PHYSICAL REVIEW D **88,** 017901 (2013)

Baryogenesis with higher dimension operators

Clifford Cheung and Koji Ishiwata

California Institute of Technology, Pasadena, California 91125, USA

(Received 5 April 2013; published 9 July 2013)

We propose a simple model of baryogenesis comprised of the standard model coupled to a singlet X via higher dimension operators $\mathcal{O}$. In the early universe, X is thermalized by $\mathcal{O}$ mediated scattering processes before it decouples relativistically and evolves into a sizable fraction of the total energy density. Eventually, X decays via $\mathcal{O}$ in an out of equilibrium, baryon number and CP violating process that releases entropy and achieves baryogenesis for a broad range of parameters. The decay can also produce a primordial abundance of dark matter. Because X may be as light as a TeV, viable regions of parameter space lie within reach of experimental probes of n-$\bar{n}$ oscillation, flavor physics, and proton decay.

DOI: 10.1103/PhysRevD.88.017901 PACS numbers: 11.30.Fs, 95.35.+d, 98.80.Cq

I. INTRODUCTION

The standard model (SM) cannot explain the observed matter-antimatter asymmetry of the Universe, and so new physics is required. In this article we propose a simple scenario for baryogenesis consisting of the SM plus an inert multiplet of states X. These states interact weakly with the SM through baryon number and CP violating higher dimension operators $\mathcal{O}$ set by the scale Λ.

The process of baryogenesis occurs in the four stages depicted in Fig. 1. In the beginning,

(i) X is thermalized with the SM plasma.
This condition is possible provided T_R, the reheating temperature, is greater than m_X, the mass of X. Hence, thermalization occurs automatically via scattering in the SM plasma mediated by $\mathcal{O}$ or the ultraviolet dynamics which generates $\mathcal{O}$. Once the Universe cools sufficiently, $\mathcal{O}$ mediated scattering goes out of equilibrium and

(ii) X decouples relativistically from the SM plasma.
Once X leaves equilibrium, it redshifts like radiation until temperatures drop below m_X, at which point X becomes nonrelativistic. Once X begins to redshift like matter,

(iii) X evolves into a large fraction of the total energy.
During this period the energy density in X is greater than that of any given relativistic species, and may even come to dominate the total energy density, sending the Universe into a matter dominated phase. The epoch of X domination terminates when

(iv) X decays, yielding a primordial baryon asymmetry.
Crucially, these out of equilibrium decays of X occur via the very same baryon number and CP violating higher dimension operators $\mathcal{O}$ that initially thermalize X in the early universe. Interference between tree and one-loop decay amplitudes generate a baryon asymmetry in the final state, as depicted in Fig. 2 for an explicit model. In certain models, X decays can also generate a primordial abundance of dark matter (DM).

Let us highlight the key features of this baryogenesis scenario. First, since this mechanism allows for low scale baryogenesis, the operators $\mathcal{O}$ may be indirectly probed through n-$\bar{n}$ oscillations, flavor violation, and proton decay. Second, this setup is quite minimal, in that only a handful of new particles X are required and the very same operators $\mathcal{O}$ that produce X initially also mediate its decay. As we will see later, a subset of the X particles can even be DM. Third, this setup exploits a cosmological "fixed point" arising because X is typically thermalized for a very broad range of reheating temperatures.

As is well known, the production and thermalization of inert particles at reheating is a ubiquitous difficulty in theories beyond the SM. This issue arises in the cosmology of gravitinos [1–8], axinos [9–12], photini [13,14], and goldstini [15,16]. Transforming this peril into a blessing is an old idea, e.g. in models linking gravitino or axino domination to baryogenesis in R-parity violating supersymmetry [17,18]. However, we argue that this mechanism applies much more broadly and is a natural byproduct of additional singlet states coupled to the SM via baryon number and CP violating higher dimension operators—the out of equilibrium condition arises from relativistic decoupling and decays of X. Alternatively, the out of

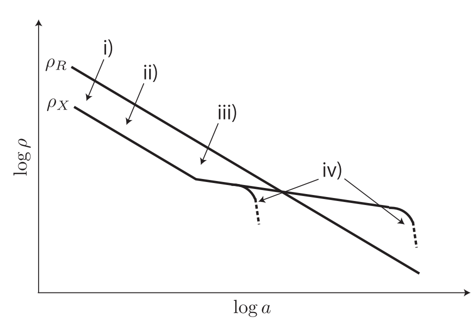

FIG. 1. The four stages of baryogenesis, shown in terms of the evolution of the energy density in SM radiation and X as a function of scale factor. The decay of X may occur before or after X grows to dominate the total energy density.

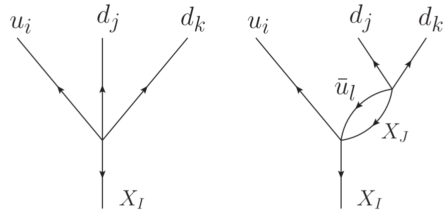

FIG. 2. Tree and one-loop diagrams for $X_I \to u_i d_j d_k$ which interfere to produce a primordial baryon asymmetry.

equilibrium condition for X can be achieved through heavy particle decays [19–23], first order phase transitions [24–27], rolling scalars [28–30], or asymmetric dark matter [31,32]. Mechanisms involving higher dimension operators have also been discussed in more specific contexts [33,34].

II. THE MODEL

In this section we present a simple theory illustrating our mechanism for baryogenesis. Consider the SM augmented by a multiplet of gauge singlet Majorana fermions X_I with mass m_I. The interaction Lagrangian for X_I is

$$\mathcal{L} = \frac{\kappa_{IJij}}{\Lambda^2}(X_I u_i)(\bar{X}_J \bar{u}_j) + \frac{\lambda_{Iijk}}{\Lambda^2}(X_I u_i)(d_j d_k) + \text{c.c.}, \quad (1)$$

where i, j, k label right-handed quark flavors. Lorentz indices are contracted implicitly among terms in parentheses, while color indices are contracted implicitly in the unique way. It is straightforward to include other Lorentz and flavor structures into the Lagrangian, but such terms will not qualitatively alter the mechanics of the model.

In terms of symmetries, baryon number is violated because X_I are Majorana, while CP is violated because the couplings κ_{IJij} and λ_{Iijk} are complex. Note that there is an exact, unbroken $\mathbb{Z}_2$ subgroup of baryon number under which X_I and the quarks are all odd.

Let us now describe the cosmological history of this model. To begin, we assume that the SM is reheated to a temperature $T_R > m_I$ shortly after inflation. If $T_R > \Lambda$, then the effective theory described in Eq. (1) does not apply, but any renormalizable ultraviolet completion of these higher dimension operators will generically induce tree level scattering processes that thermalize X. On the other hand, if $T_R < \Lambda$, then the higher dimension operator description is valid, and the interactions in Eq. (1) will mediate high energy scattering processes such as $u_i \bar{u}_j \to X_I X_J$, $u_i d_j \to X_I \bar{d}_k$, $d_j d_k \to X_I \bar{u}_i$ which also tend to thermalize X_I. The thermally averaged production cross section for X_I scales as $\langle \sigma v \rangle_I \simeq c_I T^2/\Lambda^4$, where the proportionality factor c_I depends on λ_{Iijk} and κ_{IJij}. Thus, X scattering is dominated by ultraviolet processes, and is most important at T_R. This effect is familiar from supersymmetric cosmology, where overproduction of gravitinos during reheating places a stringent limit on T_R [1]. Similar limits have been computed for a general hidden sector cosmology [35].

The critical decoupling temperature T_{D_I} defines the temperature at which these scattering processes go out of equilibrium, i.e. when $n_{\text{eq}} \langle \sigma v \rangle_I \sim H$ where n_{eq} is the equilibrium number density of X_I and H is the Hubble parameter. Together with the scaling of $\langle \sigma v \rangle_I$, this implies that $T_{D_I} \sim (\Lambda^4/m_{\text{Pl}})^{1/3}$ where $m_{\text{Pl}} \simeq 2.4 \times 10^{18}$ GeV is the reduced Planck mass. In summary, if $T_R > \Lambda$ or $\Lambda > T_R > T_{D_I}$, then X_I will be thermalized during reheating. This thermalization condition is easily satisfied for sufficiently high values of T_R, which we assume for the remainder of our discussion.

While X_I is thermalized initially, it leaves equilibrium once temperatures drop below T_{D_I}. After decoupling, the yield of X_I is given is given by

$$Y_I \simeq \frac{n_{\text{eq}}(T_{D_I})}{s(T_{D_I})}, \quad (2)$$

where s is the entropy density. The yield is constant in the absence of entropy production. Because we are interested in a case in which X_I decouples while it is relativistic, we assume throughout that $T_{D_I} > m_I$.

Once temperatures drop below m_I, X_I becomes nonrelativistic and its energy density begins to redshift like matter, since $\rho_I(T)/s(T) = m_I Y_I$ is a constant. From this point forward, the energy density of X_I will evolve to dominate that of any relativistic species. During this era, X_I may even come to dominate the total energy density, at which point the universe will enter a matter dominated phase.

This period of X_I domination ends when X_I decays via processes of the form $X_I \to u_i d_j d_k$, $\bar{u}_i \bar{d}_j \bar{d}_k$, $X_J \bar{u}_i u_j$. This final state of baryogenesis is similar to that of [23]. The associated partial decay widths are

$$\Gamma(X_I \to u_i d_j d_k) = \frac{|\lambda_{Iijk}|^2}{512\pi^3} \frac{m_I^5}{\Lambda^4}, \quad (3)$$

$$\Gamma(X_I \to X_J \bar{u}_i u_j) = \frac{|\kappa_{IJij}|^2}{1024\pi^3} \frac{m_I^5}{\Lambda^4}, \quad (4)$$

ignoring kinematic factors arising from masses of final state particles. The lifetime of X_I is constrained by number of cosmological constraints. First, if Λ is too high, the model is constrained by stringent limits from big bang nucleosynthesis (BBN) [36,37] on late time injection of electromagnetic energy. The lifetime of X_I is thus bounded by $\tau_I \lesssim 1$ s, where

$$\tau_I \simeq \frac{5.2 \times 10^{-12}\ \text{s}}{\lambda_I^2} \left(\frac{\Lambda}{10^3\ \text{TeV}}\right)^4 \left(\frac{1\ \text{TeV}}{m_I}\right)^5, \quad (5)$$

and we have defined the effective couplings $\lambda_I^2 = \sum_{ijk} |\lambda_{Iijk}|^2 + \sum_{Jij} |\kappa_{IJij}|^2/4$ where the sums range over kinematically allowed final states. Equation (5) demonstrates that BBN bounds are satisfied for a broad range of parameter space. Second, if Λ is too low, then $\langle\sigma v\rangle_I$ will be large and scattering may keep X_I in equilibrium down to temperatures of order m_I. In this case, $T_{D_I} < m_I$ and X_I decouples nonrelativistically. While a residual baryon asymmetry maybe still persist, a correct evaluation would require a full analysis of Boltzmann equations which goes beyond the scope of this work, so we restrict to the case where X_I decouples relativistically.

All but the lightest of the X_I will have CP violating decay modes of different baryon number, so their decays produce a final state baryon asymmetry through one-loop interference. The asymmetric width of $X_I \to u_i d_j d_k$ is given by interference between tree and loop diagrams depicted in Fig. 2. Ignoring kinematic factors in the final and intermediate states,

$$\Gamma(X_I \to u_i d_j d_k) - \Gamma(X_I \to \bar{u}_i \bar{d}_j \bar{d}_k) = \sum_{Jl} \frac{\mathrm{Im}(\lambda^*_{Iijk}\kappa_{IJli}\lambda_{Jljk})}{5120\pi^4}\frac{m_I^7}{\Lambda^6}, \tag{6}$$

where here the sums range over kinematically accessible final and intermediate states. We define an asymmetry parameter for each X_I decay by

$$\epsilon_I = \sum_f B(f)[\mathrm{BR}(X_I \to f) - \mathrm{BR}(X_I \to \bar{f})] = \frac{1}{20\pi}\frac{\delta_I}{\lambda_I^2}\frac{m_I^2}{\Lambda^2}, \tag{7}$$

where f sums over final states, BR denotes the branching ratio of a given process, and B is the baryon number of each final state. Here we have defined the quantity $\delta_I = \sum_{Jijkl}\mathrm{Im}(\lambda^*_{Iijk}\kappa_{IJli}\lambda_{Jljk})$ to characterize the net CP violation associated with X_I.

In order to compute the net baryon asymmetry, let us first consider the decay of a single component, X_I. The cosmology depends sensitively on the relative size of ρ_I, the energy density in X_I, and ρ_R, the total energy density in radiation, evaluated just prior to decay. If X_I decays very soon after it becomes nonrelativistic, then its energy density is of order that of a single relativistic species, which we dub the "weak domination regime," $\rho_I \ll \rho_R$. Little entropy is produced and the temperature of the radiation remains more or less constant. However, if X_I decays quite late, then it dominates the total energy density of the Universe, which we dub the "strong domination regime," $\rho_I \gg \rho_R$. Thus, X_I decays will boost the temperature of the radiation bath to an effective temperature T_I determined by the total energy density injected, $\rho_I = \pi^2 g_* T_I^4/30 = 3H^2 m_{\mathrm{Pl}}^2$ when $\tau_I \sim 1/H$, so

$$T_I \simeq \left(\frac{90}{\pi^2 g_*(T_I)}\right)^{1/4}\sqrt{m_{\mathrm{Pl}}/\tau_I}, \tag{8}$$

where g_* counts relativistic degrees of freedom.

The asymmetric baryon number generated by X_I decays is given by

$$\eta_I = \epsilon_I Y_I d_I, \tag{9}$$

where Y_I is defined in Eq. (2) and the dilution factor d_I is the ratio of the entropy density before and after X_I decays, so

$$d_I \simeq \begin{cases} 1 & \rho_I \ll \rho_R \\ \frac{3T_I}{4m_I Y_I} & \rho_I \gg \rho_R, \end{cases} \tag{10}$$

so the dilution factor is much smaller in the strong domination regime. Applying Eqs. (8)–(10) we find

$$\eta_I \simeq \frac{6.2\times 10^{-11}}{\lambda_I^2/\delta_I}\left(\frac{10^3\ \mathrm{TeV}}{\Lambda}\right)^2\left(\frac{m_I}{1\ \mathrm{TeV}}\right)^2\left(\frac{106.75}{g_*(T_{D_I})}\right), \tag{11}$$

in the weak domination regime, $\rho_I \ll \rho_R$. Meanwhile,

$$\eta_I \simeq \frac{3.5\times 10^{-9}}{\lambda_I/\delta_I}\left(\frac{10^3\ \mathrm{TeV}}{\Lambda}\right)^4\left(\frac{m_I}{1\ \mathrm{TeV}}\right)^{7/2}\left(\frac{106.75}{g_*(T_I)}\right)^{1/4}, \tag{12}$$

in the strong domination regime, $\rho_I \gg \rho_R$. Note that sphaleron processes will partially wash out the baryon asymmetry if $T_I \gtrsim m_W/\alpha_W$, where m_W is the W boson mass, $\alpha_W = g^2/4\pi$, and g is the $\mathrm{SU}(2)_L$ gauge coupling. In this case, the net baryon asymmetry is processed according to $\eta_I \to (28/79)\eta_I$.

With the expressions in Eqs. (11) and (12), it is straightforward to compute the net baryon asymmetry generated from all X_I decays. If the masses and couplings are not hierarchical, then each X_I should decay around a similar time. In this case, the X_I should either all be in the weak domination regime or all be in the strong domination regime. For the former, little entropy is produced by each X_I decay, and the net baryon asymmetry is simply given by the sum of all η_I. For the latter, entropy is substantially produced in each decay, thus diluting the asymmetry generated in earlier epochs. In this case the net baryon asymmetry is dominated by η_I from latest of the X_I decays.

Finally, let consider the issue of DM. Let us denote a stable component of the X_I multiplet by X_{DM}. If X_{DM} is lighter than the proton, then X_{DM} is exactly stable because it is the lightest odd particle under the unbroken $\mathbb{Z}_2$ subgroup of baryon number. The primordial relic abundance of X_{DM} is $\Omega_{\mathrm{DM}} = m_{\mathrm{DM}} Y_{\mathrm{DM}}^{\mathrm{tot}}(s/\rho_c)_0$, where $(\rho_c/s)_0 \simeq 3.6h^2 \times 10^{-9}$ GeV for $h \simeq 0.67$ [38]. The DM abundance arises from two sources, $Y_{\mathrm{DM}}^{\mathrm{tot}} = Y_{\mathrm{DM}}^{\mathrm{th}} + Y_{\mathrm{DM}}^{\mathrm{dec}}$ arising from thermal scattering during initial reheating and decays of heavier X_I, respectively. Assuming that baryogenesis is dominated by the decays of a single species X_I, these contributions are

$$Y_{\rm DM}^{\rm th} = \frac{\eta_I}{\epsilon_I} \frac{Y_{\rm DM}}{Y_I}, \tag{13}$$

$$Y_{\rm DM}^{\rm dec} = \frac{\eta_I}{\epsilon_I} {\rm BR}(X_I \to X_{\rm DM}), \tag{14}$$

where $Y_{\rm DM}$ and Y_I are as defined in Eq. (2). Typically, $Y_{\rm DM}$ and Y_I will be comparable, and ${\rm BR}(X_I \to X_{\rm DM}) \sim \mathcal{O}(1)$, so the contributions to the DM abundance from thermal scattering and decays will be of similar order, but both $\sim 1/\epsilon_I$ larger than the asymmetric yield.

Finally, we summarize the allowed parameter space for a single species of X_I in Fig. 3. The grey region indicates where the higher dimension operator description is invalid because $m_I > \Lambda$. For $\lambda_I = 1$, the blue region depicts the parameter space excluded by BBN limits on late decays of X_I, while the purple region depicts the parameter space excluded by requiring that X_I decouples relativistically—i.e. it is not thermalized by scattering processes at temperatures of order m_I. The purple region is very similar to the region excluded by washout from scattering processes, assuming $c_I = \lambda_I/4\pi$ for all X_I. Furthermore, taking that $\delta_I = 5$ and that X_I is the primary origin of baryogenesis, then the yellow band indicates where $10^{-11} \le \eta_I \le 10^{-10}$. Note that this choice for δ_I is not in a strong coupling regime because all couplings are really normalized to a higher dimension operator scale Λ. As noted earlier, the model also has the option of including primordial relic DM. Requiring that $\Omega_{\rm DM} h^2 \simeq 0.11$ [38] fixes the DM mass, which is denote by green dashed lines for $m_{\rm DM} = 0.1$ and 1 keV. Thus, for $\mathcal{O}(1)$ couplings, the observed baryon asymmetry is generated in the regime in which our effective theory analysis is valid, and DM can also be accommodated.

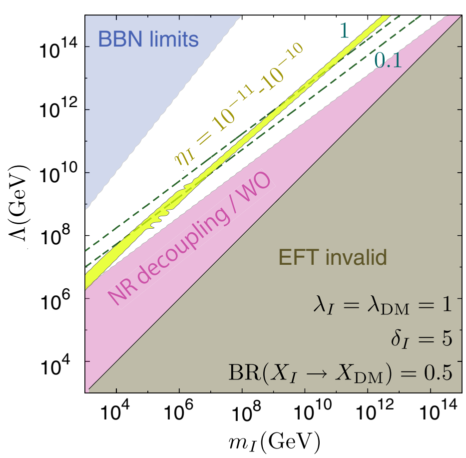

FIG. 3 (color online). Cosmologically allowed regions of parameter space. The grey region lies outside the regime of the effective theory. The blue region is disfavored by BBN, while the purple region is excluded by the requirements that X_I decouple relativistically and washout be evaded. The yellow region accommodates the observed primordial baryon asymmetry, while the green dashed lines denote the required DM mass in keV.

III. EXPERIMENTAL SIGNATURES

Our proposal offers experimentally observable consequences connected with the operators directly involved in asymmetry generation. Baryon number violating operators will typically induce highly constrained n-$\bar{n}$ oscillations via the effective operator $(ud)(dd)(ud)/M_{n\text{-}\bar{n}}^5$, where u and d are the right-handed up and down quarks, while Lorentz indices are contracted within the parentheses and color indices are contracted in the unique way. This operator is not induced at tree level in the model defined in Eq. (1), due to an accidental antisymmetry in the flavor indices of the coupling constant λ_{Iijk}. However, this operator will be induced at loop order, and more generally will be present at tree level if there are higher dimension operators in addition to those in Eq. (1). For example n-$\bar{n}$ oscillation will be induced if there are operators of the form $\lambda'_{Iijk}(X_I d_j)(d_k u_i)/\Lambda^2$. Integrating out X_I will produce the n-$\bar{n}$ operator with an effective cutoff $M_{n\text{-}\bar{n}}^5 \sim \Lambda^4 m_I/\lambda'^2_{I111}$. The characteristic time scale of n-$\bar{n}$ oscillation goes as $\tau_{n\text{-}\bar{n}} \sim M_{n\text{-}\bar{n}}^5/(3 \times 10^{-4}\ {\rm GeV}^6)$ [39,40], which together with the experimental bound, $\tau_{n\text{-}\bar{n}} \ge 2.4 \times 10^8$ s [41] implies

$$\Lambda \gtrsim 3.2 \times 10^6\ {\rm GeV} |\lambda'_{I111}|^{1/2} \left(\frac{1\ {\rm TeV}}{m_I}\right)^{1/4}, \tag{15}$$

so n-$\bar{n}$ oscillations could offer a sensitive probe of the low scale variants of this baryogenesis mechanism.

Flavor violation offers another possible probe of this model. In particular, K^0-$\bar{K}^0$ mixing is mediated by the operator $(dd)(\bar{s}\,\bar{s})/M^2_{K^0\text{-}\bar{K}^0}$, which is induced at loop level, where $M^2_{K^0\text{-}\bar{K}^0} \sim 16\pi^2 \Lambda^4/\lambda'^*_{I111}\lambda'_{I122} m_I^2$. Comparing the estimated mixing rate with the experimental bound, ${\rm Im} M_{12} \le 3.3 \times 10^{-18}$ GeV [42–45], gives

$$\Lambda \gtrsim 4.4 \times 10^4\ {\rm GeV}\, {\rm Im}(\lambda'^*_{I111}\lambda'_{I122})^{1/4} \left(\frac{m_I}{1\ {\rm TeV}}\right)^{1/2}, \tag{16}$$

which can be competitive with n-$\bar{n}$ limits.

In theories where there exists a cosmologically stable dark matter candidate $X_{\rm DM}$, there are stringent limits on proton decay via the process $p \to \pi^+ X_{\rm DM}$, whose decay rate is estimated as $\Gamma(p \to \pi^+ X_{\rm DM}) \sim \lambda_{\rm DM}^2 m_p \Lambda_{\rm QCD}^4/16\pi\Lambda^4$, where m_p is the proton mass and $\Lambda_{\rm QCD} \sim 250$ MeV is the QCD scale [46]. Then the experimental bound, $\tau_{p\to\pi^+\nu} \ge 2.5 \times 10^{31}$ yr [42], gives a very stringent limit on the cutoff

$$\Lambda \gtrsim 5.5 \times 10^{14}\ {\rm GeV} \lambda_{\rm DM}^{1/2} \left(\frac{\Lambda_{\rm QCD}}{250\ {\rm MeV}}\right). \tag{17}$$

In order to evade the proton decay bound for the DM model, we must assume a hierarchical flavor structure in the coupling $X_{\rm DM}$ to the light quarks. This can be accommodated in models of minimal flavor violation, e.g. in R-parity violating supersymmetric theories [40,47].

ACKNOWLEDGMENTS

This publication is funded by the Gordon and Betty Moore Foundation through GBMF Grant No. 776 to the Caltech Moore Center for Theoretical Cosmology and Physics. The work is also supported by the U.S. Department of Energy under Contract No. DE-FG02-92ER40701.

---

[1] T. Moroi, H. Murayama, and M. Yamaguchi, Phys. Lett. B **303**, 289 (1993).
[2] M. Y. Khlopov and A. D. Linde, Phys. Lett. **138B**, 265 (1984).
[3] M. Kawasaki and T. Moroi, Prog. Theor. Phys. **93**, 879 (1995).
[4] M. Kawasaki, K. Kohri, T. Moroi, and A. Yotsuyanagi, Phys. Rev. D **78**, 065011 (2008); K. Kohri, T. Moroi, and A. Yotsuyanagi, Phys. Rev. D **73**, 123511 (2006).
[5] M. Bolz, A. Brandenburg, and W. Buchmuller, Nucl. Phys. **B606**, 518 (2001); **B790**, 336(E) (2008).
[6] J. Pradler and F. D. Steffen, Phys. Rev. D **75**, 023509 (2007).
[7] C. Cheung, G. Elor, and L. Hall, Phys. Rev. D **84**, 115021 (2011).
[8] V. S. Rychkov and A. Strumia, Phys. Rev. D **75**, 075011 (2007).
[9] A. Brandenburg and F. D. Steffen, J. Cosmol. Astropart. Phys. 08 (2004) 008.
[10] L. Covi, H.-B. Kim, J. E. Kim, and L. Roszkowski, J. High Energy Phys. 05 (2001) 033; L. Covi and J. E. Kim, New J. Phys. **11**, 105003 (2009).
[11] A. Strumia, J. High Energy Phys. 06 (2010) 036.
[12] C. Cheung, G. Elor, and L. J. Hall, Phys. Rev. D **85**, 015008 (2012).
[13] A. Ibarra, A. Ringwald, and C. Weniger, J. Cosmol. Astropart. Phys. 01 (2009) 003.
[14] A. Arvanitaki, N. Craig, S. Dimopoulos, S. Dubovsky, and J. March-Russell, Phys. Rev. D **81**, 075018 (2010).
[15] C. Cheung, Y. Nomura, and J. Thaler, J. High Energy Phys. 03 (2010) 073.
[16] C. Cheung, J. Mardon, Y. Nomura, and J. Thaler, J. High Energy Phys. 07 (2010) 035.
[17] J. M. Cline and S. Raby, Phys. Rev. D **43**, 1781 (1991).
[18] S. Mollerach and E. Roulet, Phys. Lett. B **281**, 303 (1992).
[19] M. Fukugita and T. Yanagida, Phys. Lett. B **174**, 45 (1986).
[20] G. Lazarides and Q. Shafi, Phys. Lett. B **258**, 305 (1991).
[21] T. Asaka, K. Hamaguchi, M. Kawasaki, and T. Yanagida, Phys. Lett. B **464**, 12 (1999).
[22] K. S. Jeong and F. Takahashi, J. High Energy Phys. 04 (2013) 121.
[23] H. Davoudiasl, D. E. Morrissey, K. Sigurdson, and S. Tulin, Phys. Rev. Lett. **105**, 211304 (2010).
[24] V. A. Kuzmin, V. A. Rubakov, and M. E. Shaposhnikov, Phys. Lett. **155B**, 36 (1985).
[25] G. W. Anderson and L. J. Hall, Phys. Rev. D **45**, 2685 (1992).
[26] A. G. Cohen, D. B. Kaplan, and A. E. Nelson, Annu. Rev. Nucl. Part. Sci. **43**, 27 (1993).
[27] C. Cheung, A. Dahlen, and G. Elor, J. High Energy Phys. 09 (2012) 073.
[28] I. Affleck and M. Dine, Nucl. Phys.**B249**, 361 (1985).
[29] M. Dine, L. Randall, and S. D. Thomas, Phys. Rev. Lett. **75**, 398 (1995).
[30] M. Dine and A. Kusenko, Rev. Mod. Phys. **76**, 1 (2003).
[31] R. Kitano and I. Low, Phys. Rev. D **71**, 023510 (2005).
[32] D. E. Kaplan, M. A. Luty, and K. M. Zurek, Phys. Rev. D **79**, 115016 (2009).
[33] K. S. Babu, R. N. Mohapatra, and S. Nasri, Phys. Rev. Lett. **97**, 131301 (2006).
[34] D. J. H. Chung and T. Dent, Phys. Rev. D **66**, 023501 (2002).
[35] C. Cheung, G. Elor, L. J. Hall, and P. Kumar, J. High Energy Phys. 03 (2011) 042.
[36] M. Kawasaki, K. Kohri, and T. Moroi, Phys. Lett. B **625**, 7 (2005); Phys. Rev. D **71**, 083502 (2005).
[37] K. Jedamzik, Phys. Rev. D **74**, 103509 (2006).
[38] P. A. R. Ade *et al.* (ESA/Planck Collaboration), "Planck 2013 results. XVI. Cosmological parameters" (unpublished).
[39] J. L. Goity and M. Sher, Phys. Lett. B **346**, 69 (1995); **385**, 500(E) (1996).
[40] C. Csaki, Y. Grossman, and B. Heidenreich, Phys. Rev. D **85**, 095009 (2012).
[41] K. Abe *et al.* (Super-Kamiokande Collaboration), arXiv:1109.4227.
[42] J. Beringer *et al.* (Particle Data Group), Phys. Rev. D **86**, 010001 (2012).
[43] A. J. Buras, D. Guadagnoli, and G. Isidori, Phys. Lett. B **688**, 309 (2010); A. J. Buras and D. Guadagnoli, Phys. Rev. D **78**, 033005 (2008).
[44] J. Laiho, E. Lunghi, and R. S. Van de Water, Phys. Rev. D **81**, 034503 (2010).
[45] F. Mescia and J. Virto, Phys. Rev. D **86**, 095004 (2012).
[46] Y. Aoki, P. Boyle, P. Cooney, L. Del Debbio, R. Kenway, C. Maynard, A. Soni, and R. Tweedie (RBC-UKQCD Collaboration), Phys. Rev. D **78**, 054505 (2008).
[47] E. Nikolidakis and C. Smith, Phys. Rev. D **77**, 015021 (2008).