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Baryogenesis from Oscillations of Charmed or Beautiful Baryons

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We propose a model for CP violating oscillations of neutral, heavy-flavored baryons into antibaryons at rates which are within a few orders of magnitude of their lifetimes. The flavor structure of the baryon violation suppresses neutron oscillations and baryon number violating nuclear decays to experimentally allowed rates. We also propose a scenario for producing such baryons in the early Universe via the out-of-equilibrium decays of a neutral particle, after hadronization but before nucleosynthesis. We find parameters where CP violating baryon oscillations at a temperature of a few MeV could result in the observed asymmetry between baryons and antibaryons. Furthermore, part of the relevant parameter space for baryogenesis is potentially testable at Belle II via decays of heavy flavor baryons into an exotic neutral fermion. The model introduces four new particles: three light Majorana fermions and a colored scalar. The lightest of these fermions is typically long lived on collider timescales and may be produced in decays of bottom and possibly charmed hadrons.

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I. INTRODUCTION

The puzzle of *baryogenesis*, how the Universe came to be composed primarily of matter rather than equal amounts of matter and antimatter, has led to numerous theories about physics beyond the standard model (SM), beginning with the pioneering work of Sakharov [1]. Three ingredients are present in one form or another in any baryogenesis theory: baryon number violation, C and CP violation, and departure from thermal equilibrium. Because baryon number violation is required, initially baryogenesis was thought to involve new baryon number violating processes which are only important at very high energies, although it was later realized that anomalous electroweak processes could do the job at temperatures as low as the weak phase transition [2].

Most baryogenesis models require the Universe to reheat after inflation to a high temperature, typically well above the weak scale. However, many theories of physics beyond the SM are inconsistent with a high inflation scale or are inconsistent with a high postinflation reheat scale. Axion dark matter, if the axion is present during inflation, requires a low inflation scale in order to avoid excessive isocurvature perturbations [3]. Supersymmetry requires a low reheat scale in order to avoid overproduction of the gravitino [4]. The relaxion solution to the hierarchy problem requires a low inflation scale so that the Hubble temperature during inflation does not suppress instantons [5]. In addition, avoiding the need for a high reheat temperature or production of heavy particles during reheating means a low baryogenesis scale is consistent with a wider variety of inflationary models [6].

Lower reheat temperatures are possible provided the inflaton decays produce heavy particles which decay out of thermal equilibrium in a baryon and CP violating manner [7]. In Ref. [7] baryogenesis occurs due to the baryon number violating decays of TeV mass squarks in an R-parity violating supersymmetric model, in which the reheat temperature could be as low as an MeV, provided that the heavy squarks can be produced out of equilibrium at the end of inflation. Such squark mediated baryon number violation is consistent with the observed lifetime of the proton, due to the conservation of lepton number, and, depending on the flavor structure of the baryon number violating operators, can be consistent with the stability of heavy nuclei as well. A similar model involving the decays of Majorana fermions was considered in Ref. [8]. In Ref. [9] it was pointed out that heavy flavor baryon number violation could lead to oscillations of the Ξ_b^0 baryon at a rate comparable to its lifetime, while being consistent with the lifetime of heavy nuclei. Enhanced baryon number violation involving heavy flavors was also studied in Ref. [10].

Here we present a baryogenesis model which is consistent with a reheat temperature as low as a few MeV, and which requires no postinflationary production of any particle heavier than about 6—10 GeV. The required baryon number violation is conceivably observable via the oscillations of heavy-flavor neutral baryons, and the required CP violation is potentially of $\mathcal{O}(1)$ in such oscillations. The processes that produce the baryon asymmetry in the early Universe involve particles and phenomena which can be directly studied in the laboratory – a unique feature of our theory. Our proposal is that certain neutral

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heavy flavor baryons undergo CP and baryon number violating oscillations and decays, and are produced in the early Universe via the out of equilibrium decays of a weakly coupled neutral particle whose lifetime is of order 0.1 s, a time when the temperature is of order a few MeV. The basic scenario was outlined in Ref. [11], and the model we study has the same field content and couplings as Ref. [12]. The basic formalism for analyzing CP violation in fermion antifermion oscillations was worked out in Ref. [13].

The outline of the paper is as follows. In section II the model is introduced, and the effective operator responsible for baryon oscillations is constructed. In section III, general $\Delta B = 2$, six-quark effective operators are analyzed for their contribution to dinucleon decay, that is, the decay of two nucleons into mesons. Currently, dinucleon decay places similar or stronger constraints on all such operators than does neutron oscillations. For operators that cannot contribute to dinucleon decay at tree level, electroweak corrections to the six-quark operators are examined. In section IV, the general formalism for CP-violating oscillations of fermions is reviewed, and the oscillation parameters are calculated for the model introduced in section II. In sections V and VI, direct constraints on the masses and couplings of the new ϕ and χ particles from collider searches, and indirect constraints from rare decays of mesons and baryons are derived, respectively. Section VII contains our analysis of how in this model the baryon asymmetry of the Universe (BAU) is produced. Finally, in section VIII we conclude and point at possible directions for future work.

II. MODEL

We wish to find a theory which allows for sufficiently large baryon number and CP violation to explain baryogenesis at relatively low energy. In order to ensure sufficient stability of the proton, we assume lepton number is not violated, other than perhaps via the tiny $\Delta L = 2$ terms that could account for Majorana neutrino masses. The lowest dimension terms which violate baryon number and not lepton number are dimension 9, six-quark $\Delta B = 2$ operators. Such operators can lead to neutral baryon oscillations and conceivably CP violation [11], and can arise as an effective field theory description of physics at some higher energy scale. A minimal renormalizable model for generating such terms involves a new charge $-1/3$ color triplet scalar and two Majorana fermions, as described in Ref. [12]. A third Majorana fermion, which decays out of thermal equilibrium, allows for the fulfillment of the out-of-equilibrium Sakharov condition.

We note that this model for baryon number violation can easily be embedded in an R-parity violating supersymmetric (RPV SUSY) theory. In such theories, the neutralinos would play the role of the Majorana fermions and a down-type SU(2) singlet squark can be the colored scalar. For simplicity, we do not explore this embedding in a SUSY framework in this paper and we stick to the minimal version of the model.

Our model thus adds four new particles: three Majorana fermions, $\chi_{1,2,3}$, and a single color triplet scalar, ϕ. The interactions involving the new particles and weak SU(2) singlet SM quarks are given by

$$\mathcal{L}_{\text{int}} \supset -g^*_{ud}\phi^*\bar{u}_R d^c_R - y_{id}\phi\bar{\chi}_i d^c_R + \text{h.c.}, \tag{1}$$

along with terms involving other generations, $d \to s, b$ and $u \to c, t$. By convention we take all two component fields to transform in the left-handed representation under Lorentz transformations. d^c_R stands for the charge conjugate of the right-handed down quark field, which is in the left-handed Lorentz representation. In this expression and throughout the paper, we work in the mass basis, which is unambiguous for SU(2) singlet quarks.

The required new particles and their interactions are motivated as follows. A natural way to construct the $\Delta B = 2$ six-quark operator we require for baryon oscillations is from two $\Delta B = 1$ four-fermion interactions connected by an exotic neutral Majorana fermion. Thus we introduce an exotic, electrically-neutral, colorless, Majorana fermion, χ_1, which couples to other fermions via a four-fermion interaction of the form $u_R d_R d'_R \chi_1$ (using u and d here to represent any up- or down-type quark).[1]

Since such a four-fermion interaction is itself nonrenormalizable, we also introduce a complex, color triplet, scalar particle (diquark) ϕ to mediate the $\Delta B = 1$ interactions. Note that if χ_1 is heavier than the difference in mass between the proton and electron, $m_p - m_e = 937.76$ MeV, this interaction does not give rise to proton decay.[2] In the presence of only χ_1, there is no physical CP violation, as there is enough reparameterization freedom to remove the phases in the couplings. We introduce a second fermion, χ_2 (with $m_{\chi_2} > m_{\chi_1}$), in order to give rise to CP violation. Finally, for baryogenesis, the oscillating baryons must be produced out of thermal equilibrium. As described in Sec. VII, this is most simply accomplished by introducing a third Majorana fermion, χ_3, which decays out of equilibrium to produce the baryons whose oscillations result in baryogenesis.

Note that we only consider operators constructed out of right-handed quarks, for two reasons. Our phenomenological reason is that, as we will show in Sec. III, right handed quark operators are less constrained by dinucleon decay due to the requirement of light quark mass insertions in flavor-changing loops. Our top down theoretical reason is that, as mentioned, the interactions in Eq. (1) occur in RPV SUSY models, suggesting a possible embedding of our model into a more complete theory.

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[1] Models of baryogenesis that generate four-fermion interactions of this form and identify χ_i with right-handed neutrinos can be found in Ref. [14].

[2] The stability of ^{9}Be leads to a marginally stronger lower bound of $m_{\chi_1} > 937.9$ MeV [11].

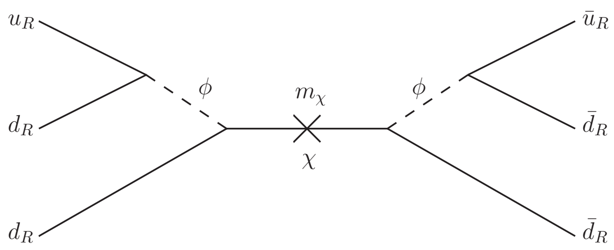

FIG. 1. The basic six-quark $\Delta B = 2$ operator generated by ϕ and χ exchange. u and d here represent any of the up- or down-type quark flavors. The quarks involved are all weak SU(2) singlets, as emphasized by the R subscripts.

Figure 1 shows the $\Delta B = 2$ six-quark operators that are generated by the interactions in Eq. (1). Such operators, which can mediate the transition of a baryon $\mathcal{B}$ to an antibaryon $\bar{\mathcal{B}}'$, can be written as

$$\mathcal{O}_{\mathcal{B}\mathcal{B}'} = \epsilon^{abc}\epsilon^{def} \times \left[(q_R)^i_a \, (q_R)_{i,b} \, (q_R)^j_c \, (q'_R)_{j,d} \, (q'_R)^k_e \, (q'_R)_{k,f} + \dots \right], \tag{2}$$

where $a, \dots, f$ are color indices, i, j, k spinor indices, and $q = u, d, s, c, b$ any of the quark flavors (because of its short lifetime, the top quark does not hadronize and is not important in the low-energy effective theory) which are all right-chiral. The ellipsis represents other possible permutations of color or spinor indices. Here, $\mathcal{B}$ denotes an arbitrary standard model baryon with the quark content qqq while $\mathcal{B}'$ contains $q'q'q'$. [We will use both the baryon name $\mathcal{B}$ or the quark content (qqq) to label the operators in question throughout this paper.] The precise index structure of $\mathcal{O}_{\mathcal{B}\mathcal{B}'}$ is not important for the purposes of this paper. Therefore, in what follows, we will suppress the indices on $\mathcal{O}_{\mathcal{B}\mathcal{B}'}$ and generically denote the operators we are interested in that appear in the effective Lagrangian via the shorthand

$$\mathcal{L}_{\text{eff}} \supset C_{\mathcal{B}\mathcal{B}'} \, (qqq) \, (q'q'q') \equiv C_{\mathcal{B}\mathcal{B}'} \mathcal{O}_{\mathcal{B}\mathcal{B}'}, \tag{3}$$

keeping in mind that the leading operators that are generated involve only right-chiral quarks.

Matching the interactions generated by Eq. (1) to the effective theory at tree-level gives the coefficient of the operator that generates oscillations between a neutral baryon and its antiparticle, $\mathcal{B} \leftrightarrow \bar{\mathcal{B}}$,

$$C_{\mathcal{B}\mathcal{B}} \sim \sum_i \frac{m_{\chi_i}}{m_{\mathcal{B}}^2 - m_{\chi_i}^2} \left(\frac{g^*_{ud} y_{id'} + g^*_{ud'} y_{id}}{m_\phi^2} \right)^2, \tag{4}$$

with u, d, and d' labeling the quarks comprising $\mathcal{B}$. For example, the operator $(ddc)^2$ would allow the processes $\bar{\Sigma}_c \leftrightarrow \Sigma_c$. Given this operator, we will find it useful to relate the coefficient to the (dispersive) transition amplitude, defined by $\delta_{\mathcal{B}\mathcal{B}} \equiv \langle \bar{\mathcal{B}} | C_{\mathcal{B}\mathcal{B}} \mathcal{O}_{\mathcal{B}\mathcal{B}} | \mathcal{B} \rangle$, with

$$\delta_{\mathcal{B}\mathcal{B}} = \kappa^2 C_{\mathcal{B}\mathcal{B}}, \tag{5}$$

where $\kappa \sim 10^{-2}$ GeV^3 [15]. In analogy with meson oscillations, when the two-state system in question is unambiguous, $\delta_{\mathcal{B}\mathcal{B}}$ can also be referred to as M_{12}.

Operators which involve different baryons of the form $\mathcal{O}_{\mathcal{B}\mathcal{B}'}$ would allow for a common decay product between $\mathcal{B}$ and $\bar{\mathcal{B}}$ and could also give rise to oscillations. For example, $(uss)\,(uds)$ would allow for Ξ^0 and $\bar{\Xi}^0$ to have a common decay product, through $\Xi^0 \to \Lambda^0$ and $\bar{\Xi}^0 \to \Lambda^0$ (ignoring any neutral meson products). However, we will find such processes are suppressed relative to their direct oscillation cousins, so we will ignore them in our analysis. There are also baryon-number–preserving operators that contribute to the masses and mixings of SM baryons. These are greatly suppressed relative to those that occur in the SM and we do not consider them further either. Therefore, in what follows, we focus on operators $\mathcal{O}_{\mathcal{B}\mathcal{B}}$ with coefficients of the form of Eq. (4).

III. DINUCLEON DECAY CONSTRAINTS

As described in the preceding section, we would like our six-quark operators to allow for the oscillation of heavy baryons in order to produce the Universe's observed baryon asymmetry. In Sec. VII, we will show that the ideal width for such an oscillation, which is dependent upon the value of $C_{\mathcal{B}\mathcal{B}}$ in Eq. (4), is a few orders of magnitude smaller than $\mathcal{B}$'s decay width. However, models with B violation by two units are certainly not a new idea, and so significant experimental effort has been put fourth into constraining $\Delta B = 2$ processes. The most immediate constraint on our six-quark operators is the lack of observed dinucleon decay, which we quantify in this section. The analysis we perform here applies to six-quark operators in general and is independent of the origin of the new physics introduced in Sec. II.

Dinucleon constraints come from underground detectors whose primary purpose is the detection of proton decay and neutrino oscillations. For example, in a nucleus, a $n \to \bar{n}$ transition will be shortly followed by the annihilation of the $\bar{n}$ with one of the other nucleons, leading to the decay of the nucleus of mass number A to a nucleus with $A' = A - 2$ plus mesons. The lack of observation of such decays can therefore bound the transition amplitude δ_{nn} [16] which is related to the coefficient of the $(udd)^2$ operator, C_{nn} [cf. Eqs. (4) and (5)]. Currently, the lower bound on the ^{16}O lifetime (in decays to pions) of 1.9×10^{32} years from the Super-Kamiokande collaboration [17] places the strongest limit, $\delta_{nn} < 1.9 \times 10^{-33}$ GeV.

Operators that also violate strangeness do not directly induce $n \to \bar{n}$ transitions in a nucleus. However, they can also lead to dinucleon decays, $A \to (A-2) + \text{mesons}$, through the reaction $NN \to \text{kaons} + X$ where N is a nucleon. For example, the diagram on the left of Fig. 2 shows how the operator $(uds)^2$ can lead to dinucleon decay to a pair of kaons. The Super-Kamiokande collaboration [18] has searched for such decays and has placed

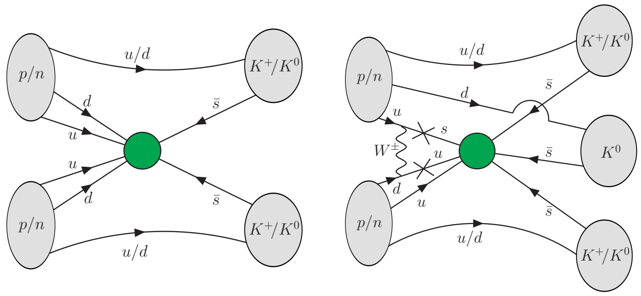

FIG. 2. Left: Dinucleon decay via the $\Delta B = \Delta S = 2$ $(uds)^2$ operator that mediates $\Lambda^0 \leftrightarrow \bar{\Lambda}^0$ oscillations. Right: Dinucleon decay mediated by the $\Delta B = 2$, $\Delta S = 4$ $(uss)^2$ operator that becomes $\Delta S = 3$ $(uds)(uss)$ operator in the presence of flavor-changing weak interactions. Because the short-distance $\Delta B = 2$ operators we consider involve weak isosinglets, this operator requires light quark chirality flips, indicated by crosses. See text for discussion of the matching of the short distance theory onto the (chiral symmetry violating) long distance theory.

an upper bound on the $pp \to K^+K^+$ decay rate by limiting the lifetime for ${}^{16}O \to {}^{14}C\, K^+K^+$ to more than 1.7×10^{32} years.

To make use of this limit, we start with the effective operator $\mathcal{O}_{\mathcal{BB}}$. The dinucleon decay rate through direct nucleon annihilation can then be roughly approximated by considering the decay rate to a meson pair [19],

$$\Gamma_{NN\to \mathrm{X}} \sim \frac{9}{32\pi}\frac{|C_{\mathcal{BB}}|^2}{m_N^2}\,|\langle 2\text{ mesons}|\mathcal{O}_{\mathcal{BB}}|NN\rangle|^2\,\rho_N \quad (6)$$

where m_N is the nucleon mass, $\rho_N \simeq 0.25$ fm^{-3} is the nucleon density, and we have ignored the masses of the final state particles. In the case of operators that can contribute at tree level, the matrix element can be estimated as roughly $\langle 2\text{ mesons}|\mathcal{O}_{\mathcal{BB}}|NN\rangle \sim \Lambda^5_{\mathrm{QCD}} \simeq (200\text{ MeV})^5$. Using this and Eq. (5), the limit on the rate for $NN \to KK$ from Super Kamiokande translates to a limit on the transition amplitude of

$$\delta_{(uds)^2} \lesssim 10^{-30}\text{ GeV}. \quad (7)$$

In what follows, we also take operators that change strangeness by one or three units to have roughly the same bound as this.

Kinematic constraints protect certain operators from contributing to dinucleon decay at leading order. Operators such as $(uss)^2$ that change strangeness by four units (i.e. $\Delta S = 4$) are kinematically forbidden from contributing to dinucleon decay at tree level since $2m_N < 4m_K$. Similarly, those that involve charm[3] or bottom quarks also do not lead to dinucleon decay at leading order. However, when combined with flavor-violating weak interactions, these operators involving heavy quarks can lead to dinucleon decay. An illustration of this is shown on the right of Fig. 2.

To properly estimate the rate for dinucleon decay from $\Delta B = 2$ operators (involving heavy flavors), we must match the UV theory involving quarks to a low energy effective theory involving baryons valid at momentum transfers below $4\pi f_\pi \sim 1$ GeV where $f_\pi = 93$ MeV. This consists of writing down an operator in the UV theory and treating the coefficient of this operator as a spurion that transforms in a particular way under the global chiral quark flavor symmetry $\mathrm{SU}(3)_L \times \mathrm{SU}(3)_R$ so as to make the operator invariant. This operator is then matched onto an operator in the effective theory that transforms in the same way under the chiral symmetry with the same spurion coefficient. In the UV, the light quarks $q_{L,R}$ transform as triplets under $\mathrm{SU}(3)_{L,R}$. In the low energy theory, the meson octet, Π, is described by a field $\Sigma = \exp(2i\Pi/f_\pi)$ which transforms under the chiral symmetry as $\Sigma \to L\Sigma R^\dagger$ where L, R are $\mathrm{SU}(3)_{L,R}$ transformations, respectively. Incorporating the baryon octet (see, e.g., Ref. [20]) can be done by defining a field $\xi = \exp(i\Pi/f_\pi)$ which transforms as $\xi \to L\xi U^\dagger$, $\xi \to U\xi R^\dagger$ under $\mathrm{SU}(3)_{L,R}$. U is an SU(3) matrix that depends nonlinearly on the meson fields. The baryon octet B is defined to transform as $B \to UBU^\dagger$. Operators in the effective theory are then constructed out of Σ, B, and ξ along with spurions from the UV theory to be invariant under the flavor symmetry. Since the chiral symmetry is dynamically broken by the strong coupling of QCD around $4\pi f_\pi$, one can use naive dimensional analysis to properly account for factors of 4π (that come from the strong coupling) and the cutoff, $4\pi f_\pi$, that appear in this matching procedure, as described in, e.g., Ref. [21].

We will first illustrate this matching procedure in our theory with interactions given by Eq. (1), assuming for now that only the light quarks u, d, and s are involved. We will deal with heavy quarks c and b below. Since distinction between chiralities is necessary, we will temporarily denote them explicitly. After integrating out the scalar, ϕ, and the Majorana fermions, χ_i, we are left with a $\Delta B = 2$ operator involving only (light) right-chiral quarks, $C_{\mathcal{BB}}(q_Rq_Rq_R)^2$. This operator must be matched onto an operator valid at long distances involving baryons at the scale of chiral symmetry breaking. The coefficient $C_{\mathcal{BB}}$ can be treated as a spurion that transforms under $\mathrm{SU}(3)_R$ in a representation that appears in the tensor decomposition of 6 triplets. For definiteness, take it to transform as an $\mathrm{SU}(3)_R$ octet. Then the object $\tilde{C}_{\mathcal{BB}} \equiv \xi C_{\mathcal{BB}}\xi^\dagger$ transforms as $\tilde{C}_{\mathcal{BB}} \to U\tilde{C}_{\mathcal{BB}}U^\dagger$ and the

[3] Depending on the nucleon binding energy, $nn \to D\gamma$ through a $\Delta C = 1$ operator is kinematically allowed for some nuclei, but due to the dependence of the amplitude on the photon momentum and coupling and phase space suppression the rate is proportional to $(\alpha/4\pi)\,(k_\gamma/m_N)^3 \sim 10^{-9}$, where k_γ is the photon energy, suppressing the rate below other decays with less constrained phase space.

operator matching is

$$C_{\mathcal{B}\mathcal{B}}(q_R q_R q_R)^2 \to \left(4\pi f_\pi^3\right)^2 \operatorname{tr} B \tilde{C}_{\mathcal{B}\mathcal{B}} B + \ldots, \tag{8}$$

where the ellipsis represents other possible orderings of the baryon octets and the spurion. Note that this gives an understanding of the size of $\kappa \sim 4\pi f_\pi^3 \simeq 10^{-2}$ GeV^3 in Eq. (5). Adjusting this analysis if $C_{\mathcal{B}\mathcal{B}}$ transforms under a different representation of $\mathrm{SU}(3)_R$ is straightforward; one inserts the required numbers of ξ and $\xi^\dagger$ into the definition of $\tilde{C}_{\mathcal{B}\mathcal{B}}$ so that it transforms in such a way as to leave $\operatorname{tr} B \tilde{C}_{\mathcal{B}\mathcal{B}} B$ invariant. For example, if $C_{\mathcal{B}\mathcal{B}}$ is a singlet then one simply takes $\tilde{C}_{\mathcal{B}\mathcal{B}} \equiv C_{\mathcal{B}\mathcal{B}}$.

Four-quark weak operators involving light quarks can be matched onto the low energy effective theory in much the same way. The coefficient of the operator $\bar{u}_L \gamma^\mu q_L^j \bar{q}_{Li} \gamma_\mu u_L$ can be viewed as a spurion that transforms as an octet under $\mathrm{SU}(3)_L$ and the strangeness changing ($\Delta S = 1$) coefficient takes a value $\propto G_F V_{us} V_{ud}^* h$ with $h^i_j = \delta^i_2 \delta^3_j$. Then $\xi^\dagger h \xi \to U \xi^\dagger h \xi U^\dagger$ and the matching is

$$\begin{aligned} &\frac{G_F}{\sqrt{2}} V_{us} V_{ud}^* \bar{u} \gamma^\mu \left(1-\gamma^5\right) s \bar{d} \gamma_\mu \left(1-\gamma^5\right) u \\ &\quad \to \frac{G_F}{\sqrt{2}} V_{us} V_{ud}^* \left(4\pi f_\pi^3\right) \operatorname{tr} \bar{B} \xi^\dagger h \xi B + \ldots, \end{aligned} \tag{9}$$

where again the ellipsis represents other possible orderings of B, $\bar{B}$, and $\xi^\dagger h \xi$.

Now we can combine a $\Delta B = 2$ operator that also changes strangeness by n units with the weak $\Delta S = 1$ operator to form a $\Delta B = 2$, $\Delta S = n-1$ operator that is given by

$$\frac{G_F}{\sqrt{2}} V_{us} V_{ud}^* f_\pi^2 \left(4\pi f_\pi^3\right)^2 \operatorname{tr} \bar{B} \tilde{C}_{\mathcal{B}\mathcal{B}} \xi^\dagger h \xi B + \ldots. \tag{10}$$

In other words, if the leading $\Delta B = 2$ operator has $\Delta S = n$, the $\Delta B = 2$, $\Delta S = n-1$ operator that is generated due to weak interactions is suppressed relative to it by the factor

$$\frac{G_F}{\sqrt{2}} V_{us} V_{ud}^* f_\pi^2 \sim 10^{-8}. \tag{11}$$

Thus, for example, the bound on a leading $\Delta S = 4$ operator $(u_R s_R s_R)^2$ from the lack of dinucleon decay is around eight orders of magnitude weaker than that on the $\Delta S = 3$ operator $(uds)(uss)$ [which we take to be comparable to that on the $\Delta S = 2$ operator $(uds)^2$],

$$\delta_{(uss)^2} \lesssim 10^{-22} \text{ GeV}. \tag{12}$$

Now, we consider the case where the leading $\Delta B = 2$ operators contain heavy quarks. Consider, for example, if after integrating out the heavy scalar ϕ and Majorana fermions χ_i, that the leading operator we generate is $C_{(udb)^2}(u_R d_R b_R)^2$. Before matching onto the theory valid after chiral symmetry breaking we must first integrate out the b quarks. In the presence of weak interactions, as shown in Fig. 3, doing so will lead to a ten-quark operator,

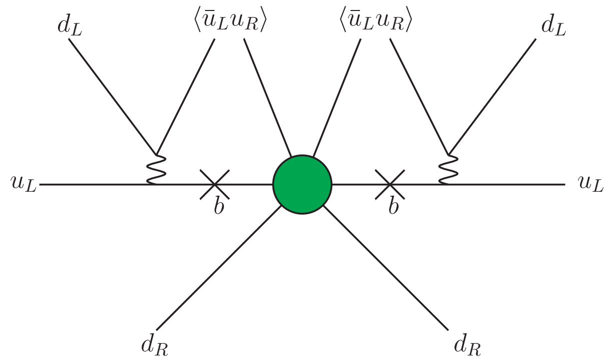

FIG. 3. The ten-quark $\Delta B = 2$ operator that results from the leading $(u_R d_R b_R)^2$ operator after integrating out the b quarks. The crosses represent chirality-flipping b quark mass insertions. We use $\langle \bar{u}_L u_R \rangle$ to indicate the pairs of light quark fields that can be replaced by the chiral condensate when matching onto the long distance theory relevant for dinucleon decay.

$$\begin{aligned} C_{(udb)^2}(u_R d_R b_R)^2 \to &\left(\frac{G_F}{\sqrt{2}} V_{ub} V_{ud}^* \frac{1}{m_b}\right)^2 \\ &\times C_{(udb)^2}(u_R d_R d_L \bar{u}_L u_L)^2. \end{aligned} \tag{13}$$

After chiral symmetry breaking, $\bar{u}_L u_R$ can be replaced by the quark condensate which is roughly $4\pi f_\pi^3$. This means that the induced $\Delta B = 2$ operator $(u_L d_R d_L)^2$ is suppressed relative to $(u_R d_R b_R)^2$ by the factor

$$\left(\frac{G_F}{\sqrt{2}} V_{ub} V_{ud}^* \frac{4\pi f_\pi^3}{m_b}\right)^2 \sim 10^{-20}. \tag{14}$$

Additionally in the case of a leading operator containing a b and c quark, e.g. $(d_R c_R b_R)^2$, there are perturbative loops that generate dimension-nine operators involving light quarks above the chiral symmetry breaking scale. In the case of the operator $(d_R c_R b_R)^2$, two such loops can be used to generate the operator $(u_L d_L d_R)^2$ with a coefficient suppressed relative to the leading one by

$$\left(\frac{G_F}{\sqrt{2}} V_{ub} V_{cd}^* \frac{m_b m_c}{4\pi^2} \log \frac{m_W^2}{m_b^2}\right)^2 \sim 10^{-16}. \tag{15}$$

In Table I, we list operators that can mediate $\mathcal{B} \leftrightarrow \bar{\mathcal{B}}$ transitions along with the number of loops required for each operator to mediate $(\Delta S = 0, 1, 2, 3)$ dinucleon decay. We show the resulting limits on the transition amplitudes $\delta_{\mathcal{B}\mathcal{B}} = |M_{12}| = \kappa^2 C_{\mathcal{B}\mathcal{B}}$ of each operator from the lack of observation of dinucleon decay, accounting for the appropriate suppression factors. In general we find that only operators which require 2 or more weak interactions to contribute to dinucleon decay can give baryon oscillations at a rate which is large enough to be relevant for either experimental searches or baryogenesis. The last column of the table gives the limit on the size of the operator that can be produced in our specific model when collider constraints on new particles are considered, which will be discussed in Sec. V.

IV. CP VIOLATION IN HEAVY BARYON OSCILLATIONS

The evolution of the $(\mathcal{B},\bar{\mathcal{B}})$ system in vacuum, assuming CPT conservation, can be described [11] by a 2×2 Hamiltonian,

$$\mathcal{H} = M - \frac{i}{2}\Gamma = \begin{pmatrix} M_{\mathcal{B}} - \frac{i}{2}\Gamma_{\mathcal{B}} & M_{12} - \frac{i}{2}\Gamma_{12} \\ M_{12}^{*} - \frac{i}{2}\Gamma_{12}^{*} & M_{\mathcal{B}} - \frac{i}{2}\Gamma_{\mathcal{B}} \end{pmatrix}. \quad (16)$$

M and Γ are both Hermitian matrices that describe the dispersive and absorptive parts of the $\mathcal{B},\bar{\mathcal{B}} \to \mathcal{B},\bar{\mathcal{B}}$ amplitude, respectively. This system is entirely analogous to the very well known case of neutral mesons and antimesons. Because of the off-diagonal terms in $\mathcal{H}$, the mass eigenstates $|\mathcal{B}_{\rm L,H}\rangle$ with masses $m_{\rm L,H}$ are linear combinations of the flavor eigenstates $|\mathcal{B}\rangle$ and $|\bar{\mathcal{B}}\rangle$,

$$|\mathcal{B}_{\rm L,H}\rangle = p|\mathcal{B}\rangle \pm q|\bar{\mathcal{B}}\rangle. \quad (17)$$

The mass difference is $\Delta m = m_{\rm H} - m_{\rm L} > 0$ and the width difference between the states is $\Delta\Gamma = \Gamma_{\rm H} - \Gamma_{\rm L}$ and can be of either sign. The flavor admixtures can be determined by

$$\left(\frac{q}{p}\right)^2 = \frac{M_{12}^{*} - (i/2)\Gamma_{12}^{*}}{M_{12} - (i/2)\Gamma_{12}}. \quad (18)$$

A state that begins at $t=0$ as a $|\mathcal{B}\rangle$ or $|\bar{\mathcal{B}}\rangle$ is at time t

$$\begin{aligned} |\mathcal{B}(t)\rangle &= g_{+}(t)|\mathcal{B}\rangle - \frac{q}{p}g_{-}(t)|\bar{\mathcal{B}}\rangle, \\ |\bar{\mathcal{B}}(t)\rangle &= g_{+}(t)|\bar{\mathcal{B}}\rangle - \frac{p}{q}g_{-}(t)|\mathcal{B}\rangle \end{aligned} \quad (19)$$

with

$$g_{\pm}(t) = \frac{1}{2}\left(e^{-im_{\rm H}t - \frac{1}{2}\Gamma_{\rm H}t} \pm e^{-im_{\rm L}t - \frac{1}{2}\Gamma_{\rm L}t}\right). \quad (20)$$

A particularly useful quantity that measures the level of CP and baryon number violation is the quantity,

$$A_{\mathcal{B}} = \frac{P_{\mathcal{B}\to\mathcal{B}} - P_{\mathcal{B}\to\bar{\mathcal{B}}} + P_{\bar{\mathcal{B}}\to\mathcal{B}} - P_{\bar{\mathcal{B}}\to\bar{\mathcal{B}}}}{P_{\mathcal{B}\to\mathcal{B}} + P_{\mathcal{B}\to\bar{\mathcal{B}}} + P_{\bar{\mathcal{B}}\to\mathcal{B}} + P_{\bar{\mathcal{B}}\to\bar{\mathcal{B}}}}, \quad (21)$$

where, e.g., $P_{\mathcal{B}\to\bar{\mathcal{B}}}$ is the time integrated probability for an intitial $\mathcal{B}$ state to oscillate into a $\bar{\mathcal{B}}$ and the other terms are defined analogously. In terms of the elements of $\mathcal{H}$, this can be concisely expressed,

$$A_{\mathcal{B}} = \frac{2\mathrm{Im}\left(M_{12}^{*}\Gamma_{12}\right)}{\Gamma_{\mathcal{B}}^2 + 4\left|M_{12}\right|^2}. \quad (22)$$

This expresses the familiar fact that CP violation requires a phase difference between the absorptive and dispersive parts of the transition amplitudes.

The dispersive part of the transition amplitude, M_{12}, is dominantly given by off-shell χ_i exchange in our model, as seen in Fig. 1. We have already written down what we need to estimate this in Eqs. (4) and (5), resulting in

$$M_{12} \sim \kappa^2 \sum_i \frac{m_{\chi_i}}{m_{\mathcal{B}}^2 - m_{\chi_i}^2}\left(\frac{g_{ud}^{*}y_{id'}}{m_\phi^2}\right)^2. \quad (23)$$

Here, u, d, and d' refer to the flavors that comprise $\mathcal{B}$ and we have assumed, if $d \neq d'$, that $g_{ud}^{*}y_{id'} \gg g_{ud'}^{*}y_{id}$. If we concentrate on the contribution due to a particular χ_i and express it in terms of its mass difference from the baryon, $\Delta m_{\mathcal{B}i} = m_{\mathcal{B}} - m_{\chi_i}$, we have

$$\begin{aligned} |M_{12}|_i &\sim \frac{\kappa^2}{2\Delta m_{\mathcal{B}i}}\left|\frac{g_{ud}^{*}y_{id'} + g_{ud'}^{*}y_{id}}{m_\phi^2}\right|^2 \\ &\simeq 8\times 10^{-16}\ \mathrm{GeV}\left(\frac{500\ \mathrm{MeV}}{\Delta m_{\mathcal{B}i}}\right) \\ &\times\left(\frac{600\ \mathrm{GeV}}{m_\phi/\sqrt{|g_{ud}^{*}y_{id'} + g_{ud'}^{*}y_{id}|}}\right)^4. \end{aligned} \quad (24)$$

The absorptive part of the transition amplitude requires an on-shell state into which both $\mathcal{B}$ and $\bar{\mathcal{B}}$ can decay. This requires at least χ_1 to be light enough for either baryon or antibaryon to decay into it. CP violation will be largest when the mass splitting between χ_1 and $\mathcal{B}$ is not too large. In this case the most important states for Γ_{12} are decays of $\mathcal{B}$ to χ_1 plus a meson. The contribution from $\chi_1\pi^0$, for instance, can be estimated using the effective Lagrangian,

$$\mathcal{L}_{\rm eff} \supset -y_{i\mathcal{B}}\pi^0\bar{\mathcal{B}}i\gamma^5\chi_i + \text{h.c.}, \quad (25)$$

where

$$y_{i\mathcal{B}} \sim \frac{4\pi\kappa}{m_{\mathcal{B}}}\frac{g_{ud}^{*}y_{id'}}{m_\phi^2}. \quad (26)$$

The factor of 4π in this expression accounts for the non-perturbative nature of the interaction, which is similar to the pion-nucleon vertex. This interaction gives a contribution to Γ_{12} of

$$\Gamma_{12} \sim \sum_i \frac{y_{i\mathcal{B}}^2 m_{\chi_i}}{32\pi}\left(1 + r_{\chi_i} - r_{\pi^0}\right)\lambda^{1/2}\left(1, r_{\chi_i}, r_{\pi^0}\right), \quad (27)$$

where $r_{\chi_i,\pi^0} = m_{\chi_i,\pi^0}^2/m_{\mathcal{B}}^2$ and $\lambda(a,b,c) = a^2+b^2+c^2-2ab-2ac-2bc$. The magnitude of the contribution to Γ_{12} from a particular χ_i is roughly

$$\begin{aligned} |\Gamma_{12}|_i &\sim \frac{y_{i\mathcal{B}}^2}{8\pi}\Delta m_{\mathcal{B}i} \sim \frac{2\pi\kappa^2}{m_{\mathcal{B}}^2}\left|\frac{g_{ud}^{*}y_{id'}}{m_\phi^2}\right|^2\Delta m_{\mathcal{B}i} \\ &\simeq 1\times 10^{-16}\ \mathrm{GeV}\left(\frac{\Delta m_{\mathcal{B}i}}{500\ \mathrm{MeV}}\right) \\ &\times\left(\frac{5\ \mathrm{GeV}}{m_{\mathcal{B}}}\right)^2\left(\frac{600\ \mathrm{GeV}}{m_\phi/\sqrt{|g_{ud}^{*}y_{id'}|}}\right)^4. \end{aligned} \quad (28)$$

Operator	$\mathcal{B}$	Weak Insertions Required	Measured Γ (GeV) [22]	Limits on $\delta_{\mathcal{B}\bar{\mathcal{B}}} = M_{12}$ (GeV) Dinucleon decay	Collider
$(udd)^2$	n	None	$(7.477 \pm 0.009) \times 10^{-28}$	10^{-33}	10^{-17}
$(uds)^2$	Λ	None	$(2.501 \pm 0.019) \times 10^{-15}$	10^{-30}	10^{-17}
$(uds)^2$	Σ^0	None	$(8.9 \pm 0.8) \times 10^{-6}$	10^{-30}	10^{-17}
$(uss)^2$	Ξ^0	One	$(2.27 \pm 0.07) \times 10^{-15}$	10^{-22}	10^{-17}
$(ddc)^2$	Σ_c^0	Two	$\left(1.83^{+0.11}_{-0.19}\right) \times 10^{-3}$	10^{-17}	10^{-16}
$(dsc)^2$	Ξ_c^0	Two	$\left(5.87^{+0.58}_{-0.61}\right) \times 10^{-12}$	10^{-16}	10^{-15}
$(ssc)^2$	Ω_c^0	Two	$(9.5 \pm 1.2) \times 10^{-12}$	10^{-14}	10^{-15}
$(udb)^2$	Λ_b^0	Two	$(4.490 \pm 0.031) \times 10^{-13}$	10^{-13}	10^{-17}
$(udb)^2$	Σ_b^{0*}	Two	$\sim 10^{-3*}$	10^{-13}	10^{-17}
$(usb)^2$	Ξ_b^0	Two	$(4.496 \pm 0.095) \times 10^{-13}$	10^{-10}	10^{-17}
$(dcb)^2$	$\Xi_{cb}^{0\ \dagger}$	Two	$\sim 10^{-12\dagger}$	10^{-17}	10^{-15}
$(scb)^2$	$\Omega_{cb}^{0\ \dagger}$	Two	$\sim 10^{-12\dagger}$	10^{-14}	10^{-15}
$(ubb)^2$	$\Xi_{bb}^{0\ \ddagger}$	Four	$\sim 10^{-13\ddagger}$	>1	10^{-17}
$(cbb)^2$	$\Omega_{cbb}^{0\ \dagger}$	Four	$\sim 10^{-12\dagger}$	>1	10^{-15}

TABLE I. Operators that mediate $\mathcal{B} \leftrightarrow \bar{\mathcal{B}}$ oscillations and the number of weak interaction insertions required for each of these to contribute to dinucleon decay. The resulting limit from dinucleon decay on the transition amplitude, defined in Eq. (5), for each operator is shown. An $*$ indicates a baryon that has not yet been observed and which has a strong decay channel open. A † (‡) indicates an unobserved baryon which primarily decays through a weak interaction of a c (b) quark.

We now see the reason we need at least two Majorana fermions. If there were only a single Majorana fermion, χ_1, that contributed to M_{12} and Γ_{12}, they would have the same phase and $A_{\mathcal{B}}$ in Eq. (22) would vanish. Thus, we use contributions from χ_1 and χ_2 exchange to obtain a physical, CP-violating phase difference between M_{12} and Γ_{12}.

The ratio of the single meson contribution to Γ_{12} from χ_1 to its contribution to M_{12} is

$$\begin{aligned} \left|\frac{\Gamma_{12}}{M_{12}}\right|_1 &\sim 4\pi \left(\frac{\Delta m_{\mathcal{B}1}}{m_{\mathcal{B}}}\right)^2 \\ &\simeq 0.1 \left(\frac{\Delta m_{\mathcal{B}1}}{500\ \mathrm{MeV}}\right)^2 \left(\frac{5\ \mathrm{GeV}}{m_{\mathcal{B}}}\right)^2 . \end{aligned} \tag{29}$$

The CP-violating quantity $A_{\mathcal{B}}$ in Eq. (22) linearly depends on $|\Gamma_{12}|$. Without finely tuning the contributions due to χ_1 and χ_2 against each other, this value of the ratio due to χ_1 alone is roughly as large as the total ratio $|\Gamma_{12}/M_{12}|$ can get.

V. COLLIDER CONSTRAINTS

To obtain a large amount of CP violation in heavy baryon oscillation, it will be clear that the lightest two Majorana fermions must have masses on the order of a few GeV along with couplings to quarks that are not too small. In this discussion, we consider the two lightest Majorana fermions. The third, χ_3 must be weakly coupled in the minimal version of the model, due to cosmological considerations as we will see in Sec. VII.

The constraints that we will discuss in this section require that ϕ have a mass of at least a few hundred GeV. In this case, χ_i decays can be analyzed by integrating out the scalar in Eq. (1) resulting in four-fermion interactions,

$$-\frac{g_{ud}y_{id'}}{m_\phi^2}\bar{\chi}_i \bar{u}_R d_R^c d_R'^c + \mathrm{h.c.} \tag{30}$$

(For $d \neq d'$, we have again assumed that $g_{ud}y_{id'} \gg g_{ud'}y_{id}$.) Both the interactions responsible for the decay of the Majorana fermions and those that source CP-violating baryon oscillations are of the same form. The quarks involved in the decay operator must be lighter than those responsible for baryon oscillations and the couplings responsible for decay must be relatively smaller, to avoid stronger dinucleon decay limits from $\Delta B = 2$ quarks involving light quarks.

The interaction in Eq. (30) allows for the decay $\chi_i \to udd'$, where u, d, and d' are up- and down-type quarks light enough for this to be kinematically allowed. It is reasonable to assume that one mode dominates their allowed branchings and in this case, their lifetimes are

$$\begin{aligned} \tau_{\chi_i} &\sim \frac{2\,(8\pi)^3}{m_{\chi_{1,2}}^5}\left|\frac{m_\phi^2}{g_{ud}\,y_{id'}}\right|^2 \\ &\simeq 10^{-6}\ \mathrm{s}\left(\frac{5\ \mathrm{GeV}}{m_{\chi_i}}\right)^5 \left(\frac{m_\phi/\sqrt{g_{ud}\,y_{id'}}}{20\ \mathrm{TeV}}\right)^4 . \end{aligned} \tag{31}$$

For $m_{\chi_i} = 5$ GeV, with couplings $g_{ud}\,y_{id'} \lesssim (m_\phi/20\ \mathrm{TeV})^2$ the $\Delta B = 2$ transition amplitude in the udd' system is less than 10^{-22} GeV, avoiding conflict with constraints from dinucleon decay (see Table I).

Furthermore, if $m_{\chi_i} = 5$ GeV, as long as $g_{ud}\, y_{id'} \gtrsim (m_\phi/350\ \text{TeV})^2$, $\tau_{\chi_i} \lesssim 0.1$ s which is a short enough lifetime to avoid spoiling successful BBN (see, e.g., [23]).

We might ask whether instead of ensuring that $\chi_{1,2}$ decay fast enough to avoid spoiling BBN, the lightest fermion χ_1 could instead be long enough lived to serve as dark matter. First we note that the range allowed for kinematic stability of both proton and χ_1 is extremely fine tuned, with the χ_1 mass between $m_p - m_e$ and $m_p + m_e$. If we assume all the χ's participate in a viable heavy flavor baryogenesis mechanism, we will see that we must require the $\chi_{1,2}$ masses to be around 3–5 GeV, and we also need sufficiently large four-fermion interactions involving $\chi_{1,2}$ and heavy flavor quarks, suppressed by a scale $\Lambda_{\text{heavy}} \equiv m_\phi/\sqrt{gy} \sim 600$ GeV. Here g and y here label couplings with the relevant flavor structure. Based on our discussion in Sec. III, at the one loop level we must generate four-fermion operators involving light quarks (into which $\chi_{1,2}$ can decay) are generated with a scale $\Lambda_{\text{light}} \gtrsim 10^4 \Lambda_{\text{heavy}} \sim 10^6$ GeV. This provides a lower bound on the strength of the light quark four-fermion operator which gives an upper bound on the $\chi_{1,2}$ lifetime of $\tau_{\chi_{1,2}} \lesssim 1000$ s in the absence of fine-tuning against some other source of this operator. There is also an unavoidable decay channel that comes from the mixing of $\chi_{1,2}$ with the heavy flavor baryon, $\mathcal{B}$, whose oscillations are responsible for the BAU, with a mixing angle $\theta \sim |M_{12}|/\Delta m$ where Δm is the mass splitting between the Majorana fermions and $\mathcal{B}$. This mixing leads to the decay of $\chi_{1,2}$ into $\mathcal{B}$'s decay channels with a partial width proportional to $\theta^2\Gamma_{\mathcal{B}}$, which is much shorter than the lifetime of the universe.

We see that the Majorana fermions are generically unstable but long-lived on the scale of collider experiments and appear as missing energy. Decay lengths on the order of 10^2 to 10^7 m are expected, potentially relevant for the recently proposed MATHUSLA detector [24] which is optimized to search for long-lived particles. In what follows, to analyze collider constraints on the new scalar ϕ we will assume that any χ_i produced at a collider is invisible and defer discussion of the displaced decay signatures at, e.g., MATHUSLA.

Now that we know that the χ_i's are invisible at colliders, we can understand how the scalars appear when produced in hadron collisions. Because ϕ is a color fundamental, if it is kinematically accessible, $\phi\phi^*$ pairs are easily produced in proton-(anti)proton collisions, and the signatures are essentially those of squarks in RPV SUSY. In addition to QCD production, (single) scalars can be resonantly produced in the presence of some nonzero g_{ud}. Once produced, the scalar decays through one of the interactions in Eq. (1), either to quark pairs with a rate

$$\Gamma_{\phi\to\bar{u}\bar{d}} \simeq \sum_{i,j} \frac{\left|g_{u_i d_j}\right|^2}{16\pi} m_\phi, \tag{32}$$

or to χ_i plus a quark,

$$\Gamma_{\phi\to\chi d} \simeq \sum_{i,j} \frac{\left|y_{i d_j}\right|^2}{16\pi} m_\phi, \tag{33}$$

where we have assumed that m_ϕ is much larger than the mass of any decay product. Therefore, these scalars can appear in searches for dijet resonances (either singly or pair produced) and (mono)jets and missing energy. Which search is most sensitive depends on m_ϕ and the branching fractions for $\phi \to \bar{u}\bar{d}$ and $\phi \to \chi d$.

Taken together, LHC searches for pair produced dijet resonances, both with [25] and without [26] heavy flavor in the final states, as well as standard SUSY searches for (b-tagged [27, 28] or not [28, 29]) jets plus missing energy rule out ϕ masses below about 400 GeV. Above this mass, limits from pair produced dijet resonances are no longer constraining while resonant production of a single ϕ with a rate proportional to $|g_{ud}|^2$ for some u, d is important [30]. We use the limits from resonant dijet production from [30] and recast searches for jets and missing energy [27–29] as well as monojets [31] to find limits on the couplings g_{ud} and $y_{id'}$ as functions of m_ϕ. We find this limit for every flavor u, d, and d', assuming that only g_{ud} and $y_{id'}$ are relevant. Given these limits, the maximum value of the product of couplings $g_{ud}y_{id'}$ at each m_ϕ can be found, and taking a value of the mass splitting between χ_i and the udd' baryon, which can be turned into an upper limit on the transition amplitude $\delta_{udd'} = M_{12}$ in the udd' system.[4] We show the upper limit on M_{12} as a function of m_ϕ for each pattern of flavors u, d, and d', assuming the dominance of one particular pair of couplings g_{ud}, $y_{id'}$ and a mass splitting between m_{χ_i} and the udd' baryon of 200 MeV in Fig. 4. We also show the largest value of M_{12} allowed from collider searches in each neutral baryon system in Table I.

Lastly, we note that the six-quark $\Delta B = 2$ operators themselves can lead to interesting signatures at the LHC. These were studied in Ref. [32].

VI. HADRON PHENOMENOLOGY

A. Hadron decays

After integrating out the heavy colored scalars, four-fermion interactions between the Majorana fermions and quarks are generated as in Eq. (30). These can lead to new decays of hadrons to final states that differ in baryon

[4] Note that we perform this scan for a single Majorana fermion. Including a second, as we must to obtain CP violation, does not change the allowed values by more than an $\mathcal{O}(1)$ factor which, given our level of precision, is unimportant. The third, χ_3 must be more weakly coupled than $\chi_{1,2}$ and can be even more safely neglected here.

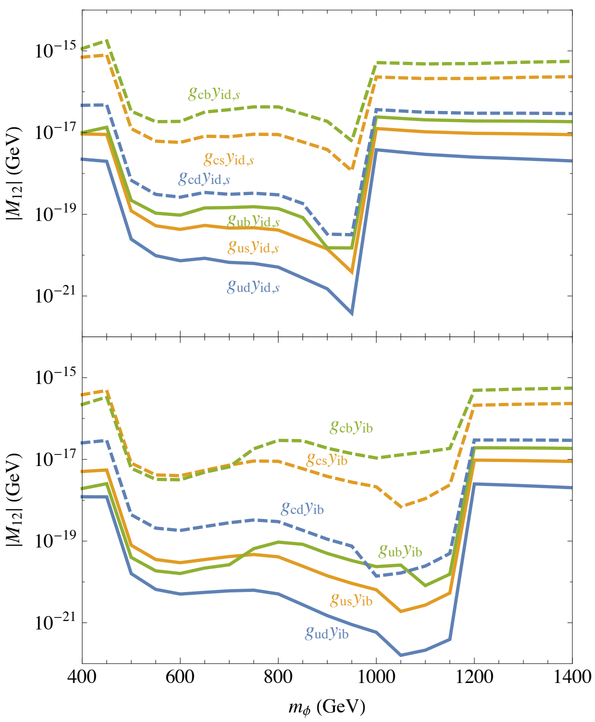

FIG. 4. Upper limits on M_{12} as functions of m_ϕ that result from collider searches for diject resonances and jets plus missing energy, assuming the dominance of the product of couplings $g_{ud}y_{id'}$ indicated, where u and $d^{(\prime)}$ label generic up- and down-type quarks, respectively. We have taken $\Delta m_{\mathcal{B}} = m_{\mathcal{B}} - m_\chi = 200$ MeV. Top: The limits when y_{id} or y_{is} are dominant. Bottom: The limits when y_{ib} is dominant. Solid curves show the limits in the case where the charge 2/3 quark involved is u while dashed lines show the limit in the case of the c quark.

number by one unit along with any kinematically accessible χ_i, e.g.,

$$\begin{aligned} &\text{meson} \to \text{baryon} + \chi_i\,[+\,\text{meson(s)}], \\ &\text{baryon} \to \text{meson(s)} + \chi_i. \end{aligned} \tag{34}$$

As we showed in Sec. V, on the scale of particle physics experiments, χ_i appear as missing energy.

For definiteness, let us focus now on four-fermion interactions that involve the b quark and the lightest Majorana fermion. This is potentially relevant to the case where baryons containing b quarks undergo CP-violating oscillations in the early Universe, producing the BAU; operators involving heavy quarks are less constrained by dinucleon decay and are therefore more promising candidates, cf. Table I. Similar considerations apply for operators involving lighter quarks.

Consider, as a definite example, b decays through the operator

$$-\frac{g_{ub}y_{1d}}{m_\phi^2}\chi_1 u_R d_R b_R, \tag{35}$$

where u and d here are the actual up and down quarks. (We have omitted the contribution to this operator from $g_{ud}y_{1b}$ which is more constrained by collider searches.) The rate for the b quark to decay through such an interaction is

$$\begin{aligned} \Gamma_{b\to\chi_1\bar{u}\bar{d}} &\sim \frac{m_b \Delta m^4}{60\,(2\pi)^3}\left(\frac{g_{ub}y_{1d}}{m_\phi^2}\right)^2 + \mathcal{O}\left(\frac{\Delta m^5}{m_b^5}\right) \\ &\simeq 2\times 10^{-15}\ \text{GeV}\left(\frac{\Delta m}{2\ \text{GeV}}\right)^4\left(\frac{1.2\ \text{TeV}}{m_\phi/\sqrt{g_{ub}y_{1d}}}\right)^4. \end{aligned} \tag{36}$$

In this expression Δm is the mass splitting between χ_1 and the bottom quark (we have ignored masses in the final state besides m_{χ_1}). We have chosen to normalize this expression on values of the mass splitting and $m_\phi/\sqrt{g_{ub}y_{1d}}$ that result in a transition amplitude of $|M_{12}| \sim 10^{-17}$ GeV in the $\Lambda_b^0 = (udb)$ baryon system, which is the rough collider limit. Given this mass splitting, this can lead to decays of B^+ mesons to a nucleon and χ_1 with a branching ratio of

$$\begin{aligned} \text{Br}_{B^\pm\to N\chi_1+X} \sim 6\times 10^{-3}&\left(\frac{\Delta m}{2\ \text{GeV}}\right)^4 \\ &\times\left(\frac{1.2\ \text{TeV}}{m_\phi/\sqrt{g_{ub}y_{1d}}}\right)^4, \end{aligned} \tag{37}$$

where X represents possible additional pions. This is not a small branching fraction, although final states of this form have not yet been searched for in B meson decays. However, the requirement that the final state hadrons carry baryon number means that this decay is kinematically forbidden if $m_{\chi_1} > m_{B^\pm} - m_p = 4.34$ GeV. Decays of bottom baryons would be allowed to proceed for splittings down to m_π, and one could expect branching ratios on the order of 10^{-3} for the parameters in Eq. (37).[5]

We also expect "wrong sign" decays of heavy baryons in this model, following a $\mathcal{B}\to\bar{\mathcal{B}}$ oscillation, with a branching fraction that is roughly

$$\frac{1}{2}\frac{|M_{12}|^2}{\Gamma_{\mathcal{B}}^2}. \tag{38}$$

Consider, e.g., the Ω_c^0. Given the constraints that appear in Table I, this branching could potentially be as large as 10^{-7}. The Belle II experiment hopes to collect ~ 50 ab^{-1} of e^+e^- data at $\sqrt{s} = 10.56$ GeV collecting about 50×10^9 B meson pairs. If Ω_c^0 baryons are produced in 2% of B

[5] The calculation of the baryon decay rate to χ_1 and a single meson is essentially the same as that of Γ_{12} in Sec. IV, modulo a factor of $m_b/m_{\chi_1} \sim \mathcal{O}(1)$.

meson decays (comparable to the measured production of Λ_c baryons), then there would be a sample of about 10^9 Ω_c^0's and $\bar{\Omega}_c^0$'s. Thus, there could be a few hundred "wrong sign" decays in the data sample. While this would be a challenging measurement, it is interesting that it is in principle observable at the next generation B-factory given current experimental limits.

We mention here that baryon-number–violating decays of baryons along these lines have been searched for by the CLAS Collaboration [33]. The branching fraction for $\Lambda \to K_S^0 + \text{inv.}$ is limited to less than 2×10^{-5} while that for $\Lambda \to \bar{p}\pi^+$ must be less than 9×10^{-7} which are sensitive to the operator $(uds)^2$. While interesting, these limit $\delta_{(uds)^2} = |M_{12}|$ to less than about 10^{-18} GeV, which is less strong than the limit on this operator from null searches for dinucleon decay. In light of the less stringent limits from dinucleon decay on operators involving heavy flavor, it would be highly desirable for searches for $\Delta B = 2$ decays of baryons with heavy quarks to be performed.

B. Meson oscillations

In addition to the decays described above, the new interactions could lead to flavor-changing oscillations of neutral mesons. The limits from these processes on this model were considered in Ref. [12]; we refer the reader to [12] and references therein for further details.

Avoiding these constraints requires a suppression of particular combinations of flavor-violating couplings. For example, considering Kaon oscillations, given $m_\phi \gtrsim 400$ GeV, y_{s1} and y_{s2} could be $\mathcal{O}(1)$ provided $y_{d1}, y_{d2} \lesssim 10^{-2}$. Similar considerations apply for flavor-violating combinations of the couplings g_{us} and g_{ud}. The constraints on charm and bottom couplings from D and B oscillations are less severe.

From the model building point of view, flavor-changing meson oscillations can be naturally avoided, e.g. charging the scalar under a symmetry so that $F_1 - F_2$ is conserved, where $F_{1,2}$ label flavor quantum numbers.

VII. COSMOLOGICAL PRODUCTION OF THE BARYON ASYMMETRY

We now answer in detail the question of how the baryon asymmetry of the Universe is produced in this model. In addition to the CP and baryon number violation described above, a nonzero asymmetry requires a departure from thermal equilibrium. The simplest possibility for this is to assume that χ_3 is very weakly coupled. It is therefore long-lived and decays out of equilibrium, producing the baryons that undergo CP- and B-violating oscillations.

At temperatures below m_{χ_3}, the equations that determine the radiation and χ_3 energy densities are

$$\frac{d\rho_{\rm rad}}{dt} + 4H\rho_{\rm rad} = \Gamma_{\chi_3}\rho_{\chi_3}, \tag{39}$$

$$\frac{d\rho_{\chi_3}}{dt} + 3H\rho_{\chi_3} = -\Gamma_{\chi_3}\rho_{\chi_3}. \tag{40}$$

H is the Hubble parameter which is related to the total energy density,

$$H = \sqrt{\frac{8\pi}{3}\frac{\rho}{M_{\rm Pl}^2}} \simeq \sqrt{\frac{8\pi}{3}\frac{\rho_{\rm rad}+\rho_{\chi_3}}{M_{\rm Pl}^2}}, \tag{41}$$

where $M_{\rm Pl} = 1.22 \times 10^{19}$ GeV is the Planck mass. In the absence of χ_3 decays, $\rho_{\rm rad}$ and ρ_{χ_3} simply redshift like radiation and matter energy densities, respectively. The right-hand sides of these equations describe how χ_3 decays cause the energy density in matter to decrease while dumping energy into the plasma.

In addition to depositing energy in the plasma, some of the χ_3 decays produce baryons and antibaryons, $\mathcal{B}$ and $\bar{\mathcal{B}}$, that can oscillate and decay, violating CP and B. For this to occur, the temperature of the Universe needs to be below the QCD confinement temperature, $T_{\rm QCD} \simeq 200$ MeV. On the timescale of the expansion of the Universe, H^{-1}, the (anti-)baryons produced this way rapidly oscillate and decay, producing a net B asymmetry. However, because of the presence of the plasma, with which they can interact, as well as their large annihilation cross section, $\mathcal{B} \leftrightarrow \bar{\mathcal{B}}$ can decohere in this environment, suppressing the asymmetry that is generated. Properly accounting for this requires a density matrix treatment, which has been used in a cosmological context for neutrino oscillations and oscillating asymmetric dark matter [34, 35]. Following Ref. [35] (see [36] for a similar analysis in the context of baryogenesis-related oscillations), we can write the Boltzmann equations that govern the evolution of the number density of the $\mathcal{B}$-$\bar{\mathcal{B}}$ system,

$$\begin{aligned}\frac{dn}{dt} + 3Hn = &-i\left(\mathcal{H}n - n\mathcal{H}^\dagger\right) - \frac{\Gamma_\pm}{2}\left[O_\pm,[O_\pm,n]\right] \\ &- \langle\sigma v\rangle_\pm\left(\frac{1}{2}\left\{n, O_\pm\bar{n}O_\pm\right\} - n_{\rm eq}^2\right) \\ &+ \frac{1}{2}\frac{\Gamma_{\chi_3}\rho_{\chi_3}}{m_{\chi_3}}{\rm Br}_{\chi_3\to\mathcal{B}}O_+,\end{aligned} \tag{42}$$

where the last term describes $\mathcal{B}$ and $\bar{\mathcal{B}}$ production through χ_3 decay. ${\rm Br}_{\chi_3\to\mathcal{B}}$ is the branching ratio for χ_3 to decay to $\mathcal{B}$ or $\bar{\mathcal{B}}$. In this equation n and $\bar{n}$ are density matrices,

$$n = \begin{pmatrix} n_{\mathcal{B}\mathcal{B}} & n_{\mathcal{B}\bar{\mathcal{B}}} \\ n_{\bar{\mathcal{B}}\mathcal{B}} & n_{\bar{\mathcal{B}}\bar{\mathcal{B}}} \end{pmatrix}, \; \bar{n} = \begin{pmatrix} n_{\bar{\mathcal{B}}\bar{\mathcal{B}}} & n_{\mathcal{B}\bar{\mathcal{B}}} \\ n_{\bar{\mathcal{B}}\mathcal{B}} & n_{\mathcal{B}\mathcal{B}} \end{pmatrix}, \tag{43}$$

and $n_{\rm eq}$ is the equilibrium density of baryons plus antibaryons. $\mathcal{H}$ is the Hamiltonian seen in Eq. (16). $\langle\sigma v\rangle_\pm$

and $\Gamma_\pm$ are thermally-averaged annihilation cross sections and scattering rates on the plasma, respectively. $O_\pm$ is a matrix

$$O_\pm = \begin{pmatrix} 1 & 0 \\ 0 & \pm 1 \end{pmatrix}. \tag{44}$$

The subscript of $\langle\sigma v\rangle_\pm$ and $\Gamma_\pm$, i.e. whether they appear with O_+ or O_- in Eq. (42), is determined by the behavior of the effective Lagrangian that gives rise to these interactions under charge conjugation of *only* the heavy baryons, $\mathcal{B} \leftrightarrow \bar{\mathcal{B}}$, $\mathcal{L}_{\rm eff} \leftrightarrow \pm\mathcal{L}_{\rm eff}$. Interactions that do not change sign are said to be flavor-blind while those that do are flavor-sensitive. For example, $\mathcal{B}$ and $\bar{\mathcal{B}}$ can scatter on light charged particles in the plasma through their magnetic moment, μ, which corresponds to a term in the effective Lagrangian of

$$\frac{i\mu}{4}\bar{\mathcal{B}}\left[\gamma^\nu, \gamma^\rho\right]\mathcal{B}F_{\nu\rho}. \tag{45}$$

Under $\mathcal{B} \leftrightarrow \bar{\mathcal{B}}$ this term changes sign, so the rate for scattering via the magnetic moment appears with O_- in the Botzmann equation.

It is useful to work in terms of the quantities

$$\begin{aligned} &\Sigma \equiv n_{\mathcal{B}\mathcal{B}} + n_{\bar{\mathcal{B}}\bar{\mathcal{B}}},\ \Delta \equiv n_{\mathcal{B}\mathcal{B}} - n_{\bar{\mathcal{B}}\bar{\mathcal{B}}}, \\ &\Xi \equiv n_{\mathcal{B}\bar{\mathcal{B}}} - n_{\bar{\mathcal{B}}\mathcal{B}},\ \Pi \equiv n_{\mathcal{B}\bar{\mathcal{B}}} + n_{\bar{\mathcal{B}}\mathcal{B}}. \end{aligned} \tag{46}$$

In this basis the Boltzmann equations are

$$\begin{aligned} \left(\frac{d}{dt} + 3H\right)\Sigma = {} & \frac{\Gamma_{\chi_3}\rho_{\chi_3}}{m_{\chi_3}}\mathrm{Br}_{\chi_3\to\mathcal{B}} - \Gamma_{\mathcal{B}}\Sigma \\ & - \left(\mathrm{Re}\,\Gamma_{12}\right)\Pi + i\left(\mathrm{Im}\,\Gamma_{12}\right)\Xi \\ & - \frac{1}{2}\Big[\left(\langle\sigma v\rangle_+ + \langle\sigma v\rangle_-\right)\left(\Sigma^2 - \Delta^2 - 4n_{\rm eq}^2\right) \\ & \quad + \left(\langle\sigma v\rangle_+ - \langle\sigma v\rangle_-\right)\left(\Pi^2 - \Xi^2\right)\Big], \\ \left(\frac{d}{dt} + 3H\right)\Delta = {} & -\Gamma_{\mathcal{B}}\Delta + 2i\left(\mathrm{Re}\,M_{12}\right)\Xi + 2\left(\mathrm{Im}\,M_{12}\right)\Pi, \\ \left(\frac{d}{dt} + 3H\right)\Xi = {} & -\left(\Gamma_{\mathcal{B}} + 2\Gamma_- + \langle\sigma v\rangle_+\Sigma\right)\Xi \\ & + 2i\left(\mathrm{Re}\,M_{12}\right)\Delta - i\left(\mathrm{Im}\,\Gamma_{12}\right)\Sigma, \\ \left(\frac{d}{dt} + 3H\right)\Pi = {} & -\left(\Gamma_{\mathcal{B}} + 2\Gamma_- + \langle\sigma v\rangle_+\Sigma\right)\Pi \\ & - 2\left(\mathrm{Im}\,M_{12}\right)\Delta - \left(\mathrm{Re}\,\Gamma_{12}\right)\Sigma. \end{aligned} \tag{47}$$

Coherent oscillations from a flavor-symmetric state to an asymmetric state proceed through $\Sigma \to \Xi, \Pi \to \Delta$. Flavor-sensitive scattering and flavor-blind annihilation suppress Ξ and Π and therefore lead to decoherence.

When they decay, $\mathcal{B}$ and $\bar{\mathcal{B}}$ create states that carry baryon number. The flavor-asymmetric configuration contributes to the difference between the baryon and antibaryon number densities,

$$\left(\frac{d}{dt} + 3H\right)\left(n_B - n_{\bar{B}}\right) = \Gamma_{\mathcal{B}}\Delta. \tag{48}$$

The dominant interaction of $\mathcal{B}$ and $\bar{\mathcal{B}}$ with the plasma is scattering on charged particles (mostly electrons at $T \lesssim 100$ MeV) via the magnetic moment term in Eq. (45). The cross section for this at temperatures well below $m_{\mathcal{B}}$ is

$$\frac{d\sigma_{\rm sc}}{d\Omega} = \alpha^2\mu^2\left(\frac{1 + \sin^2\theta/2}{\sin^2\theta/2}\right) \tag{49}$$

which diverges at small scattering angle, $\theta \to 0$. This divergence is cut off at finite temperature by the inverse photon screening length, m_γ. Using this, the total cross section can be estimated as

$$\sigma_{\rm sc} \sim 4\pi\alpha^2\mu^2\log\left(\frac{4E^2}{m_\gamma^2}\right) \tag{50}$$

where E is the electron energy. Taking $E \sim T$, $m_\gamma \sim eT/3$ [37] and $\mu \sim 1/(2m_{\mathcal{B}})$, this gives

$$\sigma_{\rm sc} \sim \frac{\pi\alpha^2}{m_{\mathcal{B}}^2}\log\left(\frac{9}{\pi\alpha}\right). \tag{51}$$

The (flavor-sensitive) scattering rate is therefore

$$\begin{aligned} \Gamma_- = \Gamma_{\rm sc} &\sim \sigma_{\rm sc}\left(n_{e^-} + n_{e^+}\right) \\ &\sim \frac{\pi\alpha^2}{m_{\mathcal{B}}^2}\log\left(\frac{9}{\pi\alpha}\right) \times \frac{3\zeta(3)}{\pi^2}T^3 \\ &\sim 10^{-11}\ \mathrm{GeV}\left(\frac{5\ \mathrm{GeV}}{m_{\mathcal{B}}}\right)^2\left(\frac{T}{10\ \mathrm{MeV}}\right)^3. \end{aligned} \tag{52}$$

At temperatures above a few MeV, as is needed for BBN, this rate is larger than a typical heavy baryon width and therefore strongly affects the $\mathcal{B}$-$\bar{\mathcal{B}}$ oscillations.

When solving the Boltzmann equations, we take an annihilation cross section that is similar to that for $p\bar{p}$ annihilation at low energies,

$$\langle\sigma v\rangle_+ + \langle\sigma v\rangle_- = 400\ \mathrm{mb}. \tag{53}$$

We will find that only the total annihilation cross section and not whether it is flavor-blind or -sensitive is important, since $\Sigma \gg \Delta, \Xi, \Pi$. Furthermore the annihilation rate is always much smaller than the scattering rate at temperatures we are interested in, so its effect on the final asymmetry is subdominant and can generally be ignored.

A. Sudden Decay Approximation

Having removed the heavy baryons from the problem due to the short timescales in their system, the evolution equations are Eqs. (39), (40), and (59). These involve only the radiation energy density, χ_3 density, and the baryon asymmetry. They can be simply studied using a sudden decay approximation to gain a rough estimate of the baryon asymmetry. We outline this estimate below.

At some high temperature above m_{χ_3}, we assume that χ_3 was in thermal equilibrium with the plasma, fixing its number density for $T \lesssim m_{\chi_3}$ to roughly

$$n_{\chi_3} \simeq \frac{3}{4}\frac{\zeta(3)}{\pi^2}T^3. \tag{54}$$

As the Universe cools the energy density in χ_3 and radiation are equal. This occurs at the temperature

$$T_{\rm eq} = \frac{45\zeta(3)}{2\pi^4 g_*(T_0)} m_{\chi_3}. \tag{55}$$

g_* is the effective number of relativistic degrees of freedom and here it is evaluated at $T_0 \gtrsim m_{\chi_3}$. This corresponds to the time

$$\begin{aligned} t_{\rm eq} &= \sqrt{\frac{45}{16\pi^3 g_*(T_{\rm eq})}}\frac{M_{\rm Pl}}{T_{\rm eq}^2} \\ &= \frac{1}{\sqrt{5\pi g_*(T_{\rm eq})}}\frac{\pi^7 g_*(T_0)^2}{135\zeta(3)^2}\frac{M_{\rm Pl}}{m_{\chi_3}^2}. \end{aligned} \tag{56}$$

After this the Universe is matter dominated and the energy density in radiation and χ_3 redshift as

$$\rho_{\rm rad} = \frac{1}{2}\rho_{\rm eq}\left(\frac{t_{\rm eq}}{t}\right)^{8/3}, \ \rho_{\chi_3} = \frac{1}{2}\rho_{\rm eq}\left(\frac{t_{\rm eq}}{t}\right)^{2}. \tag{57}$$

We then assume that all of the χ_3's decay at the time $t_{\rm dec} = 1/\Gamma_{\chi_3}$. The ratio of the energy densities just before decay is

$$\begin{aligned} \xi &\equiv \frac{\rho_{\chi_3}(t_{\rm dec}^-)}{\rho_{\rm rad}(t_{\rm dec}^-)} = (t_{\rm eq}\Gamma_{\chi_3})^{-2/3} \\ &= 15\left[\frac{g_*(T_{\rm eq})}{g_*(T_0)}\right]^{1/3}\left[\frac{50}{g_*(T_0)}\right] \\ &\quad \times \left(\frac{m_{\chi_3}}{10\ \rm GeV}\right)^{4/3}\left(\frac{10^{-22}\ \rm GeV}{\Gamma_{\chi_3}}\right)^{2/3}. \end{aligned} \tag{58}$$

We use $t_{\rm dec}^-$ here to indicate the time infinitesimally before decay.

The dominance of the scattering rate over other scales in the problem allows us to make some simplifications of the evolution equations that are useful here. In this limit we can ignore the Hubble rate as well as annihilation and the equations governing $\mathcal{B}$ and $\bar{\mathcal{B}}$ in (47) can be integrated. This results in the evolution equation for the difference between baryon and antibaryon densities, Eq. (48), becoming

$$\begin{aligned} \left(\frac{d}{dt} + 3H\right)(n_B - n_{\bar{B}}) &= \frac{\Gamma_{\chi_3}\rho_{\chi_3}}{m_{\chi_3}} \\ &\quad \times \frac{2{\rm Im}\left(M_{12}^*\Gamma_{12}\right){\rm Br}_{\chi_3\to\mathcal{B}}}{\Gamma_{\mathcal{B}}\left(\Gamma_{\mathcal{B}} + 2\Gamma_-\right) + 4\left|M_{12}\right|^2} \\ &\simeq \frac{\Gamma_{\chi_3}\rho_{\chi_3}}{m_{\chi_3}}\frac{\Gamma_{\mathcal{B}}}{2\Gamma_-}\epsilon, \end{aligned} \tag{59}$$

which is valid for the cases we consider with $|M_{12}| \ll \Gamma_{\mathcal{B}} \ll \Gamma_-$. We have defined

$$\epsilon \equiv \frac{2{\rm Im}\left(M_{12}^*\Gamma_{12}\right)}{\Gamma_{\mathcal{B}}^2}{\rm Br}_{\chi_3\to\mathcal{B}} \simeq A_{\mathcal{B}}{\rm Br}_{\chi_3\to\mathcal{B}}, \tag{60}$$

with $A_{\mathcal{B}}$ from Eq. (22).

Using ϵ, we can then relate the baryon asymmetry to the χ_3 number density at decay,

$$\begin{aligned} \eta_B &= \frac{n_B - n_{\bar{B}}}{s(t_{\rm dec}^+)} = \frac{n_{\chi_3}(t_{\rm dec}^-)}{s(t_{\rm dec}^-)}\left[\frac{T(t_{\rm dec}^-)}{T(t_{\rm dec}^+)}\right]^3\frac{\Gamma_{\mathcal{B}}}{2\Gamma_-}\epsilon \\ &= \frac{3}{4}\frac{T(t_{\rm dec}^-)}{m_{\chi_3}}\xi\left[\frac{T(t_{\rm dec}^-)}{T(t_{\rm dec}^+)}\right]^3\frac{\Gamma_{\mathcal{B}}}{2\Gamma_-}\epsilon. \end{aligned} \tag{61}$$

Here, $t_{\rm dec}^+$ is the time just after decay. The ratio of the temperatures just before and after decay is determined by $\rho_{\rm rad}(t_{\rm dec}^+) = (1+\xi)\rho_{\rm rad}(t_{\rm dec}^-)$ so that

$$\frac{T(t_{\rm dec}^-)}{T(t_{\rm dec}^+)} = (1+\xi)^{-1/4} \simeq \xi^{-1/4}, \tag{62}$$

and

$$\eta_B \simeq \frac{3}{4}\frac{\xi^{1/4}T(t_{\rm dec}^-)}{m_{\chi_3}}\frac{\Gamma_{\mathcal{B}}}{2\Gamma_-}\epsilon. \tag{63}$$

The temperature just before decay can be arrived at by evolving the radiation energy density, resulting in

$$\eta_B \simeq \frac{3}{8}\sqrt{\frac{3}{\pi}}\left[\frac{5}{2\pi g_*(T_{\rm dec})}\right]^{1/4}\frac{\sqrt{M_{\rm Pl}\Gamma_{\chi_3}}}{m_{\chi_3}}\frac{\Gamma_{\mathcal{B}}}{2\Gamma_-}\epsilon. \tag{64}$$

Using the expression for the scattering rate in Eq. (52) evaluated at $T(t_{\rm dec}^+)$,

$$\begin{aligned} \eta_B &\simeq \frac{\pi^3}{3\zeta(3)}\sqrt{\frac{\pi g_*(T_{\rm dec})}{10}}\frac{\Gamma_{\mathcal{B}}\epsilon}{\sigma_{\rm sc}m_{\chi_3}\Gamma_{\chi_3}M_{\rm Pl}} \\ &\approx 9\times 10^{-11}\left[\frac{g_*(T_{\rm dec})}{50}\right]^{1/2}\left(\frac{m_{\mathcal{B}}}{5\ \rm GeV}\right)^2\left(\frac{\Gamma_{\mathcal{B}}}{10^{-13}\ \rm GeV}\right) \\ &\quad \times \left(\frac{8\ \rm GeV}{m_{\chi_3}}\right)\left(\frac{10^{-22}\ \rm GeV}{\Gamma_{\chi_3}}\right)\left(\frac{\epsilon}{10^{-5}}\right). \end{aligned} \tag{65}$$

Therefore we see that a baryon asymmetry of the required size is possible for a heavy baryon system with $\epsilon \sim 10^{-5}$, which requires $|M_{12}|/\Gamma_{\mathcal{B}} \sim 10^{-2}$ with $|M_{12}|/|\Gamma_{12}|$ not small.

B. Full Solution of the Boltzmann Equations

To get a more precise estimate of the baryon asymmetry, we numerically solve the system in Eqs. (39), (40), and (47). As mentioned above, we need $|M_{12}|/\Gamma_{\mathcal{B}}$ to not be much smaller than around 10^{-2}. Looking at Table I, one potential candidate is the Ω_{cb}^0 where the dominant coupling involves the operator $(dcb)^2$. In Fig. 5

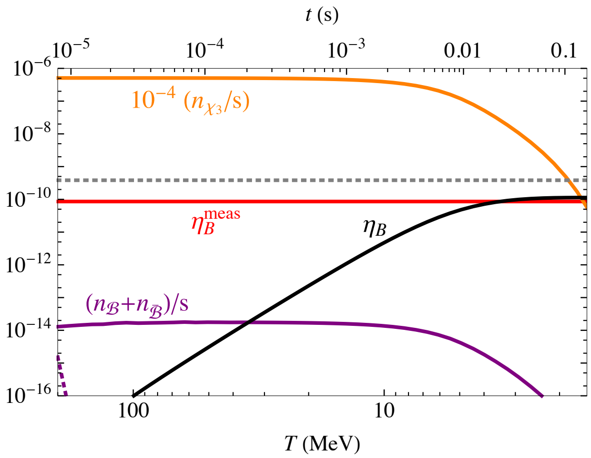

FIG. 5. $\eta_B = (n_B - n_{\bar{B}})/s$ (solid, black) as a function of the temperature or time from a numerical solution of Eqs. (39), (40), and (47) for parameters relevant to the Ω^0_{cb}-$\bar{\Omega}^0_{cb}$ system: $m_{\mathcal{B}} = 7$ GeV, $\Gamma_{\mathcal{B}} = 3\times10^{-12}$ GeV, $|M_{12}| = 3\times10^{-15}$ GeV, $|\Gamma_{12}/M_{12}| = 0.3$, $\arg(M^*_{12}\Gamma_{12}) = \pi/2$, $m_{\chi_3} = 7.5$ GeV, $\Gamma_{\chi_3} = 3\times10^{-23}$ GeV, and $\mathrm{Br}_{\chi_3\to\mathcal{B}} = 0.35$. We have taken the rate for heavy baryon scattering on the plasma from Eq. (52) and the annihilation cross section to be 400 mb. This can be compared against the value of η_B (solid, gray) from a solution of Eqs. (39), (40), and (59) as well as using the sudden decay approximation (dashed, gray) in Eq. (65). Also shown are the ratio of the number density of χ_3 to the entropy density (multiplied by 10^{-4}, solid, orange) and the ratio of the $\mathcal{B}$ plus $\bar{\mathcal{B}}$ number densities to the entropy density (solid, purple). The dashed purple line shows the equilibrium $\mathcal{B}$ and $\bar{\mathcal{B}}$ density (in units of the entropy density). The measured value of $\eta_B = 8.8\times10^{-11}$ is given by the solid red line.

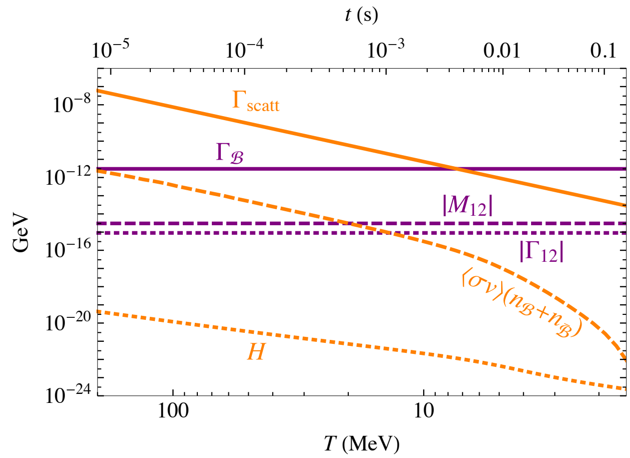

FIG. 6. The temperature dependence of the rates involved in the numerical solution of Eqs. (39), (40), and (47). The parameters are the same as in Fig. 5. In orange, from top to bottom are the scattering, annihilation, and Hubble rates. The purple lines indicate the rates relevant to the $\mathcal{B}$-$\bar{\mathcal{B}}$ system itself, $\Gamma_{\mathcal{B}}$, $|M_{12}|$, and $|\Gamma_{12}|$, from top to bottom, respectively.

we show the value of η_B as a function of temperature in the case of the asymmetry being sourced by the Ω^0_{cb}-$\bar{\Omega}^0_{cb}$ system, taking $m_{\mathcal{B}} = 7$ GeV, $\Gamma_{\mathcal{B}} = 3\times10^{-12}$ GeV [38], $|M_{12}| = 3\times10^{-15}$ GeV, $|\Gamma_{12}/M_{12}| = 0.3$, $\arg(M^*_{12}\Gamma_{12}) = \pi/2$, $m_{\chi_3} = 7.5$ GeV, $\Gamma_{\chi_3} = 3\times10^{-23}$ GeV, and $\mathrm{Br}_{\chi_3\to\mathcal{B}} = 0.35$. We have used an annihilation cross section of 400 mb (the results do not depend on whether it is flavor-blind or -sensitive) and the scattering rate given in Eq. (52).

In addition, the temperature dependence of the scattering and annihilation rates is compared to the expansion rate of the Universe as well as to the rates governing the baryon-antibaryon system in Fig. 6. As mentioned before, the (decohering) scattering is the dominant process above temperatures of about 1 MeV and, in particular, is always much larger than the annihilation rate.

At high temperatures, the heavy baryon density tracks its equilibrium value and it begins to deviate from its equilibrium value when χ_3's begin to decay. Although not directly evident from the plots (except through the change in the temperature vs. time), the out-of-equilibrium χ_3 particles actually come to dominate the energy density of the Universe prior to their decay. After the χ_3 decays, which we assume happens in less than ~ 0.1s, the Universe undergoes a transition from being matter-dominated to radiation-dominated, reheating to a temperature above a few MeV.

We have numerically confirmed the rough accuracy of the sudden decay approximation prediction for η_B over much of the parameter space. Maximal CP violation, and thus more baryon asymmetry per oscillation, occurs for $\arg(M^*_{12}\Gamma_{12}) = \pi/2$ and larger values of $|M_{12}|$ and $|\Gamma_{12}|$. A larger branching ratio, $\mathrm{Br}_{\chi_3\to\mathcal{B}}$, would produce more oscillating baryons per Majorana decay. The value of η_B that is generated is maximized if χ_3 decays when the Universe's temperature is about 10 MeV, i.e. $\tau_{\chi_3} = 1/\Gamma_{\chi_3} \sim 10^{-2}$ s. If it decays earlier than this, heavy baryon scattering on the plasma leads to decoherence, suppressing the asymmetry. If if decays later, the Universe does not have a sufficient baryon asymmetry at the time the neutrinos begin to decouple, when the Universe is around 3 MeV.

Given the constraints on the transition amplitudes in Table I, the most promising baryon that could allow for a large enough transition amplitude to source the BAU is the as yet unobserved Ω^0_{cb}. A relatively large value for $|M_{12}|$ is needed in this case, not far from the collider limit, unless $\mathrm{Br}_{\chi_3\to\Omega^0_{cb}}$ were rather large. It should be noted that the collider limits discussed in Sec. V which appear in Table I depend on the specific model that we considered. It is conceivable that the model could be extended in a way that makes the standard collider searches that we considered less constraining. For example, one could imagine making the ϕ decay to a large number of relatively soft jets by coupling to a heavy vector-like quark and a singlet which decay to a large number of colored objects. Relaxing these limits could allow for other heavy flavor baryons to source the BAU, potentially even observed baryons like the Ω^0_c, Λ^0_b, and Ξ^0_b. On the other hand, since they involve low-energy effective operators, the dinucleon decay constraints are less model dependent. Weakening them would require significant tuning of tree-level operators against those induced by weak interactions.

VIII. SUMMARY AND OUTLOOK

We have presented a model for producing the observed baryon asymmetry of the Universe which avoids high reheat temperatures. The asymmetry is generated through CP and B-violating oscillations of baryons occurring late in the hadronization era. Our model minimally introduces three neutral Majorana fermions and a single colored scalar, and could potentially be embedded into RPV SUSY.

The $\Omega_{cb} \sim (scb)$ baryon emerges as our most promising candidate when constraints due to collider data and dinucleon decay are taken into account. Note that the constraints from colliders are more model-dependent than those from the absence of dinucleon decay. Considering only the constraints from dinucleon decay, additional baryons, e.g., $\Omega_c^0 \sim (ssc)$, $\Lambda_b^0 \sim (udb)$, and $\Xi_b^0 \sim (usb)$, become viable candidates for baryogenesis via their oscillation. An interesting avenue for future work would be constructing models that are less constrained by collider experiments while preserving a large baryon oscillation rate.

Interesting signatures of this scenario could be present in the large dataset of the upcoming Belle II experiment. If the lightest Majorana fermion is sufficiently light, one possible signature would be decays of heavy flavor hadrons that violate baryon number and involve missing energy. Additionally there could be heavy flavor baryons that oscillate into their antiparticles at potentially measurable rates. Exploring the experimental prospects of this model at high luminosity, lower energy colliders in more detail will be left for future work.

Constraints from the LHC and the lack of dinucleon decay observation are quite important, suggesting the possibility of the detection of a signal in one or both areas. Dinucleon decays are a more model-independent consequence of this scenario, and because of the requirement of baryon number violation involving heavy flavors, it is likely to assume that dinucleon decay to kaons would be dominant. In the case of the LHC, a particular combination of signals in dijet resonances (singly and pair produced) along with an excess in jets plus missing energy should be expected. We should mention in this case that a long-lived neutral particle, χ_1, that decays hadronically is a generic prediction of this model. The typical χ_1 decay length is in the range of 10^{2-7} m, which could be well probed by the MATHUSLA detector that was recently proposed. The signal of a long-lived but unstable particle at this experiment could help disentangle this scenario from others that lead to excesses in jets plus missing energy.

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