

Baryogenesis from B meson oscillations

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We show how CP violating B meson oscillations in conjunction with baryon number violating decays can generate the cosmological asymmetry between matter and antimatter, and explore the parameter space of a simple, self-contained model, which can be tested via exotic B meson decays, and via the charge asymmetry in semileptonic decays of neutral B mesons.

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I. INTRODUCTION

Baryogenesis—generating the cosmological asymmetry between matter and antimatter—requires physics beyond the Standard Model (SM). The pioneering work of Sakharov [1] found three necessary conditions: baryon number violation, C and CP violation, and departure from thermal equilibrium. Baryon number violation occurs nonperturbatively in the Standard Model [2]. CP violation also occurs, however the standard model CP violation appears to be too small to explain the observed baryon asymmetry. Finally, the minimal Standard Model contains no mechanism for departure from thermal equilibrium.

Recent work [3,4] has shown the possibility for low energy baryogenesis via the oscillations of neutral hadrons, in conjunction with new sources of CP and baryon number violation. In Ref. [4], the oscillating hadrons were *mesinos*—bound states of a quark and an antiquark. In that work a relatively long-lived squark decayed into antiquarks via baryon number violating R -parity violating decays. A minimal model to capture this physics was studied in detail—that model contained three neutral Majorana fermions (“neutralinos”) and a color triplet scalar (“squark”). The same model, in a different parameter region with lighter neutralinos, was shown to lead to baryogenesis via potentially observable baryon and CP violating neutral heavy flavor baryon oscillations [3,5]. A similar model, in which baryon number is conserved but also carried by dark matter, was shown to be capable of producing both the

visible matter-anti-matter asymmetry and asymmetric dark matter [6] via B -meson decays. In the present work, we reexamine the simpler model of Ref. [4], and show that for a different parameter range that was not considered in the previous work, baryon number violating decays of B^0 mesons are allowed by experiment, potentially observable, and could be the explanation for baryogenesis.

The baryogenesis scenario described here begins in the pre-nucleosynthesis early universe with the decays of a long lived scalar into b -quarks and antiquarks. These decays are assumed to take place late enough and at low enough temperature to allow hadronization, but before nucleosynthesis. Most of the b -quarks form B mesons. The neutral B mesons then oscillate and decay, sometimes to baryons or antibaryons, resulting in the observed asymmetry.

Our model is similar to the one used in Ref. [5] to produce the baryon asymmetry through oscillations of baryons [3,4], although even simpler. It is worth noting that even though we use a similar model, the parameter space is different and we are less constrained from dinucleon decay. The mechanism we describe is similar to the one found in a slightly more elaborate model [6], which could generate a dark matter relic along with baryon number. In the current work we have no dark matter sector. We assume that whatever the dark matter is, it is very weakly coupled and has no effect on baryogenesis. It would be a straightforward matter to include, for instance, axion dark matter. Note that the presence of a late decaying heavy particle can have an effect on the allowed axion parameter range and substructure [7].

This paper is organized as following: In Sec. II we give an overview of B meson physics and how it could be related to baryogenesis. A more detailed description of the model and phenomenology is given in Sec. III. In Sec. IV we show that the baryon asymmetry may be produced with parameters which are allowed by experiment. In Sec. V we conclude and sketch ideas for future work.

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II. CHARGE ASYMMETRY IN OSCILLATING B-MESONS AND SIGN OF THE MATTER-ANTI-MATTER ASYMMETRY

We briefly review the physics of neutral B -meson oscillations and describe how new physics in the decays and mixing can lead to baryogenesis. The oscillations are described by an effective 2 state Hamiltonian:

$$\mathcal{H} = M - \frac{i}{2}\Gamma = \begin{pmatrix} M & M_{12} \\ M_{12}^* & M \end{pmatrix} - \frac{i}{2} \begin{pmatrix} \Gamma & \Gamma_{12} \\ \Gamma_{12}^* & \Gamma \end{pmatrix}. \quad (1)$$

Here M and Γ are respectively the dispersive part and the absorptive part of the transition amplitude. The CP violating phase $\arg(\Gamma_{12}/M_{12})$ is reparameterization invariant and observable in the semi-leptonic charge asymmetry of neutral meson decays. This phase is crucial for our baryogenesis mechanism, as it determines whether there are more b quarks or b anti-quarks at the time of decay. With an asymmetry between b quarks and antiquarks, the baryon number violating decays of the b quark into two lighter antiquarks and a Majorana fermion, or the b anti-quark into two quarks and a Majorana fermion can produce the observed matter antimatter asymmetry. In order for baryon number violating decays to produce more matter than antimatter, there must be more b -antiquarks at the time of decay. Because the semi-leptonic decays of b -antiquarks produce positively charged leptons, this baryogenesis mechanism requires a positive charge asymmetry in either B^0 or B_s meson oscillations, or both. The charge asymmetry depends on the phase of the absorptive term (Γ_{12}) relative to the dispersive term (M_{12}) in B meson mixing. The magnitude of Γ_{12} and M_{12} are measured and are in agreement with the Standard Model in both neutral meson systems. In the Standard Model, the phase $\arg(\Gamma_{12}/M_{12})$ is predicted to be very small for both the B^0 and the B_s due to the unitarity of the CKM matrix and the relatively small mass of the charm quark, leading to a very small prediction for the charge asymmetry. Experimentally, the charge asymmetry has not yet been distinguished from zero in either system. The prediction in the standard model is that the charge asymmetry is negative for the B^0 and positive but very small for the B_s . Our detailed analysis will show that because more b quarks fragment into B^0 mesons than B_s mesons, and because of the small size of the standard model asymmetry in B_s mesons, new physics in the mixing amplitude may be needed for baryogenesis from B mesons. New physics which makes a small contribution to M_{12} and is consistent with experimental constraints can have a significant effect on the phase $\arg(\Gamma_{12}/M_{12})$ in either system. When experimental constraints on such new contributions are taken into account, the most promising case for a positive charge asymmetry which is large enough

for baryogenesis is a new contribution to mixing in the B_s system.

III. NEW PHENOMENOLOGY FROM OUR MODEL

A minimal renormalizable model for this baryogenesis mechanism contains three new particles: One Majorana fermion, χ , in the mass range 2–3 GeV, one new charge $-1/3$ color triplet scalar ϕ , with mass in the range 500 GeV to 1.9 TeV, and a very weakly coupled particle, Φ , which decays out of equilibrium into b -quarks and antiquarks. The upper bound on the ϕ mass is from the need for a large enough branching ratio for exotic B meson decay, as we explain later. Φ might be the inflaton or a string modulus. As Φ is very weakly coupled it is not experimentally accessible. Detailed computation shows that we need the mass of Φ to be between 11 GeV and about a hundred GeV. We note that it is to be expected that a scalar in this mass range would mainly decay into b -quarks. A simple way to achieve this goal is coupling Φ to the Higgs boson with a very small coupling constant g : $\Delta\mathcal{L} = g\Phi H^\dagger H$. The vacuum expectation value (VEV) of the Higgs boson would naturally give rise to $g\nu\Phi h$, which results in a tiny amount of Φ - H mixing and a decay channel to b -quarks. The neutral Majorana fermion χ and color triplet ϕ will interact with $SU(2)$ singlet quarks through the following terms:

$$\mathcal{L}_{\text{int}} \supset -g_{ij}\phi^* \bar{u}_R^i d_R^j - y_j \phi \bar{\chi} d_R^j + \text{H.c.} \quad (2)$$

where i and j are flavor indices which run over the three generations of up- and down-type quarks.

The Majorana fermion χ is neutral and its mass should be greater than the mass difference between proton and electron, $m_p - m_e = 937.75$ MeV, otherwise this model will give rise to proton decay [3]. The mass of χ will be taken to be 2–3 GeV in this paper. Mediated by ϕ , the Majorana fermion χ could decay to three quarks. The lifetime must be less than 0.1s to avoid spoiling successful BBN [8]. This requires that $g_{ij}y_j \gtrsim (5 \text{ GeV}/m_\chi)^5 (m_\phi/350 \text{ TeV})^2$ [5], with $i = u$ or c and $j, j' = d$ or s .

The colored scalar ϕ will either decay to antiquark pairs or into χ plus a quark, and could appear in searches for dijet resonances and jets and missing energy. The mass of ϕ should be at least about 500 GeV to pass collider constraints [9]. There will also be many constraints from nucleon oscillations and dinucleon decays [5], which give us upper bounds on various flavor combinations of the g_{ij} and y_j couplings. The strongest bounds, which are on combinations of couplings that allow dinucleon decays [5], requires that $g_{ij}y_j \lesssim (m_\phi/34 \text{ TeV})^2$ for $m_\chi = 2$ GeV, which are consistent with a cosmologically acceptable χ lifetime if the mass of χ is greater than 2 GeV. The upper bound on the χ mass comes from branching ratio of exotic

B meson decay, which will be discussed in the following subsection III A. The properties of new particles we introduce are summarized in Table 1.

Long-lived but unstable particles like Φ and χ are extremely hard to detect because the coupling constant is too small, which appears as missing energy in colliders. But the χ particle is produced in b -quark decays and will eventually decay to SM particles, which gives an opportunity for detection in long lived particle searches [10,11].

This model allows B mesons to decay into a baryon (plus mesons) plus an unstable Majorana fermion. As the latter will equally likely decay into a baryon or an anti baryon plus mesons, when combined with a CP violating charge asymmetry, such decays will result in a net baryon number. The box diagrams involving the new particles and the y_d couplings can also modify the phase $\arg(\Gamma_{12}/M_{12})$, producing a nonstandard charge asymmetry in B oscillations.

A. Exotic B meson decays

There are potentially observable consequences for B meson physics. First, there will be a new decay channel for B mesons, which violates baryon number, and involves a relatively long lived exotic Majorana fermion. The Feynman diagram for this decay process is shown in Fig. 1.

The Majorana fermion will be present as missing energy in most collider searches. Searches for long lived particles,

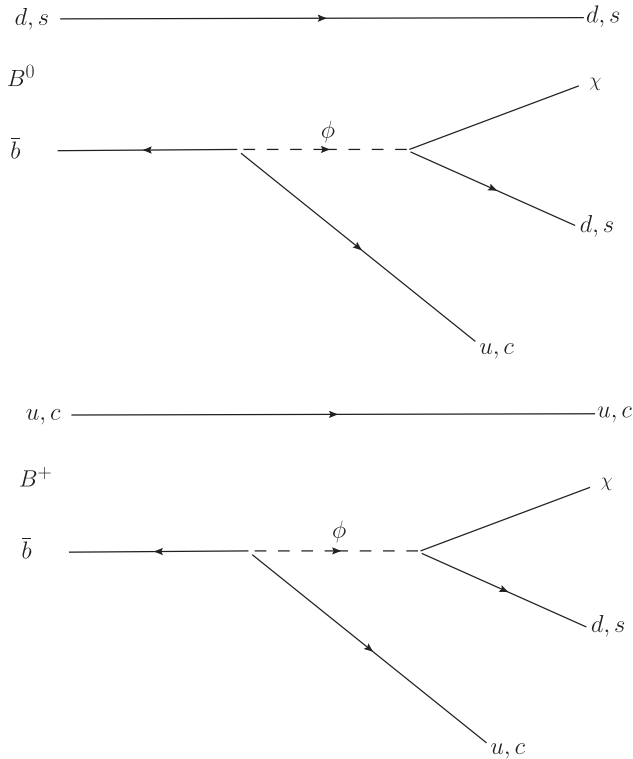


FIG. 1. These are the Feynman diagrams for a new b decay channel, which violate baryon number.

TABLE I. This table briefly summarizes the properties of the new particles we introduced. The masses and lifetimes are mainly constrained by colliders and cosmological observations, which will be discussed in more detail. The colored scalar ϕ will decay rapidly to anti-quark pairs or χ plus a quark through tree-level couplings.

New Particles in Our Model						
	Spin	Mass $z(\text{GeV})$	Lifetime	SU(3)	SU(2)	Q
Φ	0	11–100	0.2–20 ms	Singlet	Singlet	0
ϕ	0	500–1900		Triplet	Singlet	$-1/3$
χ	1/2	2–3	<0.1 s	Singlet	Singlet	0

e.g., using a MATHUSLA-like detector, have a chance to find it [10].

In estimating the branching fraction for the exotic decay of B mesons, we assume that the χ mass is low enough that the momenta in the decay are all larger than the QCD scale so that we may treat QCD perturbatively. Then decay may be approximated by the rate for a heavy b quark to decay into 3 lighter fermions, while the light quark in the B meson acts as a spectator.

The effective interaction term is

$$\Delta\mathcal{L} = \frac{y_s g_{ub}^*}{m_\phi^2} \bar{b} u s \chi. \quad (3)$$

In this limit, neglecting the masses of the light quarks, the decay rate is approximately [5]:

$$\Delta\Gamma \sim \frac{|y_s g_{ub}^*|^2}{60(2\pi)^3 m_\phi^4} m_b \Delta m^4, \quad (4)$$

where Δm is the mass splitting between the χ and the bottom quark. The B^0, B^+ , and B_s could all possibly decay to a baryon plus mesons plus the χ fermion, with the χ appearing either as missing transverse momentum or as a long lived particle. As χ can decay into either 3 quarks or 3 antiquarks, it will appear as either a baryon plus mesons or an antibaryon plus mesons. Note that y_s is not strongly constrained by colliders or by the oscillations of baryons, therefore this branching fraction can be relatively large.

To produce the required baryon number, we will find in Sec. IV that the branching ratio $\text{Br}_{B \rightarrow B}$ must be in the range of $\text{Br}_{B \rightarrow B} \sim 10^{-3} - 10^{-1}$. From the decay rate we have estimated, the branching ratio is

$$\text{Br}_{B \rightarrow B} \sim 10^{-3} \left(\frac{\Delta m}{2 \text{ GeV}} \right)^4 \left(\frac{1 \text{ TeV}}{m_\phi} \frac{\sqrt{y_s g_{ub}}}{0.53} \right)^4. \quad (5)$$

The branching ratio has to be greater than 10^{-3} to generate sufficient baryon number, which imposes an upper bound $m_\phi \lesssim 1.9 \text{ TeV}$ if we take $\Delta m \sim 2 \text{ GeV}$, $y_b g_{us} \sim 1$. This is compatible with collider constraints on a colored scalar,

including the constraints on resonant single ϕ production [12–14]. The signature of exotic B decays is a baryon with missing energy [15].

B. New contributions to $B^0 - \bar{B}^0$ mixing

In our model, baryogenesis can result from baryon number violation in B meson decay. The baryon number will depend on CP violating phases determining the charge asymmetry in B_d and B_s meson mixing. In the SM the CP violating phase in $B^0 - \bar{B}^0$ mixing is determined by the CKM matrix. Our new model could also contribute to $B^0 - \bar{B}^0$. As $|M_{12}|$ is measured with high precision and agrees well with SM predictions, the new contribution should be much smaller than the standard box diagram. It generically will have a different phase. Because the CP violating phase is small, the width difference $\Delta\Gamma_q$ and Δm_q may be used to estimate the magnitude of Γ_{12} , M_{12} , with $\Delta\Gamma_q = 2|\Gamma_{12}|$ and $\Delta m_q = 2|M_{12}|$, where the subscript q represents arbitrary quarks. Δm_q is well measured [16–18]:

$$\begin{aligned}\Delta m_d &= 0.5064 \pm 0.0019 \text{ ps}^{-1}, \\ \Delta m_s &= 17.757 \pm 0.021 \text{ ps}^{-1}\end{aligned}\quad (6)$$

The decay width difference $\Delta\Gamma_s$ is given by [19]:

$$\Delta\Gamma_s = 0.085 \pm 0.015 \text{ ps}^{-1}. \quad (7)$$

$\Delta\Gamma_d$ is not well measured because $\Delta\Gamma_d/\Gamma_d$ is too small and the uncertainties are relatively very large. However, in the approximation of negligible CP violation in mixing, the ratio $\Delta\Gamma_q/\Delta m_q$ is equal to the small quantity $|\Gamma_{12}/M_{12}|$, which is independent of CKM matrix elements and could be used to determine $\Delta\Gamma_d$:

$$\Delta\Gamma_d = 0.0026 \text{ ps}^{-1} \quad (8)$$

The CP asymmetry in semileptonic B decays is defined as:

$$A_{\text{SL}} = \frac{\Gamma[\bar{B}^0(t) \rightarrow l^+ X] - \Gamma[B^0(t) \rightarrow l^- X]}{\Gamma[\bar{B}^0(t) \rightarrow l^+ X] + \Gamma[B^0(t) \rightarrow l^- X]} \quad (9)$$

where X stands for any other particles produced in this inclusive process. This asymmetry would be determined by the relative phase between the absorptive and dispersive parts of the transition amplitude (For calculations of transition amplitude in Standard Model, see [20–22]), $A_{\text{SL}} = \text{Im}\Gamma_{12}/M_{12}$. Assuming only M_{12} receives a new contribution from new physics, the experimental searches for A_{SL} gives a range of [19,23–25]:

$$\begin{aligned}A_{\text{SL}}^d &\in (-5.9 \times 10^{-3}, -4 \times 10^{-4}), \\ A_{\text{SL}}^s &\in (-1.11 \times 10^{-3}, 8.8 \times 10^{-4}).\end{aligned}\quad (10)$$

The SM predictions for A_{SL} of both B_s^0 and B_d^0 are very small [19,26]:

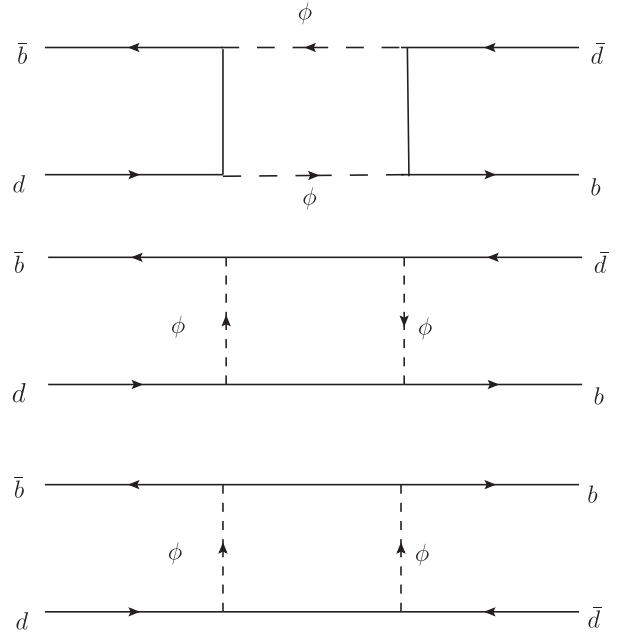


FIG. 2. These are the Feynman diagram for a nonstandard B meson oscillations.

$$\begin{aligned}A_{\text{SL}}^d &= (-4.7 \pm 0.6) \times 10^{-4} \\ A_{\text{SL}}^s &= (2.22 \pm 0.27) \times 10^{-5}.\end{aligned}\quad (11)$$

As mentioned, A_{SL} has to be positive to give rise to baryogenesis, which means that we will need a positive asymmetry in the B_s system. We will find that new physics in the mixing is favored to make this sufficiently positive. There is still room for new physics which could make an order of magnitude change in A_{SL} in either system.

The second term of the Lagrangian, will directly result in an extra contribution to $B^0 - \bar{B}^0$ transition amplitude, which can be seen from the box Feynman diagram 2:

The dispersive part will be given by the exchange of off-shell particles, while the absorptive part of transition amplitude will be given by the exchange of on-shell particles. Since ϕ is heavy, there will not be any absorptive contribution from the new diagrams as long as the mass of χ is greater than half of B meson mass, and even if χ is lighter than this any such contribution would be tiny. For our purpose, we only need to affect the small phase $\arg(\Gamma_{12}/M_{12})$, which does not require any new contribution to Γ_{12} . Any contribution to M_{12} from new physics which has a phase different from the CKM phase could achieve this. The extra transition amplitude would be given by (see Appendix for detailed calculation):

$$\Delta M_{12} = 0.66 \frac{f_B^2 m_B B_B}{16\pi^2 m_\phi^2} (y_b y_d^*)^2 \quad (12)$$

where f_B is the decay constant of B mesons, m_B is B meson mass and B_B is a bag parameter, which is order one. The subscript in y_d^* still represents the down type quark.

It is worth noting that the y_d^i are the only coupling constants from the new model that affect transition amplitude. However, the constraints for the parameters from dinucleon decay and heavy flavor baryon oscillations only set an upper bound for combinations of $g_{ud}^{ij}y_d^k$. From another point of view, the magnitude of transition amplitude is measured with high precision, which agrees with SM predictions roughly and that imposes constraints on combinations of y_d :

$$\begin{aligned} |y_b y_d^*|^2 &\lesssim 1.02 \times 10^{-6} \left(\frac{m_\phi}{1 \text{ TeV}} \right)^2, \\ |y_b y_s^*|^2 &\lesssim 7.33 \times 10^{-7} \left(\frac{m_\phi}{1 \text{ TeV}} \right)^2. \end{aligned} \quad (13)$$

There are other constraints from $\Delta F = 1$ and $\Delta F = 2$ observables [15]. Those constraints will not affect this model because we only need $y_s g_{ub}$ to be large enough. Noting that $(y_b y_s^*)^2$ is generally a complex number, if it has large phase difference with transition amplitude from SM predictions, this term could give rise to positive value of semileptonic asymmetry. The B meson transition amplitude gives the strongest constraint on the combinations $(y_b y_d^*)^2$ and $(y_b y_s^*)^2$, therefore we may fit these parameters to the baryon asymmetry.

IV. COSMOLOGICAL PRODUCTION OF THE BARYON ASYMMETRY

The massive particles, Φ , which mainly decay out of equilibrium to $b\bar{b}$ quarks, provide the necessary departure from thermal equilibrium to produce the baryon asymmetry. The $b\bar{b}$ quarks will quickly hadronize then decay into other lighter particles. Since the time scale for $b\bar{b}$ quark decay is much shorter than the age of the universe, in considering the thermal evolution of the universe it is valid to neglect the brief existence of B mesons or baryons and consider the Φ decay products to be radiation. The radiation produced by Φ decays will thermalize on a timescale that is very short compared with the lifetime of the Φ . Therefore, the evolution of the energy density of radiation and Φ can be described by the following equations:

$$\begin{aligned} \frac{d\rho_\Phi}{dt} + 3H\rho_\Phi &= -\Gamma_\Phi \rho_\Phi \\ \frac{d\rho_r}{dt} + 4H\rho_r &= \Gamma_\Phi \rho_\Phi \end{aligned} \quad (14)$$

where ρ_Φ is the energy density of Φ , ρ_r is the energy density of radiation, and Γ_Φ is the decay rate of Φ . The energy density of ρ_r could directly determine the temperature of the radiation through:

$$\rho_r = \frac{\pi^2}{30} g_*(T) T^4 \quad (15)$$

where $g_*(T)$ is the effective number of degrees of freedom. The Hubble parameter is given by:

$$H = \sqrt{\frac{8\pi\rho_r + \rho_\Phi}{3M_{\text{pl}}^2}} \quad (16)$$

where $M_{\text{pl}} = 1.22 \times 10^{19}$ GeV is the Planck mass. Solving these equations numerically we can obtain the thermal history of the early universe during the early matter dominated era.

When the temperature is below a scale of order $\Lambda_{\text{QCD}} \sim 200$ MeV, the quarks will hadronize. The b quarks produced by Φ decay will mainly form B_d , B^\pm and B_s mesons. The fragmentation ratio of b -quarks to B_d^0 , B^\pm and B_s^0 is taken to be 4:4:1, which is roughly consistent with observation in Z decays and $p\bar{p}$ collisions [27] (Note that the ratio is production-mode dependent with a slightly higher ratio of B_s mesons produced in $p\bar{p}$). The charged B mesons play no role in baryogenesis, while the neutral B mesons, as described in Sec. III, will undergo CP and flavor changing oscillations, and also sometimes decay into baryons and anti-baryons, as shown in Fig. 1.

Since B_d^0 and \bar{B}_d^0 oscillate coherently while also potentially undergoing decoherence from scattering, a density matrix treatment is useful for treatment of these states in the corresponding Boltzmann equations. Accounting for the interaction with plasma and the annihilation between B_d^0 and \bar{B}_d^0 , the Boltzmann equations are [28–30]:

$$\begin{aligned} \frac{dn}{dt} + 3Hn &= -i(\mathcal{H}n - n\mathcal{H}^\dagger) - \frac{\Gamma_\pm}{2} [O_\pm, [O_\pm, n]] \\ &\quad - \langle \sigma v \rangle_\pm \left(\frac{1}{2} \{n, O_\pm \bar{n} O_\pm\} - n_{\text{eq}}^2 \right) \\ &\quad + \frac{1}{2} \frac{\Gamma_\Phi \rho_\Phi}{m_\Phi} \text{Br}_{\Phi \rightarrow B} O_+ \end{aligned} \quad (17)$$

where the last term describes B_d^0 and \bar{B}_d^0 production during the decay of Φ .

Here $\text{Br}_{\Phi \rightarrow B}$ is the branching ratio for $\Phi \rightarrow B + X$, which is assumed to be one because Φ mainly decays to b quarks and \bar{b} quarks mainly hadronize to light B mesons B_d^0 and B_s^0 . This is actually the very reason that we need B meson oscillations to explain baryogenesis from the theoretical point of view. In this equation n and \bar{n} are density matrices,

$$n = \begin{pmatrix} n_{BB} & n_{B\bar{B}} \\ n_{\bar{B}B} & n_{\bar{B}\bar{B}} \end{pmatrix}, \quad \bar{n} = \begin{pmatrix} n_{\bar{B}\bar{B}} & n_{\bar{B}B} \\ n_{B\bar{B}} & n_{BB} \end{pmatrix}, \quad (18)$$

and n_{eq} is the equilibrium density of B mesons plus anti B mesons. \mathcal{H} is the Hamiltonian for $B\bar{B}$ mixing, see in Eq. (1). $\langle \sigma v \rangle_\pm$ are thermally-averaged annihilation cross section for B meson and anti B mesons, and Γ_\pm are the scattering rates between B mesons and charged particles in the plasma. It turns out the annihilation is negligible for B mesons, while the scattering would be important because

there is a charge radius for B mesons which is different for mesons and antimesons and interacts with the e^\pm particles in the plasma. O_\pm is a matrix

$$O_\pm = \begin{pmatrix} 1 & 0 \\ 0 & \pm 1 \end{pmatrix}. \quad (19)$$

The subscript \pm is determined by the behavior of effective Lagrangian that gives rise to the interaction under charge conjugation of B mesons, $B \leftrightarrow \bar{B}$, $\mathcal{L}_{\text{eff}} \leftrightarrow \pm \mathcal{L}_{\text{eff}}$. Interactions that do not change sign are called flavor-blind interactions while those that change are flavor-sensitive interactions. For the concern of this work, the only interaction that is important to us is the charge radius which gives the neutral B mesons a photon coupling and allows for scattering off of e^\pm particles.

$$\begin{aligned} \Sigma &\equiv n_{BB} + n_{\bar{B}\bar{B}}, & \Delta &\equiv n_{BB} - n_{\bar{B}\bar{B}}, \\ \Xi &\equiv n_{B\bar{B}} - n_{\bar{B}B}, & \Pi &\equiv n_{B\bar{B}} + n_{\bar{B}B}. \end{aligned} \quad (20)$$

After hadronization, b quarks form B_d^0 , B_s^0 and other B mesons, with fragmentation ratio $\text{Br}(\bar{b} \rightarrow B_d^0) = 0.4$ and $\text{Br}(\bar{b} \rightarrow B_s^0) = 0.1$ [27]. Considering the flavor-sensitive interaction only, we can write the Boltzmann equations as:

$$\begin{aligned} \left(\frac{d}{dt} + 3H\right)\Sigma &= \frac{\Gamma_\Phi \rho_\Phi}{m_\Phi} \text{Br}_{\Phi \rightarrow B} - \Gamma_B \Sigma - (\text{Re}\Gamma_{12})\Pi \\ &\quad + i(\text{Im}\Gamma_{12})\Xi, \\ \left(\frac{d}{dt} + 3H\right)\Delta &= -\Gamma_B \Delta + 2i(\text{Re}M_{12})\Xi + 2(\text{Im}M_{12})\Pi, \\ \left(\frac{d}{dt} + 3H\right)\Xi &= -(\Gamma_B + 2\Gamma_{\text{sc}})\Xi + 2i(\text{Re}M_{12})\Delta \\ &\quad - i(\text{Im}\Gamma_{12})\Sigma, \\ \left(\frac{d}{dt} + 3H\right)\Pi &= -(\Gamma_B + 2\Gamma_{\text{sc}})\Pi - 2(\text{Im}M_{12})\Delta \\ &\quad - (\text{Re}\Gamma_{12})\Sigma, \end{aligned} \quad (21)$$

where Γ_B is the decay rate of B mesons, M_{12} , Γ_{12} are the off-diagonal terms of Hamiltonian described in Eq. (1), and Γ_Φ is the decay rate of Φ , which can define the reheating temperature:

$$\Gamma_\Phi = 3H(T_{rh}), \quad (22)$$

where the universe is assumed to be dominated by radiation with temperature T_{rh} .

From these equations we can see coherent oscillations will cause transition $\Pi \rightarrow \Delta$, which will produce B meson asymmetry and thus baryon asymmetry as long as B mesons will decay to baryons. Flavor-sensitive scattering will suppress this process.

The scattering rate could be estimated as [6]:

$$\Gamma_{\text{sc}} \sim 10^{-11} \left(\frac{T}{0.02 \text{ GeV}}\right)^5 \text{ GeV} \quad (23)$$

When T is below 0.01 GeV, the scattering rate is small and the decoherence caused by scattering is no longer significant. The production of the baryon asymmetry mainly takes place below this temperature.

The baryon asymmetry is directly determined by B meson asymmetry:

$$\left(\frac{d}{dt} + 3H\right)\delta_B = \text{Br}_{B \rightarrow \mathcal{B}} \Gamma_B \Delta, \quad (24)$$

where δ_B is the number density of baryon asymmetry and $\text{Br}_{B \rightarrow \mathcal{B}}$ is the branching ratio for $B \rightarrow \mathcal{B} + X$.

It is worth noting that the lifetime and oscillation period of B mesons are much shorter than the Hubble time, so the Hubble term and the time derivative term could be ignored. Under this approximation, we can compute the ratio Δ/ρ_Φ , which is a function of transition amplitudes and scattering rate:

$$\begin{aligned} \frac{\Delta}{\rho_\Phi} &= 2|M_{12}||\Gamma_{12}| \sin(\phi_\Gamma - \phi_M) \Gamma_\Phi (\Gamma_B + 2\Gamma_{\text{sc}}) \\ &\quad \times \{-4|M_{12}|^2 |\Gamma_{12}|^2 \cos(\phi_\Gamma - \phi_M) \\ &\quad + (\Gamma_B^2 + 2\Gamma_B \Gamma_{\text{sc}})(4|M_{12}|^2 - |\Gamma_{12}|^2) \\ &\quad + \Gamma_B^2 (\Gamma_B + 2\Gamma_{\text{sc}})^2\}^{-1}, \end{aligned} \quad (25)$$

where ϕ_Γ , ϕ_M is the phase of Γ_{12} and M_{12} . Under the approximation of $\cos(\phi_\Gamma - \phi_M) \approx 1$, Δ/ρ_Φ is proportional to A_{sl} . Given reheating temperature of the late decaying particle Φ , the number density of net baryons could be calculated by:

$$n_B = \int_{t_0}^{t_1} \text{Br}_{B \rightarrow \mathcal{B}} \frac{\Delta}{\rho_\Phi}(t) \frac{\rho_\Phi(t)}{m_\Phi} \Gamma_B \frac{R(t)^3}{R(t_1)^3} dt, \quad (26)$$

where $R(t)$ is the scale factor of the universe. Δ/ρ_Φ is a function of transition amplitudes and scattering rate, and only scattering rate depends on temperature, which is a function of time. Therefore Δ/ρ_Φ itself is also a function of time. Here t_0 is the time when hadronization begins and t_1 is some time when the universe is dominated by radiation again. The baryon asymmetry would be given by:

$$Y = \frac{n_B}{\frac{2\pi^2}{45} g_{*s}(t_1) T(t_1)^3}, \quad (27)$$

where g_{*s} is the effective degree of freedom at present. The baryon asymmetry can be represented as:

$$Y = \left(\frac{\text{Br}_{B \rightarrow \mathcal{B}}}{10^{-2}}\right) \left(\frac{100 \text{ GeV}}{m_\Phi}\right) (\alpha_d(T) A_{\text{SL}}^d + \alpha_s(T) A_{\text{SL}}^s), \quad (28)$$

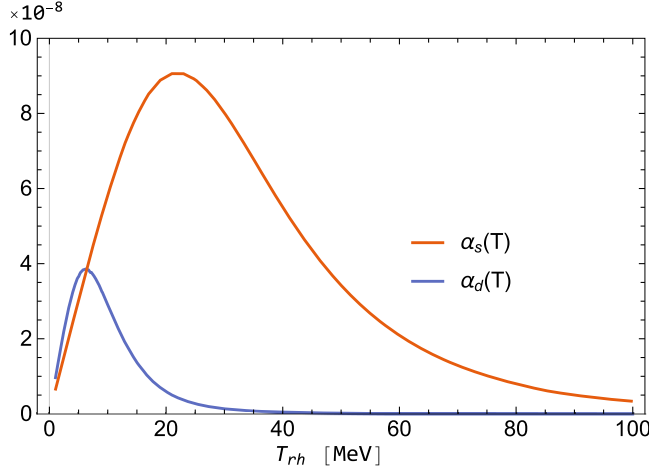


FIG. 3. This figure shows the ability of B_d^0 and B_s^0 to produce baryon number at a given reheating temperature. The scattering rate is greater at higher temperature, which leads to more significant decoherence and suppression of the baryon asymmetry. However, $|M_{12}|$ and $|\Gamma_{12}|$ are very different for B_d^0 and B_s^0 , which leads to a lower characteristic temperature when decoherence is significant for the B_d than the B_s . Thus typically the baryon asymmetry is mostly produced in B_s oscillations.

where $\alpha_d(T)$ and $\alpha_s(T)$ are coefficients as a function of reheating temperature T , whose exact values request numerical study.

The baryon asymmetry, with measurement from cosmic microwave background (CMB) [31,32] and big bang nucleosynthesis (BBN) [27,33], is given by:

$$Y = (8.718 \pm 0.004) \times 10^{-11} \quad (29)$$

From Fig. 3 we can conclude that SM predictions of A_{SL} is too small to produce the expected baryon number. Besides, the SM CP violation predicts that the B_d makes a negative contribution to baryon number. The fact that the B_d contribution is suppressed down to lower temperature means one might hope to produce the correct sign from B_s oscillations, however detailed computation shows that the net effect from the SM CP violation is to give the wrong sign for the baryon asymmetry.

We require A_{SL} to be greater than the SM prediction and positive for the B_s , and not too negative for the B_d , which is a testable feature of our mechanism. However, the exact constraints sensitively depend on the reheating temperature, as shown in Fig. 3. When reheating temperature is at 5 MeV, the A_{SL} should satisfy:

$$\left(\frac{\text{Br}_{B \rightarrow B}}{10^{-2}}\right) \left(\frac{100 \text{ GeV}}{m_\phi}\right) (0.42A_{\text{SL}}^d + 0.35A_{\text{SL}}^s) \approx 10^{-3}. \quad (30)$$

When reheating temperature is higher, the contribution from B_d would be suppressed. For reheating temperature at 25 MeV, the A_{SL} should satisfy:

$$\left(\frac{\text{Br}_{B \rightarrow B}}{10^{-2}}\right) \left(\frac{100 \text{ GeV}}{m_\phi}\right) (0.03A_{\text{SL}}^d + 1.02A_{\text{SL}}^s) \approx 10^{-3}. \quad (31)$$

Given that A_{SL} is always negative for the B_d , a high reheating temperature around 25 MeV is needed in order to suppress the CP violation in the B_d oscillations. Combining with the range of allowed A_{SL} in Eq. (10), we find that a branching ratio $\text{Br}_{B \rightarrow B}$ larger than about 10^{-3} is required for baryon violating b quark decays. Also, to produce the expected baryon number, we expect A_{SL}^s to be positive and greater than the SM predictions.

V. DISCUSSION

In this work, we have shown that B meson oscillations could solve the puzzle of baryogenesis within a simple renormalizable model containing three new particles. This model predicts an exotic baryon number violating B meson decay mode. We also predict new contributions to the semileptonic asymmetry in B meson oscillations. Specifically, we link the sign of the matter asymmetry to a new positive contribution to the semileptonic symmetry in B_s meson decays. Another prediction is that immediately prior to nucleosynthesis in the early universe, the energy density is dominated by a massive late decaying particle. This may have implications for forming small clumps of dark matter in the early universe.

The baryon asymmetry is proportional to the branching fraction for exotic B meson decays into a baryon and missing energy. These features could be searched at LHCb and Belle-II, as discussed in [6]. The missing energy is carried by a long-lived Majorana fermion χ , which will decay into a baryon or antibaryon, and could be found in dedicated searches for long-lived particles [10,11].

The cosmological constraints for the corresponding reheating temperature is $T_{\text{RH}} > 4.7$ MeV [34]. The thermal history of early universe before BBN is hard to probe. A late decaying ϕ particle does have some implications for axion miniclusters [7] and substructure formation [35]. It happens that the reheating temperature that is favored by axion minicluster also tends to generate baryon number efficiently with our mechanism. Looking for axion miniclusters may also provide evidence for early matter domination.

Embedding this model into a R-parity violating supersymmetric model would provide more motivation and phenomenological implications.

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APPENDIX: NEW CONTRIBUTIONS TO NEUTRAL B MESON OSCILLATIONS

In this appendix we will give explicit formulas for the extra contributions to M_{12} from new physics, whose Feynman diagrams are already shown in Fig. 2. They are very similar diagrams with those from Standard Model except the gamma matrices. Therefore we are doing similar computations as in Refs. [20–22].

The first box diagram will give us the following effective Hamiltonian:

$$H_{\text{eff}}^{(1)} = \bar{b}_v(1-\gamma^5)\gamma^\alpha d_u \bar{d}_v(1-\gamma^5)\gamma^\beta b_u p_\alpha(p-q)_\beta I_{\alpha\beta}^{(1)} + \text{H.c.} \quad (\text{A1})$$

where p is the momentum of b quark and q is the loop momentum. The integral $I_{\alpha\beta}$ is given as:

$$I_{\alpha\beta}^{(1)} = \frac{1}{4}(y_b y_d^*)^2 \int \frac{d^4 q}{(2\pi)^4} q_\alpha(p-q)_\beta \frac{1}{(q^2 + m_\phi^2)} \times \frac{1}{[(p-q)^2 + m_\phi^2]} \frac{1}{(q^2 + m^2)[(p-q)^2 + m^2]}, \quad (\text{A2})$$

where m denotes the mass of Majorana fermion. We have ignored the momentum distributed to u, s quarks because they can be approximated as massless particles. The second Feynman diagram will give us similar results. However, the contribution from the third diagram is negligible. After manipulating those gamma matrices, we will find that it is proportional to m^2 , which is very small compared to m_ϕ^2 . The effective Hamiltonian from the second diagram could be written as:

$$H_{\text{eff}}^{(2)} = \bar{b}_v(1-\gamma^5)\gamma^\alpha d_v \bar{b}_u(1-\gamma^5)\gamma^\beta d_u(p-q)_\alpha q_\beta I_{\alpha\beta}^{(2)} + \text{H.c.} \quad (\text{A3})$$

Switching indices on $I_{\alpha\beta}^{(1)}$ will give us $I_{\alpha\beta}^{(2)}$:

$$I_{\alpha\beta}^{(2)} = I_{\beta\alpha}^{(1)}. \quad (\text{A4})$$

There is a trick on calculating this Feynman integral, which utilize the symmetry of the Feynman diagrams.

$$I_{\alpha\beta}^{(1)} = q_\alpha(p-q)_\beta \frac{1}{(m_\phi^2 - m^2)} \int \frac{d^4 q}{(2\pi)^4} \times \left\{ \frac{1}{(q^2 + m^2)[(p-q)^2 + m^2]} + \frac{1}{(q^2 + m_\phi^2)[(p-q)^2 + m_\phi^2]} - \frac{1}{(q^2 + m_\phi^2)[(p-q)^2 + m^2]} - \frac{1}{(q^2 + m^2)[(p-q)^2 + m_\phi^2]} \right\}. \quad (\text{A5})$$

Applying Feynman's formula:

$$\frac{1}{(q^2 + m_\phi^2)[(p-q)^2 + m_\phi^2]} = \int_0^1 dx \{ (q-xp)^2 + x(1-x)p^2 + xm_\phi^2 + (1-x)m_\phi^2 \}^{-2}. \quad (\text{A6})$$

Since the momentum is given by on-shell relation: $p^2 = -m_b^2$, it is convenient to denote:

$$\begin{aligned} D_1 &= -x(1-x)m_b^2 + m_\phi^2 \\ D_2 &= -x(1-x)m_b^2 + m^2 \\ D_3 &= -x(1-x)m_b^2 + xm^2 + (1-x)m_\phi^2 \\ D_4 &= -x(1-x)m_b^2 + xm_\phi^2 + (1-x)m^2 \end{aligned} \quad (\text{A7})$$

The Hamiltonian could be represented as:

$$\mathcal{H}_{\text{eff}} = A \bar{d}_\alpha \gamma^\mu (1 + \gamma^5) b_\alpha \bar{d}_\beta \gamma_\mu (1 + \gamma^5) b_\beta + B \bar{d}_\alpha (1 - \gamma^5) b_\alpha \bar{d}_\beta (1 - \gamma^5) d_\beta + \text{H.c.} \quad (\text{A8})$$

Coefficients A, B can be written as:

$$\begin{aligned} A &= \frac{-1}{16\pi^2(m_\phi^2 - m^2)^2} \int_0^1 dx \sum_{i=1}^4 {}' D_i \ln \frac{D_i}{m_\phi^2} \\ B &= \frac{m_b^2}{32\pi^2(m_\phi^2 - m^2)^2} \int_0^1 dx \sum_{i=1}^4 {}' \ln D_i, \end{aligned} \quad (\text{A9})$$

where $\sum_{i=1}^4 {}' = \sum_{i=1}^2 - \sum_{i=3}^4$. Since D_i are always positive for the whole range of x , there is no contribution to Γ_{12} . The contribution to M_{12} is

$$\Delta M_{12} = \langle B^0 | \mathcal{H}_{\text{eff}} | \bar{B}^0 \rangle. \quad (\text{A10})$$

The matrix elements are corresponding to the nonperturbative effects, which could be estimated as:

$$\begin{aligned} \langle B^0 | \bar{d}_\alpha \gamma^\mu (1 + \gamma^5) b_\alpha \bar{d}_\beta \gamma_\mu (1 + \gamma^5) b_\beta | \bar{B}^0 \rangle &= \frac{8}{6} f_B^2 m_B B_B^2 \\ \langle B^0 | \bar{d}_\alpha (1 - \gamma^5) b_\alpha \bar{d}_\beta (1 - \gamma^5) d_\beta | \bar{B}^0 \rangle &= -\frac{5}{6} f_B^2 m_B B_B^2, \end{aligned} \quad (\text{A11})$$

where f_B is the decay constant of B mesons, m_B is B meson mass and B_B is a bag parameter, which is order one. For $m = 2$ GeV, which is the parameter space we are interested in, the result is

$$\Delta M_{12} = 0.66 \frac{f_B^2 m_B B_B^2}{16\pi^2 m_\phi^2} (y_b y_d^*)^2. \quad (\text{A12})$$

Lattice study can numerically give us the value of decay constant and bag parameter, $B_B = 0.87$, $B_{B_s} = 0.9$, $f_B = 0.192$ GeV, $f_{B_s} = 0.228$ GeV. [36,37].

The ratio to the experimental value of M_{12} is

$$\begin{aligned}\frac{\Delta M_{12}}{M_{12}} &= 3689 \left(\frac{1 \text{ TeV}}{m_\phi} \right)^2 (y_b y_d^*)^2, \\ \frac{\Delta M_{12s}}{M_{12s}} &= 161 \left(\frac{1 \text{ TeV}}{m_\phi} \right)^2 (y_b y_s^*)^2.\end{aligned}\quad (\text{A13})$$

Therefore y_b, y_d, y_s must be much smaller than 1 to agree with experimental results. This will not cause a problem because ΔM_{12} is the only observable that is only determined by y_d . It is worth noting that the calculations we present here are very much the same with box diagram calculations in a supersymmetric theory [38], but with very different parameters.

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