe while being safe from constraints today.

The “morphing” of the mediator mass was achieved through generation of domain walls in which another scalar field acquires near-degenerate vacuum expectation values. In half of the domains, the mediator is lighter and the BAU is copiously produced. However, the domains with the heavier mediator is at the global minimum of the scalar

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¹For other examples of utilizing the Cabibbo–Kobayashi–Maskawa (CKM) phase for baryogenesis by dynamical effects in the early Universe, see Refs. [8–16].

potential and eventually the domain walls annihilate. This intriguing domain wall mechanism, unfortunately, might suffer from constraints from big bang nucleosynthesis (BBN) [21]. Even though the domain wall annihilation happens before BBN, at temperatures $O(10 \text{ MeV})$, the resulting baryon asymmetry, which is only correctly produced in certain domains, might not be able to homogenize before BBN and this inhomogeneity could lead to incorrect light element abundances.

In this paper, we discuss a different mechanism that will allow for a significant exotic B -meson branching fraction in the early Universe while being safe from such constraints today. Instead of morphing of the mediator mass, here we focus on the dark sector fermion,² ψ_B , into which the B meson decays, together with SM quarks (see Sec. II, below). We follow Refs. [23,24] to envision a scenario in which coupling an ultralight scalar to SM charged leptons lowers the mass of the dark fermion in the early Universe, compared to today.

In what follows, we will first briefly outline the basic ingredients of mesogenesis. Next, we will provide a model that makes the key B meson decays that involve the dark sector fermion ψ_B dominant while the Universe was at temperatures of $O(10 \text{ MeV})$, but kinematically forbidden at later times. This is achieved by making the ψ_B mass dependent on an ultralight scalar that can be sourced by charged leptons in the early Universe. We will also discuss how the needed effective operators can be generated by certain color charged scalars that can evade current LHC bounds, but may be expected to be probed in the coming years. Finally, we will outline some of the potential signals of our scenario that may arise in laboratory and astrophysical settings.

II. MESOGENESIS: A BRIEF OUTLINE

We start with a brief summary of the mesogenesis mechanism [17–20]. In general terms, mesogenesis can be described by dim-6 effective operators between SM quarks and a new dark fermion ψ_B carrying antibaryon number, $B = -1$, giving rise to a (B -conserving) effective operator

$$\mathcal{O}_{d_k, u_i, d_j} = \mathcal{C}_{d_k, u_i, d_j} \epsilon_{\alpha\beta\gamma} (\bar{\psi}_B d_k^\alpha) (\bar{d}_j^{c\beta} u_i^\gamma), \quad (2)$$

where i, j and α, β, γ are flavor and color indices, respectively. Note that in order to kinematically forbid proton decay through the $\psi_B udd$ interaction, we require all dark sector baryons to have a mass $\gtrsim 1 \text{ GeV}$ at any given time.

A straightforward UV model that gives rise to Sec. II is one with a (heavy) triplet scalar $\mathcal{Y}$ carrying electric charge

²A similar approach was pursued very recently in Ref. [22] using different mechanisms for changing the ψ_B mass.

$-1/3$ and $B = -2/3$. ($\mathcal{Y}$ can be identified with a squark [19].) Then, we get the following Lagrangian terms:

$$\mathcal{L}_{\mathcal{Y}} = -\sum_{i,j} y_{u_i, d_j} \mathcal{Y}^* \bar{u}_{iR} d_{jR}^c - \sum_k y_{\psi d_k} \bar{\psi}_B \mathcal{Y} d_{kR}^c + \text{H.c.} \quad (3)$$

Of course, just as squarks, any color-triplet scalar is typically constrained by the LHC to be well above the weak scale $\sim 100 \text{ GeV}$. Let us first explain the generation of the baryon asymmetry. Afterward, we will discuss the collider and B -physics constraints which will then guide us in choosing which specific operators to use among $\mathcal{O}_{d_k, u_i, d_j}$.

The model first requires a scalar particle Φ that decays at a temperature below the QCD scale, $T_d < T_{\text{QCD}} \sim 100 \text{ MeV}$, creating out-of-equilibrium B mesons. B mesons oscillate, giving rise to CP violation when they decay. The effective operators in Sec. II facilitate the exotic decay $B_{d,s}^0 \rightarrow \bar{\psi}_B \mathcal{B}_{\text{SM}}$, into the dark antibaryon and a SM baryon $\mathcal{B}_{\text{SM}}$. These decays generate an equal and opposite baryon asymmetry between the dark and SM sectors. (Afterward, ψ_B decays into the stable dark matter (DM). This is discussed in Sec. IV.) The resulting baryon asymmetry is calculated to be [18,19]

$$Y_B \simeq 5 \times 10^{-5} \sum_{i=d,s} [\text{Br}(B_i^0 \rightarrow \bar{\psi}_B \mathcal{B}_{\text{SM}}) A_{sl}^i] \alpha_i(T_d), \quad (4)$$

where

$$A_{sl}^d|_{\text{SM}} = (-4.7 \pm 0.4) \times 10^{-4}, \quad (5a)$$

$$A_{sl}^s|_{\text{SM}} = (2.1 \pm 0.2) \times 10^{-5}, \quad (5b)$$

are the semileptonic asymmetries in B^0 decays generated by the SM CP violation [25]. The factor $\alpha_i(T_d) \in [0, 1]$ accounts for decoherence effects and depends on the temperature at which B mesons were created through the scalar Φ decays. It is calculated by solving the Boltzmann equations for the B -meson density matrix in the early Universe [18].

From Eq. (1) and Sec. II, it is obvious that we need $\text{Br}(B_i^0 \rightarrow \bar{\psi}_B \mathcal{B}_{\text{SM}}) \sim 1$ if we want to use the SM CP violation in Eq. (5). However, recast LEP constraints on the Wilson coefficients $\mathcal{C}_{d_k, u_i, d_j}$ [26] and dedicated mesogenesis searches by Belle-II [27] and BABAR [28] collaborations put stringent limits on these branching fractions, $\text{Br}(B_i^0 \rightarrow \bar{\psi}_B \mathcal{B}_{\text{SM}}) \lesssim 10^{-4}$. (See Tables I and II in Ref. [7] for a summary of all available constraints.)

One way to utilize the SM CP violation in mesogenesis while still being safe from branching fraction constraints is to only allow the new B meson decays in the early Universe while shutting them off today. In Ref. [7] this was achieved by changing the mass of the color-triplet scalar mediator $\mathcal{Y}$ between $T \sim O(10 \text{ MeV})$ and now. It was shown that, with

$$M_{\mathcal{Y}} \lesssim 620 \text{ GeV} \sqrt{y_{\psi s} y_{cb}} \quad (\text{successful mesogenesis}), \quad (6)$$

the operator $O_{s,cb}$ can generate the correct BAU using only the SM CP violation when only B_s mesons are the source [20]; see Fig. 1 of [7]. We will focus on this operator as it allows for the heaviest $\mathcal{Y}$ given a set of Yukawa couplings. (All other operators require much lighter $\mathcal{Y}$ masses to generate the BAU.) Below, we discuss how such a light $SU(3)$ -triplet scalar could be allowed.

III. COLLIDER CONSTRAINTS AND OPERATOR SELECTION

In the following section, we describe a model in which ψ_B mass changes from a few GeV in the early Universe to $\gtrsim 10$ GeV today. With this final mass, the B meson decays into ψ_B final states are not allowed. Hence, the previously mentioned Belle-II [27] and BABAR [28] constraints do not apply.

In Ref. [26], constraints on the mediator mass $M_{\mathcal{Y}}$ and the coupling constants $y_{u,d_j}, y_{\psi d_k}$ were calculated for single $\mathcal{Y}$ production followed by $\mathcal{Y}$ decays to either multijet or jets + missing transverse energy (MET) final states, assuming a single operator, of the form $\mathcal{Y}^* \bar{u}_i d_j^c$, is turned on at a time. (See Fig. 1 for example channels.) Given that the model produces the baryon asymmetry through B meson decays, the analysis only included operators that involve the b quark. A more recent (effective field theory) analysis was carried out in Ref. [29]. For the $O_{s,cb}$ operator mentioned above, the LHC constraints on jets + MET on single $\mathcal{Y}$ channels translate to

$$M_{\mathcal{Y}} \geq 2.4 \text{ TeV} \sqrt{y_{\psi s} y_{cb}} \quad (\text{bjet} + \text{MET}), \quad (7)$$

which is in clear conflict with Eq. (6). ($O_{s,cb}$ is O_{232}^{uddN} in [29].) The above limit applies to $m_{\psi} \lesssim 7.5$ GeV, in which case ψ_B is collider stable and considered as missing energy. If ψ_B is heavier, it can decay into hadron final states with a width approximately given by [29]

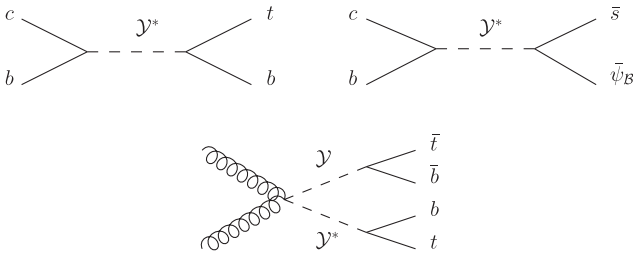


FIG. 1. Some production and decay channels for the $SU(3)$ -triplet scalar mediator $\mathcal{Y}$ at the LHC.

$$\begin{aligned} \Gamma(\psi_B \rightarrow cbs) &\approx \frac{(y_{\psi s} y_{cb})^2 m_{\psi}^5}{1024 \pi^3 M_{\mathcal{Y}}^4} \\ &\simeq 2 \times 10^{-11} \text{ GeV} (y_{\psi s} y_{cb})^2 \\ &\quad \times \left(\frac{m_{\psi}}{10 \text{ GeV}} \right)^5 \left(\frac{600 \text{ GeV}}{M_{\mathcal{Y}}} \right)^4. \end{aligned} \quad (8)$$

The lack of missing energy relaxes the constraint to $M_{\mathcal{Y}} \gtrsim 1 \text{ TeV} \sqrt{y_{\psi s} y_{cb}}$. For $7.5 \text{ GeV} \lesssim m_{\psi} \lesssim 10 \text{ GeV}$, displaced vertex searches are projected to be sensitive to such final states. Above $m_{\psi} = 10 \text{ GeV}$, multijet constraints give³

$$M_{\mathcal{Y}} \gtrsim 900 \text{ GeV} \sqrt{y_{\psi s} y_{cb}} \quad (\text{multijet}), \quad (9)$$

which is still in conflict with Eq. (6).

It is reasonable that two or more of the operators in Sec. II will be present in a UV model. In that case, there are more stringent constraints coming from $\Delta F = 2$ processes that affect the mixing parameters in neutral meson systems. However, these only apply to different operators that couple to the same quark, e.g. $|y_{td} y_{ts}^*| < 3 \times 10^{-4} (M_{\mathcal{Y}}/1.5 \text{ TeV})$ coming from the kaon mixing parameter ϵ_K [see Eq. (43) in Ref. [26]].

With these obstacles in mind, we use the following Lagrangian

$$\mathcal{L}_{\text{new}} = -y_{\psi s} \mathcal{Y} \bar{\psi}_B s_R^c - y_{cb} \mathcal{Y}^* \bar{c}_R b_R^c - y_{td_i} \mathcal{Y}^* \bar{t}_R d_{iR}^c + \text{H.c.} \quad (10)$$

for $d_i = d, s, b$. We first note that the addition of top quark interactions do not alter the baryogenesis scenario even with the $\mathcal{Y} \bar{t} b$ operator since the B mesons cannot kinematically decay to a top quark and chain decays would be suppressed by $(m_b/m_t)^2$. In order to relax the collider constraints, we require that $y_{td_i} \gg y_{cb}, y_{\psi s}$. For example, choosing $y_{td_i} \simeq 5 y_{cb}$ can lower the branching fraction of $\mathcal{Y} \rightarrow cb$ decays to less than 5%, relaxing the relevant constraints from this channel enough to allow for $M_{\mathcal{Y}} \sim 600 \text{ GeV}$.⁴ Note that allowing for all three channels, $\bar{t} d, \bar{t} s, \bar{t} b$ suffers from the flavor constraints from $\Delta F = 2$ processes described above. Hence, only one such coupling can be large and the rest needs to be suppressed.

When the dominant interaction is through the $\mathcal{Y} \bar{t} b$ operator, its contribution to single production of $\mathcal{Y}$ is

³We take this limit from O_{232}^{uddN} as shown in Fig. 8 of [29]. The operator of interest, O_{232}^{uddN} , is expected to be less constrained.

⁴In order to keep the same exclusion strength, assuming the s -channel contributions corresponding to a resonance dominate, a rough estimate is $(\text{Br} \times \sigma)_{900 \text{ GeV}} = (\text{Br} \times \sigma)_{600 \text{ GeV}}$, where $\sigma \propto \text{PDF}(c) \text{PDF}(b) / M_{\mathcal{Y}}^4$. A conservative estimate is that the relevant parton distribution functions (PDFs) roughly double going from 900 to 600 GeV. Hence, the new branching fraction to the cb final states need to be $\text{Br} \lesssim (600 \text{ GeV}/900 \text{ GeV})^4 / 4 \sim 0.05$.

heavily suppressed due to small top quark PDFs. Although in our case $\mathcal{Y}$ can still be produced through the cb channel, the dominant final state is now multijets, instead of jets + MET, which are much less constrained. Of course, as an $SU(3)$ -triplet scalar, $\mathcal{Y}$ can be pair-produced through its gluon couplings. There are several LHC searches on pair production of stops with R-parity violating (RPV) decays. The most relevant one for this case is a CMS experiment search [30] that excludes stop masses up to 670 GeV (for 100% branching fractions into top final states). (See Ref. [31] for a recent review of RPV searches.)

We also note that by adding a large coupling constant, $y_{td_i} \sim 5$, we are effectively broadening the $\mathcal{Y}$ width. As the width of the resonance increases, the background (in a collider search) that one integrates over also increases, hence lowering the signal-to-background ratio. Hence, we expect that the above constraint could relax to $M_{\mathcal{Y}} \simeq 620$ GeV from 670 GeV. In Ref. [32] broad and narrow resonance scenarios were discussed for particles with masses $O(\text{TeV})$ with widths up to $\Gamma/M = 30\%$, showing that the background could increase by an order of magnitude from a narrow resonance scenario. The effect is less prominent for smaller masses. It would be interesting to do a dedicated LHC analysis for this particular scenario, with large Yukawas, which we leave for future work.

IV. A MODEL FOR COSMOLOGICAL EVOLUTION OF DARK BARYON MASS

In this section, we provide a simple model for making the mass of the dark sector fermion ψ_B dependent on the ambient SM particle density. In particular, we will make m_{ψ} dependent on the number density of charged leptons $\ell = e, \mu, \tau$. Given the typical temperatures for mesogenesis $T_M \sim O(10 \text{ MeV})$, this is mostly achieved via the $e^{\pm}$ and $\mu^{\pm}$ thermal populations. To do this, we introduce an ultralight scalar φ of mass m_{φ} that couples to ψ_B and charged leptons

$$(m_{\psi}^0 + g_{\psi}\varphi)\bar{\psi}_B\psi_B + \sum_{\ell=e,\mu,\tau} g_{\ell}\varphi\bar{\ell}\ell, \quad (11)$$

where g_{ψ} and g_{ℓ} are the couplings of φ to ψ_B and ℓ , respectively. Here, m_{ψ}^0 is the mass of ψ_B in vacuum, corresponding to $\varphi = 0$. Given the above setup, the mass of ψ_B would have a dependence on φ via

$$m_{\psi} = m_{\psi}^0 + g_{\psi}\varphi. \quad (12)$$

We would like to lower the mass of ψ_B for $T_M \sim 10 \text{ MeV}$ so that $\text{Br}(B_i^0 \rightarrow \bar{\psi}_B \mathcal{B}_{\text{SM}})$ can proceed effectively during mesogenesis, but the mass later reaches its larger vacuum value, shutting off these decays after mesogenesis is over. In order for mesogenesis to proceed via our proposed operator, a prompt $B_s \rightarrow \psi_B \Omega_c^0$ decay must occur.

This can only happen when $m_{\psi} < m_{B_s} - m_{\Omega_c^0} = 2.7 \text{ GeV}$. We also require a prompt decay of ψ_B to dark sector particles ϕ, ξ to shut off back-reactions and possible baryon number washout, $m_{\psi} > m_{\phi} + m_{\xi} \gtrsim 1.5 \text{ GeV}$ [18]. To safely accommodate these requirements, we conservatively require $m_{\psi} < 2.5 \text{ GeV}$ during mesogenesis, and we choose a benchmark such that $m_{\psi} > 2 \text{ GeV}$ for its entire evolution.

To achieve a reduction of the ψ_B mass, we assume $g_{\psi}, g_{\ell} > 0$. We note that ψ_B eventually decays to DM particles—a scalar ϕ and a fermion ξ —via an interaction

$$y_{\psi}\phi\bar{\psi}_B\xi + \text{H.c.} \quad (13)$$

Through these interactions, ψ_B can decay to the dark sector faster than washout processes that would repopulate B mesons. To see this, note that the Hubble rate is given by $H(T) = \sqrt{8\pi^3 g_*(T)/90} T^2 / M_{\text{P}}$, where $g_*(T)$ is the number of relativistic degrees of freedom at temperature T and $M_{\text{P}} \approx 1.22 \times 10^{19} \text{ GeV}$ is the Planck mass. $T \sim 10 \text{ MeV}$ corresponds to timescales of $O(0.01 \text{ s})$. The decay width for $\psi \rightarrow \xi\phi$, ignoring final state masses, is roughly given by

$$\Gamma(\psi_B \rightarrow \xi\phi) \approx \frac{y_{\psi}^2 m_{\psi}}{32\pi}. \quad (14)$$

Hence, for a ψ_B of mass $\sim$ a few GeV, it only needs to couple with a strength $y_{\psi} \gg 10^{-10}$ to DM final states in order to decay fast on the relevant cosmological time scales. For $y_{\psi} \lesssim 10^{-5}$, the prompt cosmological decay of ψ_B into DM is compatible with its late time decay dominantly into hadrons.

The initial value of φ is sourced by the lepton thermal population and once the Hubble rate falls roughly below $m_{\varphi}/3$, the scalar begins to oscillate and its amplitude redshifts as $T^{3/2}$. As φ changes, the mass of the dark fermion ψ_B also changes. Since we require that m_{ψ} stays small until $T \lesssim 10 \text{ MeV}$, we need the φ oscillations to be delayed until after mesogenesis as well. Therefore, the mass of φ should generally not be much heavier than $m_{\varphi} \sim O(10^{-13} \text{ eV})$, making φ an ultralight scalar, with a Compton wavelength exceeding $\sim 10^3 \text{ km}$. For lighter φ , the oscillations start later and hence the increase in m_{ψ} happens later as well. This might result in a period of time where ψ_B becomes lighter than the proton allowing proton decay. This effect can give a lower bound on the φ mass. To be conservative, we will take $m_{\varphi} \sim O(10^{-13} \text{ eV})$ as a safe benchmark value.

Thermal evolution of φ is governed by its coupling to particles in the primordial plasma. In our case, these are charged leptons ℓ . In the regime where ℓ is relativistic, the evolution of φ is given by (see, for example, Ref. [33])

$$\frac{T^6}{4M_*^2} \frac{d^2\varphi}{dT^2} + \left(m_{\varphi}^2 + \frac{g_{\ell}^2}{6} T^2 \right) \varphi + \frac{g_{\ell}}{6} m_{\ell} T^2 = 0, \quad (15)$$

where

$$M_* \equiv \frac{M_P}{2} \sqrt{\frac{90}{8\pi^3 g_*(T)}}. \quad (16)$$

The corresponding equation for a nonrelativistic lepton ℓ is [33]

$$\frac{T^6}{4M_*^2} \frac{d^2 \varphi}{dT^2} + \left[m_\varphi^2 - \frac{g_\ell^2 n_{0\ell}}{T} \left(1 - \frac{3T}{2m_\ell} \right) \right] \varphi + g_\ell n_{0\ell} = 0, \quad (17)$$

where

$$n_{0\ell} = g \left(\frac{m_\ell T}{2\pi} \right)^{3/2} e^{-m_\ell/T}, \quad (18)$$

with $g = 4$ for a Dirac fermion and its antifermion conjugate ($\ell, \bar{\ell}$). Technically, in these equations of motion, ψ_B interactions and decays could also be included. However, as discussed above, ψ_B decays to DM particles very fast compared to the other time scales. After these decays, there will be a nonthermal density of ξ (and ϕ) particles which can also shift the φ field value through loop interactions. These effects are also negligible; see discussions around Eqs. (21) and (23).

We will take $m_\varphi = 9 \times 10^{-14}$ eV as a reference value which suffices to show the feasibility of the mechanism. We would like $m_\psi^0 + g_\psi \langle \varphi \rangle$ to be less than a few GeV during mesogenesis, but larger than ~ 5 GeV today in order to shut off two-body B meson decays into an on-shell ψ_B and an SM baryon.

A φ coupling to muons of $g_\mu = 10^{-24}$ is unconstrained by about 5 orders of magnitude after accounting for loop processes that can couple φ to electrons and protons⁵ [34–38]. Even though the muons are quite Boltzmann suppressed at $T \sim T_M$, we still find a relatively large background of φ that is sourced by them.

As can be seen in Fig. 2, at $T \sim T_M$ we have $\varphi \sim 10^{16}$ MeV. Hence, if we want to change the mass of ψ_B by $\mathcal{O}(10^3)$ MeV, we need to choose $g_\psi \sim 10^{-13}$. Note that with our reference choices above, φ starts to oscillate around $T_1 \approx 4.8$ MeV (first peak in Fig. 2) with an amplitude of $\varphi \approx 8.4 \times 10^{15}$ MeV, corresponding to an energy density

$$\rho_\varphi(T_1) = \frac{1}{2} m_\varphi^2 \varphi^2 \approx 2.8 \times 10^{17} \text{ eV}^4, \quad (19)$$

⁵The two-loop diagram that couples φ to a proton has been recently explicitly calculated in [34], which yields $g_p = \alpha^2 g_\mu / 2 \sim 3 \times 10^{-5} g_\mu$. The current upper bound on proton coupling to a long range force is $g_p \sim 10^{-24}$ [35,36].

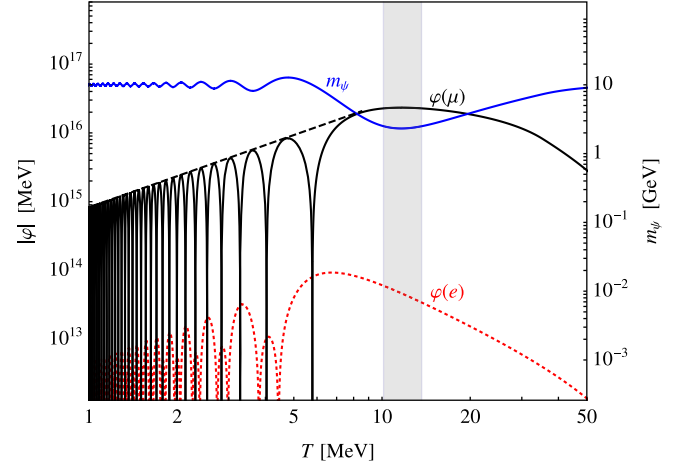


FIG. 2. Evolution of the absolute value of φ via muon interactions only (black/solid curve, in MeV) and electron interactions only (red/dotted) with temperature (in MeV). We also show the variation of m_ψ (blue curve, in GeV) with temperature. We have set $g_\psi = 3.3 \times 10^{-13}$, $g_\mu = 10^{-24}$, $g_e = 10^{-25}$, $m_\varphi = 9 \times 10^{-14}$ eV, $m_\psi^0 = 10$ GeV. The dashed line is the envelope of the φ amplitude that redshifts like $T^{3/2}$, as expected for a component of cold dark matter. The gray shaded region marks the interval where mesogenesis can take place, corresponding to the regime where a B meson can decay into ψ_B on-shell.

which will consequently redshift as T^3 . Near the matter-radiation equality at $T_{\text{eq}} \approx 1$ eV, we get $\rho_\varphi(T_{\text{eq}}) \approx 2.6 \times 10^{-3} \text{ eV}^4$, well below the radiation energy density of $\mathcal{O}(\text{eV}^4)$. Hence, with the above choices for our parameters we will not overclose the Universe, and φ will be a small subcomponent of DM.

Here, we note that we could have also used the coupling of φ to electrons to induce $\langle \varphi \rangle \neq 0$ in the thermal bath. However, this contribution would be subdominant to that of the muon under generic assumptions, that is, if we assume a coupling to the muon that is not smaller than the electron's, which one may expect on general grounds. The coupling of a long range scalar to electrons is constrained to satisfy $g_e \lesssim 10^{-25}$ [35–38]. (Although this is a tighter constraint than the one on g_p , it does not give a correspondingly tighter bound on g_μ because a similar two-loop interaction is chirality suppressed [34].) As can be seen in Fig. 2, including the electron contribution, according to Eq. (15), would only result in a negligible contribution to the shift of the ψ_B mass. This is also partly due to the suppressed effect of a relativistic source for φ . Qualitatively, this can partly be understood as follows. A relativistically invariant fermion source for φ can be formulated by the product of its number density n_f and its inverse boost factor $\sqrt{1 - v_f^2}$, where v_f is a typical velocity. Hence, electrons have a suppressed effect at high temperatures (see, for example, Ref. [39]).

One may worry that the much larger coupling of φ to ψ_B may lead to a back reaction and drive φ to values larger than depicted in Fig. 2. This could push the mass of ψ_B to small values and hamper its decay into ϕ, ξ through the interactions given in Eq. (13). However, taking $y_\psi \gg 10^{-10}$, we have assumed that the decay of ψ_B is prompt on the cosmological evolution time scales and a ψ_B would quickly decay into DM states.

A loop generated interaction can still couple φ to DM, for example $g_\xi \varphi \bar{\xi} \xi$. The loop-induced coupling can be roughly estimated by

$$g_\xi \sim \frac{g_\psi y_\psi^2}{16\pi^2}. \quad (20)$$

One may also want to make sure that for our benchmark parameters the induced φ couplings to DM in Eq. (20) are sufficiently small in order to avoid a large back-reaction on the misalignment of φ . To check this, let us assume the initial decaying modulus Φ , which produced the required B mesons, reheated the Universe to $T_M \sim 10$ MeV and that $\text{Br}(\Phi \rightarrow B\bar{B}) = 1$. We then roughly have

$$m_B n_B \sim T_M^4 \Rightarrow n_B \sim 10^{-3} T_M^3. \quad (21)$$

Given that we want an $\mathcal{O}(1)$ fraction of B mesons to decay into the dark sector during mesogenesis, we infer that $n_\xi \sim 10^{-3} T_M^3$, initially.

To make sure that the initial population of DM (before annihilation down to its asymmetric relic density) does not dominate the misalignment of φ and push m_ψ below $\mathcal{O}(\text{GeV})$, we demand

$$g_\psi \delta\varphi \lesssim m_\psi^0 \Rightarrow \frac{g_\psi g_\xi n_\xi}{m_\varphi^2} \lesssim m_\psi^0, \quad (22)$$

where $\delta\varphi$ is the contribution of DM state ξ . For $m_\varphi \sim 10^{-13}$ eV and $m_\psi^0 \sim 10$ GeV, as in Fig. 2, and using Eq. (21) and (20), we get

$$g_\psi y_\psi \lesssim 10^{-16}. \quad (23)$$

Hence, for our benchmark value of $g_\psi \sim 3 \times 10^{-13}$, the above relation is readily satisfied with $y_\psi \lesssim 10^{-5}$, which is also required to relax the LHC bounds we discuss in the next section. See also the discussion below Eq. (14).

V. POSSIBLE SIGNALS

In this section we discuss possible signals for this mechanism beyond LHC searches for color-triplet scalars. We focus instead on signals for the dark sector, such as displaced vertices, as well as long-range muonic forces arising from the ultralight scalar. Although we will not discuss it here, another signal can come from the physics

facilitating the annihilation of the symmetric population of DM.

A. Potential collider signals

In previous works on mesogenesis, where ψ_B is light, exotic decays of B mesons were extensively studied. The branching fraction that can be constrained is given as [18]

$$\begin{aligned} \text{Br}(B \rightarrow \psi_B + \text{baryon}) \\ \simeq 0.5(y_{cb} y_{\psi s})^2 \left(\frac{m_B - m_\psi}{3 \text{ GeV}} \right)^4 \left(\frac{600 \text{ GeV}}{m_\psi} \right)^4. \end{aligned} \quad (24)$$

Belle II has the potential to probe branching fractions $\text{Br}(B \rightarrow \psi_B + \text{baryon}) \sim 3 \times 10^{-6}$ [26].

When the mass of the ψ_B is too heavy to be a final state in a $B \rightarrow \text{inv.} + \mathcal{B}_c$ decay, there is still a possibility of the three-body $B \rightarrow \psi_B^* + \mathcal{B}_c \rightarrow \xi\phi + \mathcal{B}_c$, where the ψ_B^* is off-shell. (Here $\mathcal{B}_c$ is a charm baryon.) Compared to the above estimate, this decay rate will be suppressed by y_ψ^2 , an extra phase space factor of $8\pi^2$ and a factor of $(m_B/m_\psi^0)^2 \lesssim 1/4$. [Note also that the mass difference $(m_B - m_\psi)$ will be replaced by $m_B - (m_\xi + m_\phi)$.] This exotic decay can be observable in the future for $y_\psi \gtrsim 10^{-2}$. However, as we require $\psi_B \rightarrow \text{jets}$ to be dominant in order to relax the LHC constraints [29], we typically need $y_\psi \lesssim 10^{-5}$, as discussed earlier, making the missing energy channel unobservable.

Although we take $m_\psi = 10$ GeV in this work as a benchmark value, masses around $7.5 \text{ GeV} \lesssim m_\psi \lesssim 10 \text{ GeV}$ have been identified in [29] as a potential target for displaced vertex searches at the LHC. This mass regime can easily be accommodated in our scenario by only choosing $g_\psi \sim 2 \times 10^{-13}$.

B. Long range muonic force probes

For values of $g_\mu \sim 10^{-18}$, experiments searching for new long range forces can in principle be sensitive to our parameter space [35–38]. An intriguing possibility is that a long-range muonic force can be probed by gravitational wave signal resulting from neutron star mergers [40,41]. This is due to the presence of large numbers of muons, $N_\mu \sim 10^{55}$ (see, for example, Ref. [42]) in the dense environment of a neutron star. It is estimated that LIGO could probe long-range muonic forces with couplings below 10^{-20} , but a robust bound is subject to LIGO analysis.

A question related to the above setup is to what extent one can expect the mass of ψ_B to be lowered within a neutron star due to the significant muon abundance in the stellar environment. This could in principle lead to proton decay, if ψ_B becomes lighter than a nucleon. A simple estimate for the background value of the scalar in the star is

$$\varphi_{\text{NS}} \sim \frac{g_\mu N_\mu}{4\pi R} \sim g_\mu 10^{37} \text{ MeV}, \quad (25)$$

where $R \sim 10$ km is a typical length scale for the neutron star.

Since the mass of ψ_B is given by $m_\psi^0 + g_\psi \varphi$, we see that for $m_\psi^0 = 10$ GeV, we need

$$g_\mu g_\psi \lesssim 10^{-33}, \quad (26)$$

assuming $g_\mu, g_\psi > 0$, in order to have $m_\psi \gtrsim 1$ GeV. Therefore, our benchmark values, used in Fig. 2, would not lead to the possibility of nucleon decay in a neutron star.

VI. CONCLUSIONS

What is the SM CP violation good for? Although one of the fundamental conditions for creating a baryon asymmetry is the existence of CP violation in a theory, the effects of the CKM phase get suppressed in electroweak baryogenesis models. Mesogenesis with a morphing mediator model [7] is the only model that utilizes the observed SM CP violation in B meson systems in order to create the baryon asymmetry of the Universe.

Mesogenesis models mainly rely on effective operators given in Sec. II which can arise from the UV interactions in Sec. II between SM quarks and a dark sector fermion ψ_B , which carries baryon number, mediated by an $SU(3)$ triplet scalar $\mathcal{Y}$. On the one hand, in order for the SM CP violation to be enough, the effective interactions need to be sufficiently large. On the other hand, large enough interactions with quarks are heavily constrained by the LHC and, specifically for mesogenesis, by B factories. In order to alleviate these constraints, previous work [7] explored a mechanism where the mediator mass changes from ~ 600 GeV to a TeV after the baryon asymmetry is created. However, this realization of mesogenesis, where the baryon asymmetry is created in certain domains with a light mediator between the SM and a dark sector, is argued to suffer from BBN constraints due to inhomogeneities [21].

In this work, we introduced an alternative mechanism to realize mesogenesis using still only the SM CP violation. Instead of focusing on the mediator, we developed a model in which the dark sector fermion ψ_B , in vacuum, is heavier than the B meson but becomes temporarily light enough to allow for mesogenesis in the early Universe. In this scenario, since ψ_B is heavy currently, exotic B meson decay channels are shut off at colliders. This mass evolution is realized via couplings to an ultralight scalar whose initial field value is driven by the thermal number density of SM charged leptons [24]. As benchmark values, we show that a

scalar with mass $\mathcal{O}(10^{-13}$ eV) and coupling constants of $10^{-10} - 10^{-24}$ are able to realize this scenario. In this work, we focused on muon-induced thermal misalignment, but one could alternatively employ electrons, for different choices of parameters.

Here, we also included multiple terms from the mesogenesis Lagrangian in Sec. II in order to dilute the branching fraction of the mediator $\mathcal{Y}$ into quark and missing energy final states. Successful mesogenesis, if we insist on using only the SM CP violation, requires $M_{\mathcal{Y}} \lesssim 600$ GeV. We examined several relevant LHC searches to argue that with couplings $y_{cb} \sim 1, y_{tb} \sim 5$ such a light $SU(3)$ triplet scalar could still be viable. These are very sizable couplings, giving rise to a broad resonance. It is interesting to explore further search strategies at the LHC for such a scenario, which we leave for future work. However, we expect that the underlying physics mediator $\mathcal{Y}$ would be within the reach of the LHC, as more data get analyzed and accumulated.

In the benchmark scenario we considered, $m_\psi = 10$ GeV, the dark fermion ψ_B decays to jets at the LHC. For slightly lighter ψ_B , these decays can produce potentially detectable displaced vertex signals. The ultralight scalar can also be uncovered in terrestrial searches for long-range forces as well as neutron star mergers.

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DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

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