MIT Libraries | DSpace@MIT

MIT Open Access Articles

Generating the observed baryon asymmetry from the inflaton field

The MIT Faculty has made this article openly available. ***Please share*** how this access benefits you. Your story matters.

Citation: Hertzberg, Mark P., and Johanna Karouby. "Generating the Observed Baryon Asymmetry from the Inflaton Field." Phys. Rev. D 89, no. 6 (March 2014). © 2014 American Physical Society

As Published: http://dx.doi.org/10.1103/PhysRevD.89.063523

Publisher: American Physical Society

Persistent URL: http://hdl.handle.net/1721.1/88983

Version: Final published version: final published article, as it appeared in a journal, conference proceedings, or other formally published context

Terms of Use: Article is made available in accordance with the publisher's policy and may be subject to US copyright law. Please refer to the publisher's site for terms of use.

Massachusetts Institute of Technology

PHYSICAL REVIEW D **89,** 063523 (2014)

Generating the observed baryon asymmetry from the inflaton field

Mark P. Hertzberg[*] and Johanna Karouby[†]

Center for Theoretical Physics and Department of Physics Massachusetts Institute of Technology, Cambridge, Massachusetts 02139, USA

(Received 30 August 2013; published 20 March 2014)

We propose a mechanism by which the inflaton can generate baryogenesis, by taking the inflaton to be a complex scalar field with a weakly broken global symmetry, and present a new version of the Affleck–Dine mechanism. The smallness of the breaking is motivated both by technical naturalness and a requirement for inflation. We study inflation driven by a quadratic potential for simplicity and discuss generalizations to other potentials. We compute the inflationary dynamics and find that a conserved particle number is obtained toward the end of inflation. We then explain in detail the later decay to baryons. We present two promising embeddings in particle physics. (i) The first is using high-dimension operators for a gauge singlet; we find this leads to the observed asymmetry for decay controlled by the ∼grand unified theory scale, and this is precisely the regime where the effective field theory applies. (ii) The second is using a colored inflaton, which requires small couplings. We also point out two observational consequences: a possible large-scale dipole in the baryon density and a striking prediction of isocurvature fluctuations for which the amplitude is found to be just below current limits and potentially detectable in future data.

DOI: 10.1103/PhysRevD.89.063523 PACS numbers: 98.80.Cq, 11.30.Fs, 12.60.-i, 14.20.-c

I. INTRODUCTION

The Standard Model of particle physics is a major success of modern physics, accurately describing all laboratory observations, and recently confirmed by LHC data [1,2]. On the other hand, there are several cosmological observations that are left unaccounted for. One is the asymmetry between matter and antimatter.

Cosmological observations in the CMB, and other measurements, have revealed that the asymmetry between matter and antimatter is small. It is normally quantified by the parameter η, which is the average baryon-to-photon ratio in the Universe, with the small value

$$\eta_{\rm obs} \approx 6 \times 10^{-10}. \tag{1}$$

As discovered by Sakharov [3], any particle physics mechanism that seeks to account for this, starting from a state of symmetry, requires (i) the violation of baryon number, (ii) the breaking of C and CP symmetry, and (iii) out-of-equilibrium processes. These conditions are not effectively satisfied by the Standard Model of particle physics, strongly suggesting new physics.

One might try to avoid this problem by simply imagining that the Universe began with the asymmetry. However, such a proposal appears both unsatisfying and unlikely due to cosmological inflation, a phase of exponential expansion in the early Universe [4,5]. Evidence is continually mounting, including recent CMB data [6,7], that this paradigm explains not only the large-scale homogeneity and isotropy but also the density inhomogeneities and temperature anisotropies. An important consequence is that such an exponentially growing phase would wipe out any initial baryon number.

Ordinarily, to avoid this problem due to cosmological inflation, one introduces a collection of new fields that enter well after the inflationary phase in order to generate a baryon number. There exist many proposals for baryogenesis of this sort; for reviews, see, e.g., Refs. [8–10]. One popular mechanism is to introduce new physics around the electroweak scale and achieve baryogenesis at the electroweak phase transition; e.g., see Refs. [11–14]. As we have yet to see any new physics at the LHC, which is probing this energy regime, it is important to consider the possibility that baryogenesis is associated with much higher energies. Perhaps the only known probe to physics at very high energies is through inflation.

In this paper, and accompanying paper [15], we show that, although inflation wipes out any initial matter/antimatter asymmetry, the asymmetry can still be generated by the inflaton itself. The key reason this is possible is that the inflaton acquires a type of vacuum expectation value (vev) during inflation, and this information is not wiped out by the inflationary phase. To connect this to baryogenesis, we will put forward a new variation on the classic Affleck–Dine [16] mechanism for baryogenesis, which uses scalar field dynamics to obtain a net baryon number. In the original proposal, Affleck and Dine used a complex scalar field, usually thought to be unrelated to the inflaton but possibly a spectator field during inflation, to generate baryons in the radiation or matter eras.

Many interesting versions and constraints, often including connections to supersymmetry, have been found for

[*]mphertz@mit.edu
[†]karoubyj@mit.edu

 © 2014 American Physical Society

these types of Affleck–Dine models; some interesting works appears in Refs. [17–35]. For instance, Ref. [33] considers an interesting leptogenesis model. In Ref. [22] an interesting extradimensional model was put forward. In Ref. [24] radiative corrections were considered. In Ref. [25] connections to graviton dark matter were investigated. In Ref. [29] corrections from moduli were computed. Many works have been in the context of supergravity, such as the very interesting work in Ref. [20], which has overlap with our work here. However, with no current evidence for (low-scale) supersymmetry, it is very useful to consider less restrictive frameworks. Here we provide a general model using only the tools of effective field theory, constrained by the latest cosmological data. Our model will be minimal and predictive.

Here we propose a new model where the aforementioned complex scalar field is the inflaton itself. We study both particle physics and cosmological aspects of our model, including current observational constraints. This work is a full treatment of the basic idea summarized in our accompanying paper [15]. Our key ideas and findings are summarized as follows:

(1) We propose that the inflaton is a complex scalar field with a weakly broken global $U(1)$ symmetry. For simplicity, we consider inflation driven by a symmetric quadratic potential, plus a subdominant symmetry-breaking term.
(2) We find that a nonzero particle number is generated in the latter stage of inflation. After inflation this can decay into baryons and eventually produce a thermal Universe.
(3) We propose two promising particle physics models for both the symmetry breaking and the decay into baryons:
   (a) Using high-dimension operators, which is preferable if the inflaton is a gauge singlet.
   (b) Using low-dimension operators, which is natural if the inflaton carries color.
(4) We find that model i predicts the observed baryon asymmetry if the decay occurs through operators controlled by $\sim$ grand unified theory (GUT) scale, and this is precisely the regime where the effective field theory applies.
(5) We find that model ii requires small couplings to obtain the observed baryon asymmetry; which may be interesting in supersymmetric contexts.
(6) We find that the model allows a large-scale baryon dipole in the Universe, for which the amplitude depends on the number of e-foldings of inflation.
(7) We find that the model predicts a baryon isocurvature fluctuation at a level consistent with the latest CMB bounds and is potentially detectable in future generations of data.

The basic outline of this paper is as follows: In Sec. II we specify the model and computational techniques, in Sec. III we obtain results for the scalar field asymmetry, in Sec. IV we discuss the decay into baryons, in Sec. V we discuss constraints from inflation, in Sec. VI we discuss particle physics realizations, in Sec. VII we mention possible observational consequences, in Sec. VIII we discuss our results, and in the Appendix we present further analytical results.

II. COMPLEX SCALAR MODEL

Consider a complex scalar field ϕ, coupled to gravity, with dynamics governed by the standard two-derivative action (signature $+---$, units $\hbar = c = 1$)

$$S = \int d^4x\sqrt{-g}\left[\frac{1}{16\pi G}\mathcal{R} + |\partial\phi|^2 - V(\phi,\phi^*)\right]. \quad (2)$$

With a canonical kinetic sector, our entire freedom comes from the specification of the potential function V. For our purposes, it will be useful to decompose the potential into a "symmetric" piece V_s and a "breaking" piece V_b, with respect to a global $U(1)$ symmetry $\phi \to e^{-i\alpha}\phi$,

$$V(\phi,\phi^*) = V_s(|\phi|) + V_b(\phi,\phi^*). \quad (3)$$

To describe inflation we assume that the symmetric piece V_s dominates, even at rather large field values where inflation occurs. By expanding ϕ around 0, we assume that the symmetric piece contains a quadratic mass term

$$V_s(|\phi|) = m^2|\phi|^2, \quad (4)$$

which clearly respects a $\phi \to e^{-i\alpha}\phi$ global symmetry. It is well known that a purely quadratic potential (without higher-order corrections to V_s) will establish large field, or "chaotic" inflation [36]. This is a simple model of inflation that will provide a useful pedagogical tool to describe our mechanism for baryogenesis. As we will discuss in Sec. V, such a model is in good agreement with the spectrum of density fluctuations in the Universe, though it is in small tension with recent Planck constraints on tensor fluctuations [7]. However, generalizing to other symmetric inflationary potentials is straightforward, such as models that "flatten" at large field values.

The global symmetry is associated with a conserved particle number. So to generate a nonzero particle number (that will decay into baryons), we add a higher-dimension operator that explicitly breaks the global $U(1)$ symmetry

$$V_b(\phi,\phi^*) = \lambda(\phi^n + \phi^{*n}), \quad (5)$$

with $n \geq 3$. Note that other terms, such as $\sim(\phi^{n-1}\phi^* + \phi^{*n-1}\phi)$, would also break this $U(1)$ symmetry. Since we do not expect any qualitative differences to our central results, we restrict ourselves to the simple symmetry breaking potential given in Eq. (5). Another

motivation of focusing purely on this simple potential is that it is protected by a discrete $\mathbb{Z}_n$ symmetry of the form $\phi \to e^{2\pi i/n}\phi$. This symmetry forbids the generation of many other operators, such as $\sim(\phi^2 + \phi^{*2})$, which would cause the $U(1)$ symmetry to be broken even at low energies. Instead the $\mathbb{Z}_n$ protects the structure of the Lagrangian under renormalization group flow. Possible corrections that can be generated include $\sim\lambda^2(\phi^{2n} + \phi^{*2n})/M_{\rm Pl}^4$, which is entirely negligible in our regime of interest. In short, our model includes the set of leading operators that respect a discrete $\mathbb{Z}_n$ symmetry and weakly breaks a global continuous $U(1)$ symmetry.

We assume that the breaking parameter λ is very small (in appropriate units to be discussed in Sec. V) so that the global symmetry is only weakly broken. This assumption of very small λ is motivated by two main reasons:

(1) Since λ is responsible for the breaking of a symmetry, it is technically natural for it to be small according to the principles of effective field theory.
(2) The smallness of λ is an essential requirement on any inflationary model so that such higher-order corrections do not spoil the flatness of the potential V_s. (The breaking term adds steep positive and negative parts to the potential in the complex plane, unless λ is small.)

As should be expected, the smallness of λ favors a small baryon-to-photon ratio, as we examine in detail later.

Although the global symmetry $\phi \to e^{-i\alpha}\phi$ is explicitly broken by the higher-dimension operator $\lambda(\phi^n + \phi^{*n})$, this theory does respect the $\phi \leftrightarrow \phi^*$ symmetry, which is the charge conjugation symmetry. However, in order to satisfy the Sakharov conditions for baryogenesis, this charge conjugation (and CP) symmetry can be broken spontaneously. Affleck and Dine [16] assumed it was broken spontaneously by some light field that acquires a vev in the radiation or matter eras.

In this work, we identify ϕ with the inflaton and use the inflationary phase to obtain a vev for ϕ. As a result we are unifying inflation with baryogenesis.

A. Particle/antiparticle asymmetry

For the above class of models, we would like to compute the details of the inflationary phase, the postinflationary phase, and ultimately the transfer of energy to radiation including baryons. To begin, we note that since $n \geq 3$, then at late times the inflaton ϕ becomes small, the $\phi \to e^{-i\alpha}\phi$ symmetry violating term becomes negligible, and the symmetry becomes respected. By Noether's theorem this is associated with a conserved particle number,

$$\Delta N_\phi = N_\phi - N_{\bar\phi} = i\int d^3x\sqrt{g_s}(\phi^*\dot\phi - \dot\phi^*\phi), \quad (6)$$

where $d^3x\sqrt{g_s}$ is the spatial volume measure. N_ϕ ($N_{\bar\phi}$) is the number of ϕ particles (antiparticles), and later we will relate ΔN_ϕ to baryon number. Since ϕ is taken to be the inflaton, then we know to excellent approximation that ϕ is homogeneous on large scales. So in a Friedmann Robertson Walker (FRW) universe, the spatial integral can be immediately performed in terms of some comoving volume $V_{\rm com}$ and scale factor $a(t)$, giving

$$\Delta N_\phi = N_\phi - N_{\bar\phi} = iV_{\rm com}a^3(\phi^*\dot\phi - \dot\phi^*\phi). \quad (7)$$

To be self-consistent, we ignore spatial gradients, and the equation of motion for ϕ is

$$\ddot\phi + 3H\dot\phi + m^2\phi + \lambda n\phi^{*n-1} = 0, \quad (8)$$

where $H = \dot a/a$ is the Hubble parameter. The evolution of the field, including the early-time behavior during slow-roll inflation and the late-time behavior as the field acquires elliptic motion around the origin, is shown in Fig. 1 for two different initial conditions.

By taking a time derivative of ΔN_ϕ and using the equation of motion, it is simple to obtain an alternate expression for ΔN_ϕ at some final time t_f in terms of its value at some initial time t_i,

$$\Delta N_\phi(t_f) = \Delta N_\phi(t_i) + i\lambda V_{\rm com}n\int_{t_i}^{t_f} dt a(t)^3(\phi(t)^n - \phi^*(t)^n). \quad (9)$$

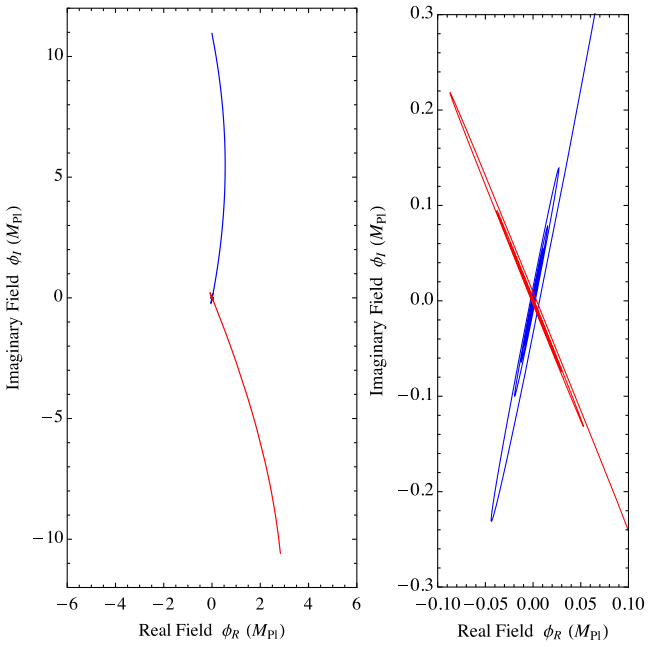

FIG. 1 (color online). Field evolution in the complex ϕ plane for $n = 3$ and $\lambda M_{\rm Pl}/m^2 = 0.006$, with initial condition $\rho_i = 2\sqrt{60}M_{\rm Pl}$. Left is zoomed out and shows early-time behavior during slow-roll inflation. Right is zoomed in to $\phi = 0$ and shows late-time elliptic motion. Blue (upper) curve is for initial angle $\theta_i = \pi/2$, and red (lower) curve is for initial angle $\theta_i = -5\pi/12$.

It is appropriate that the integral has a prefactor of λ, as the particle number must be conserved in the $\lambda \to 0$ limit.

We now rewrite the complex field in polar coordinates

$$\phi(t) = \frac{1}{\sqrt{2}}\rho(t)e^{i\theta(t)}. \tag{10}$$

Substituting into Eq. (9) leads to

$$\Delta N_\phi(t_f) = \Delta N_\phi(t_i) - \lambda \frac{V_{\rm com} n}{2^{\frac{n}{2}-1}} \int_{t_i}^{t_f} dt a(t)^3 \rho(t)^n \sin(n\theta(t)). \tag{11}$$

We will find this integral representation to be very convenient, as we explain in the next subsection.

B. Small coupling approximation

At leading order in λ, we can reduce the complexity of the problem significantly. Since the expression in Eq. (11) already has a factor of λ in front of the integral, then to first order in λ, we only need to evaluate the quantities in the integral to zeroth order in λ.

By assuming the field begins deep in the inflationary era, then any initial number of ϕ particles are diluted by inflation, and so we can ignore $\Delta N_\phi(t_i)$. Then as the field evolves at zeroth order in λ, it evolves *radially* in the complex ϕ plane. This means that we can simply take $\theta(t) = \theta_i$ as a constant and allow $\rho(t)$ to oscillate; i.e., we allow $\rho(t)$ to take the sign of the field, either positive or negative. This reduces the problem significantly to solving only a single ordinary differential equation. At first order in λ, ΔN_ϕ is simply

$$\Delta N_\phi(t_f) = -\lambda \frac{V_{\rm com} n}{2^{\frac{n}{2}-1}} \sin(n\theta_i) I(t_i, t_f), \tag{12}$$

where

$$I(t_i, t_f) \equiv \int_{t_i}^{t_f} dt a(t)^3 \rho_0(t)^n. \tag{13}$$

Here ρ_0 is a real valued function satisfying the quadratic potential version of the equation of motion

$$\ddot{\rho}_0 + 3H_0\dot{\rho}_0 + m^2\rho_0 = 0, \tag{14}$$

with corresponding Friedmann equation (we assume flat FRW)

$$H_0^2 = \frac{\varepsilon_0}{3M_{\rm Pl}^2}, \qquad \varepsilon_0 = \frac{1}{2}\dot{\rho}_0^2 + \frac{1}{2}m^2\rho_0^2, \tag{15}$$

where $M_{\rm Pl} \equiv 1/\sqrt{8\pi G}$ is the reduced Planck mass. So by solving for a single degree of freedom in a quadratic potential, we have an expression for the particle number in the small λ regime.

It is worthwhile to note that for particular values of the initial angle θ_i, such that $\theta_i = \frac{p\pi}{n} | p \in \mathbb{Z}$, no asymmetry is generated due to the $\sim \sin(n\theta_i)$ factor. In this case, the classical motion of the field is exactly radial as we can assume $\dot{\theta}_i = 0$ (nonzero $\dot{\theta}$ is quickly erased by Hubble friction). Near these special values of the initial angle, there is a large isocurvature fluctuation [see eq. (79)]. In the following, since we are interested in baryogenesis, we consider θ_i to be a typical generic value rather than these special ones.

C. Dimensionless quantities

Although ΔN_ϕ is dimensionless, it is extrinsic, depending on the size of the Universe. It is useful to define a related intrinsic quantity, namely, the particle density

$$\Delta n_\phi = \frac{\Delta N_\phi}{V_{\rm com} a^3}. \tag{16}$$

But this is now dimensionful. To obtain a dimensionless, intrinsic measure of asymmetry, we divide by the energy density ε_0:

$$A \equiv \frac{m\Delta n_\phi}{\varepsilon_0}. \tag{17}$$

This is appropriate because at late times the energy density is provided by a gas of nonrelativistic ϕ and anti-ϕ particles with energy density

$$\varepsilon_0 = m(n_\phi + n_{\bar{\phi}}). \tag{18}$$

This means that at late times, the dimensionless asymmetry variable A reduces to

$$A = \frac{\Delta n_\phi}{n_\phi + n_{\bar{\phi}}} = \frac{n_\phi - n_{\bar{\phi}}}{n_\phi + n_{\bar{\phi}}}, \tag{19}$$

which is clearly a good measure of the asymmetry.

Moreover, we would like to introduce dimensionless variables. We introduce a dimensionless time variable τ, a dimensionless field variable $\bar{\rho}$, and a dimensionless Hubble parameter $\bar{H}$ as follows:

$$\tau \equiv mt, \qquad \bar{\rho} \equiv \frac{\rho_0}{M_{\rm Pl}}, \qquad \bar{H} \equiv \frac{H_0}{m}. \tag{20}$$

In terms of these new variables, we can express A as

$$A = -f_n(\tau_i, \tau_f) \frac{\lambda M_{\rm Pl}^{n-2}}{m^2} \sin(n\theta_i), \tag{21}$$

where the dimensionless function f_n is given by

$$f_n(\tau_i,\tau_f) = \frac{n}{2^{\frac{n}{2}-1}} \frac{\bar{I}(\tau_i,\tau_f)}{a(\tau_f)^3 \bar{\varepsilon}(\tau_f)}, \tag{22}$$

where

$$\bar{I}(\tau_i,\tau_f) \equiv \int_{\tau_i}^{\tau_f} d\tau a(\tau)^3 \bar{\rho}(\tau)^n. \tag{23}$$

The dimensionless field $\bar{\rho}$ satisfies the dimensionless version of the differential equation (14). While $a(\tau), \bar{\varepsilon}$ satisfy the dimensionless version of the Friedmann equation (15). We have now extracted all the dimensionful parameters of the theory, leaving the task of solving for f_n. To do so, we perform numerics and then give an analytical estimate. From now on, we only use the dimensionless variables introduced in this section.

III. ASYMMETRY RESULTS

A. Numerical findings

Although we cannot exactly solve this nonlinear differential equation analytically in general, we can solve it numerically. (See the Appendix for exact analytical results in the Affleck–Dine regime).

In the limit in which we take τ_i very early during slow-roll inflation and we take τ_f very late after inflation (a type of ϕ-matter dominated era), then f_n becomes independent of both τ_i and τ_f. It becomes a constant only dependent on the power n. The reason is the follows: At very early times during slow-roll inflation, $a(\tau)$ is exponentially small, so the lower part of the integral is negligible. At very late times during the ϕ-matter dominated era, $\bar{\rho}$ redshifts away, so the upper part of the integral becomes negligible. Instead only the "middle" part of the integral is important, the latter stage of inflation just before the transition to matter. This is seen in Fig. 2 where we plot the integrand of $\bar{I}$ for $n = 3$ as a function of τ. So most of the ϕ (or anti-ϕ) particles are generated in the latter stage of inflation.

It is useful to note that for quadratic inflation an approximate time for the start of the matter era is given by $\tau - \tau_i \sim \sqrt{3/2}\bar{\rho}_i - 1$, as will be derived in the next subsection. Let us now compare to Fig. 2. In the figure we chose $\bar{\rho}_i = 2\sqrt{60}$ (corresponding to $N_e \approx 60$), which gives $\tau - \tau_i \approx 18$. Indeed this corresponds to the end of the sharp rise and fall of the integrand; it then begins to oscillate and redshift away in the matter era. This region is shifted to slightly earlier times in the inflationary era as we increase n, which we elaborate on in the next subsection.

Moreover, the denominator of f_n is easily shown to be independent of τ_f at late times. In summary this means that f_n approaches a constant that we denote c_n that is independent of the initial and final times in this limit:

$$c_n = f_n(\tau_i \to -\infty, \tau_f \to \infty). \tag{24}$$

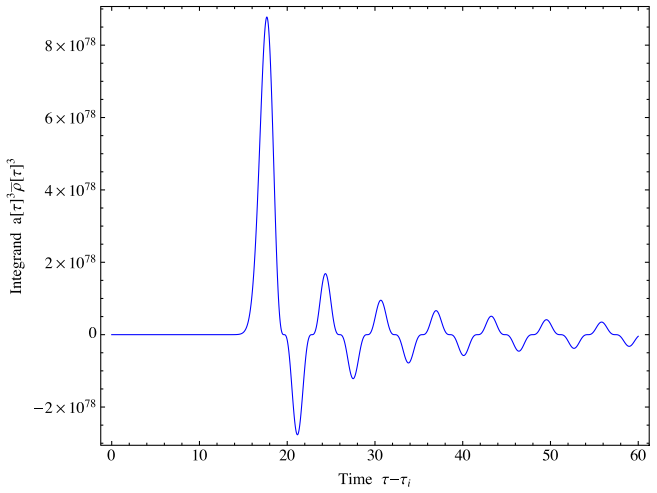

FIG. 2 (color online). The integrand that appears in the asymmetry parameter A [see Eqs. (21), (22), (23)], which is proportional to the total number of $\phi - \bar{\phi}$ particles. In this plot we have $n = 3$ and initial conditions $\bar{\rho}_i = 2\sqrt{60}$, $a_i = 1$. The large peak is in the latter phase of inflation; so this is where most of the ϕ (or anti-ϕ) particles are produced.

So our leading-order approximation for the asymmetry takes on the simple form

$$A = -c_n \frac{\lambda M_{\rm Pl}^{n-2}}{m^2} \sin(n\theta_i). \tag{25}$$

Numerically solving the dimensionless ordinary differential equation for $\bar{\rho}$ and then integrating leads to the following results for the coefficient c_n for the first few n:

$$\begin{aligned} &c_3 \approx 7.0, \quad c_4 \approx 11.5, \quad c_5 \approx 14.4, \quad c_6 \approx 21.8, \\ &c_7 \approx 34.8, \quad c_8 \approx 59.3, \quad c_9 \approx 107, \quad c_{10} \approx 201. \end{aligned} \tag{26}$$

In the next subsection, we derive an analytical expression for the coefficients c_n, valid in the large n regime; see Eq. (41). In Fig. 3 we plot the coefficients c_n for a range of $n \geq 3$, comparing the numerical results to the upcoming approximate analytical results.

We have also solved the full coupled complex system numerically; see Fig. 1 for a representative plot of the field evolution. In Fig. 4 we plot the full numerical result for A as a function of $\lambda_d \equiv \lambda M_{\rm Pl}/m^2$ for $n = 3$, and we compare it to this leading-order approximation. We have taken the field to begin during slow-roll inflation at $\rho_i = 2\sqrt{60}M_{\rm Pl}$ (corresponding to $N_e \approx 60$), which is relevant for the full numerical solution, though it is irrelevant to the leading-order approximation as we discussed earlier.

In Fig. 4 we have presented the results of two different initial angles θ_i. We note that the $\theta_i = \pi/2$ case is considerably more accurate than the $\theta_i = 5\pi/12$ case. We can understand this difference as follows: In addition

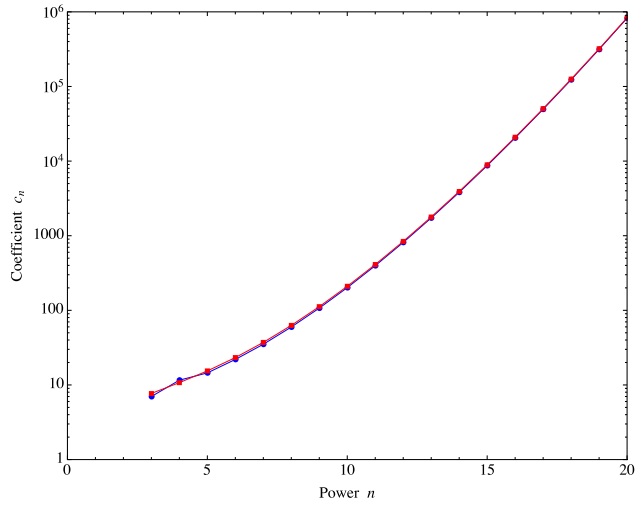

FIG. 3 (color online). A plot of the asymmetry coefficient c_n as a function of n. The blue is the exact numerical result, and the red is the approximate analytical result given in Eq. (41) and derived in Sec. III B.

to the leading-λ approximation, we expect higher-order corrections. These higher-order corrections should involve higher harmonics in $3\theta_i$. So while the first term is proportional to $\sin(3\theta_i)$, higher terms will include pieces proportional to $\sin(6\theta_i)$, etc. For special values of θ_i, such as $\theta_i = \pi/2$, this piece proportional to $\sin(6\theta_i)$ vanishes, while it does not vanish for a generic value of θ_i, such as $\theta_i = 5\pi/12$. Furthermore, higher-order terms will involve higher powers of λ; this explains the departure from the analytical result as we increase λ.

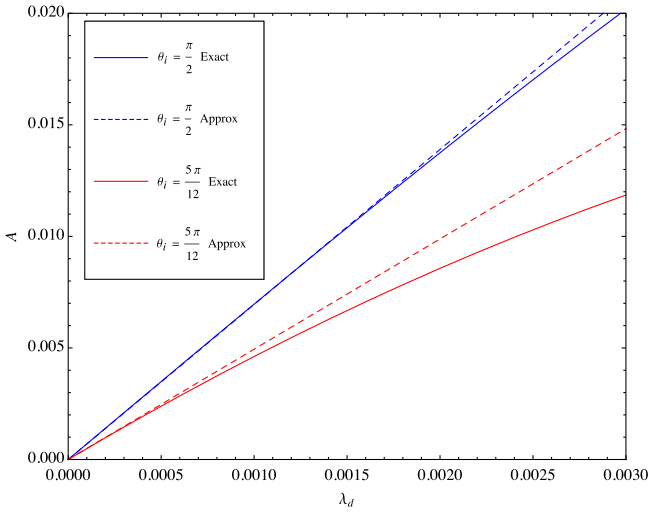

FIG. 4 (color online). A plot of A vs $\lambda_d \equiv \lambda M_{\rm Pl}/m^2$ for $n = 3$ with an initial amplitude during inflation of $\rho_i = 2\sqrt{60}M_{\rm Pl}$. The solid curves are the exact numerical results, and the dashed curves are the small λ analytical results given by Eq. (25) with $c_3 = 7.0$. The blue (upper) curves are for an initial angle of $\theta_i = \pi/2$, and the red (lower) curves are for an initial angle of $\theta_i = 5\pi/12$.

B. Analytical estimate

In the previous subsection, we reported on our numerical results for the coefficient c_n in Eq. (26) for $3 \le n \le 10$. Later, in Sec. VI A we discuss an interesting possibility that high n may be of interest. So here we would like to calculate the c_n for high n. We will find that our result is surprisingly accurate even for small n.

To compute the asymmetry, we need $a(\tau)$ and $\bar{\rho}(\tau)$, which come from solving the dimensionless versions of Eqs. (14),(15). In the slow-roll regime, the $\ddot{\bar{\rho}}$ is small in the equation of motion for $\bar{\rho}$, and the kinetic term $\dot{\bar{\rho}}^2/2$ is small in the equation for $\bar{H}$. To make this explicit, we rewrite these equations with factors of δ next to these small terms, i.e.,

$$\delta\ddot{\bar{\rho}} + 3\bar{H}\,\bar{\rho} + \bar{\rho} = 0, \tag{27}$$

$$\bar{H}^2 = \frac{1}{3}\left(\frac{\delta}{2}\frac{(d\bar{\rho}^2}{d\tau)} + \frac{1}{2}\bar{\rho}^2\right). \tag{28}$$

This allows for a consistent power series expansion as follows:

$$\bar{\rho} = \bar{\rho}_a + \delta\bar{\rho}_b + \cdots \tag{29}$$

$$\bar{H} = \bar{H}_a + \delta\bar{H}_b + \cdots \tag{30}$$

In the final result, we will replace $\delta \to 1$. For an approximation to $\bar{I}$, we only need the first few terms in the expansion. After a straightforward calculation, we find the solution

$$\bar{\rho}_a(\tau) = \bar{\rho}_i - \sqrt{\frac{2}{3}}(\tau - \tau_i), \tag{31}$$

$$\bar{\rho}_b(\tau) = \frac{1}{3}\left(\frac{1}{\bar{\rho}_a(\tau)} - \frac{1}{\bar{\rho}_i}\right), \tag{32}$$

$$\bar{H}_a(\tau) = \frac{\bar{\rho}_a(\tau)}{\sqrt{6}}, \tag{33}$$

$$\bar{H}_b(\tau) = \frac{1}{3\sqrt{6}}\left(\frac{2}{\bar{\rho}_a(\tau)} - \frac{1}{\bar{\rho}_i}\right). \tag{34}$$

Recall that the integrand of $\bar{I}$ is $\sim\bar{a}(\tau)^3\bar{\rho}(\tau)^n$. Since it is only a power law in $\bar{\rho}$, the first term $\bar{\rho}_a$ will suffice, while for $\bar{a}$ the first and second terms in $\bar{H}$ are required, as the scale factor is an exponential of this, so higher accuracy is required. After integrating the Hubble parameter $\bar{H}_a + \bar{H}_b$, we find the following expression for the scale factor in the slow-roll regime, including a subleading correction:

$$a(\tau) \approx \frac{a_i \exp\left(\frac{\tau-\tau_i}{\sqrt{6}}\left(\bar{\rho}_i - \frac{1}{3\bar{\rho}_i}\right) - \frac{(\tau-\tau_i)^2}{6}\right)}{\left(1 - \sqrt{\frac{2}{3}}\frac{\tau-\tau_i}{\bar{\rho}_i}\right)^{1/3}}. \tag{35}$$

For large n the integral $\bar{I}(\tau_i, \tau_f)$ [Eq. (39)] is dominated by the late stage of inflation rather than the matter era. For a concrete approximation, let us use the slow-roll regime from the start of inflation until the slow-roll parameter reaches $\epsilon = 3$, which corresponds to the equation of state $w = 0$, the start of the matter era. This corresponds to using an end point of integration of

$$\tau_f = \tau_i + \sqrt{\frac{3}{2}}\bar{\rho}_i - 1 \tag{36}$$

and ignoring the contribution from the matter era as it is relatively small for large n.

We now perform a change of variable from τ to b, where

$$b \equiv \sqrt{\frac{3}{2}}\bar{\rho}_i - (\tau - \tau_i). \tag{37}$$

The corresponding approximate value of the integral becomes

$$\bar{I} \approx a_i^3 \left(\frac{2}{3}\right)^{(n-1)/2} \frac{\bar{\rho}_i e^{3\bar{\rho}_i^2/4}}{\sqrt{e}} \times \int_1^{\sqrt{\frac{3}{2}}\bar{\rho}_i} dbb^{n-1} \exp\left(-\frac{b^2}{2} + \frac{b}{\sqrt{6}\bar{\rho}_i}\right). \tag{38}$$

Since we are assuming the Universe began deep in the slow-roll regime, we can take $\bar{\rho}_i$ large and ignore the $b/\bar{\rho}_i$ term inside the exponential and send the upper limit of the integral to ∞. We are then able to rewrite the integral as

$$\bar{I} \approx a_i^3 2^{n-3/2} 3^{-(n-1)/2} \frac{\bar{\rho}_i e^{3\bar{\rho}_i^2/4}}{\sqrt{e}} \Gamma_{\frac{1}{2}}(n/2), \tag{39}$$

where Γ_a is the incomplete gamma function, defined as

$$\Gamma_a(z) \equiv \int_a^\infty dt t^{z-1} e^{-t}. \tag{40}$$

We now form the coefficient c_n defined by Eqs. (22), (24). To do so, we need to divide $\bar{I}$ by $a(\tau)^3\bar{\varepsilon}(\tau)$. This cancels the factor of $a_i^3\bar{\rho}_i \exp(3\bar{\rho}_i^2/4)$ in Eq. (39), as this is just an extensive factor and independent of n. The final result for c_n is then

$$c_n \approx \tilde{c} 2^{n/2} 3^{-n/2} n \Gamma_{\frac{1}{2}}(n/2), \tag{41}$$

where $\tilde{c}$ absorbs any $\mathcal{O}(1)$ factors that are independent of n; it accounts for the transition of $\bar{\rho}$, a from inflation to matter eras, and we find its value to be $\tilde{c} \approx 6.64$. This result for c_n is asymptotically correct for large values of n and also provides a good estimate for moderate values of n; see Fig. 3 for comparison to the exact numerical results.

IV. DECAY INTO BARYONS

We would now like to compute the relationship between this asymmetry parameter A defined in Eq. (19) and the final productions of baryons. Recall that the baryon asymmetry is defined as the ratio of baryon difference to photon number at late times,

$$\eta = \frac{(n_b - n_{\bar{b}})_f}{(n_\gamma)_f} = \frac{(N_b - N_{\bar{b}})_f}{(N_\gamma)_f}, \tag{42}$$

where f indicates the late-time, or "final," value after decay and thermalization. We would like to relate this to the ϕ asymmetry parameter A and the details of how ϕ decays.

We associate with each ϕ particle a baryon number b_ϕ; for instance, $b_\phi = 1$ or $b_\phi = 1/3$ in simple models (see Sec. VI). We assume that the decay of ϕ and all subsequent interactions is baryon number conserving, so we can relate the final number to the initial number as

$$(N_b - N_{\bar{b}})_f = b_\phi (N_\phi - N_{\bar{\phi}})_i, \tag{43}$$

where i indicates the early-time, or "initial," value before decay and thermalization (but well after the baryon-violating processes have stopped). We can now rewrite η as

$$\eta = b_\phi \frac{(N_\phi - N_{\bar{\phi}})_i}{(N_\gamma)_f} = b_\phi A \frac{(N_\phi + N_{\bar{\phi}})_i}{(N_\gamma)_f}. \tag{44}$$

So we need to evaluate the number of ϕ particles at early times and the number of photons at late times.

At early times we can relate the number of ϕ particles to the energy density as

$$(N_\phi + N_{\bar{\phi}})_i = V_{\mathrm{com}}(a^3\varepsilon)_i/m, \tag{45}$$

and the energy density is related to the Hubble parameter as

$$(\varepsilon)_i = 3M_{\mathrm{Pl}}^2 (H^2)_i. \tag{46}$$

On the other hand, at late times we can relate the number of photons to the number density as

$$(N_\gamma)_f = V_{\mathrm{com}}(a^3 n_\gamma)_f, \tag{47}$$

and the number density is related to the temperature as

$$(n_\gamma)_f = \frac{2\zeta(3)}{\pi^2}(T^3)_f. \tag{48}$$

Hence, we express η as

$$\eta = \frac{3\pi^2 b_\phi}{2\zeta(3)} \frac{M_{\rm Pl}^2 A}{m} \frac{(a^3H^2)_i}{(a^3T^3)_f}. \quad (49)$$

Note that both the numerator and denominator here are separately time independent. It is nontrivial to exactly compute these final parameters as a function of initial conditions as it depends on the details of the decay and thermalization. However, a good approximation arises by assuming that the thermalization is rapid. This means that we can simply evaluate both the initial and final quantities around the time of decay.

We denote the decay rate of the ϕ field as Γ_ϕ. Thermalization occurs around $H \approx \Gamma_\phi$ [37]. The energy density in the radiation era is

$$\varepsilon = \frac{\pi^2}{30} g_* T^4, \quad (50)$$

where g_* is the number of relativistic degrees of freedom, which is typically $g_* \sim 10^2$. Then by setting $\varepsilon = 3H^2 M_{\rm Pl}^2 \approx 3\Gamma_\phi^2 M_{\rm Pl}^2$ and solving for T, we have an estimate of the reheat temperature:

$$T_r \approx \left(\frac{90}{g_* \pi^2}\right)^{1/4} \Gamma_\phi^{1/2} M_{\rm Pl}^{1/2}. \quad (51)$$

Substituting T with the reheat temperature T_r and $H \approx \Gamma_\phi$ into Eq. (49), we obtain

$$\eta \approx \frac{\beta \pi^{7/2} g_*^{3/4} b_\phi}{2^{7/4} 3^{1/2} 5^{3/4}} \frac{A \Gamma_\phi^{1/2} M_{\rm Pl}^{1/2}}{m}, \quad (52)$$

where β is an $\mathcal{O}(1)$ fudge factor that accounts for the details of the transition from the ϕ era to the thermal era. The precise value of β is not important for our qualitative conclusions, but we do expect it to be of order 1.

Finally, we insert the expression for A from Eq. (25) to obtain our result for the baryon-to-photon ratio:

$$\eta \approx -c_n \frac{\beta \pi^{7/2} g_*^{3/4} b_\phi}{2^{7/4} 3^{1/2} 5^{3/4}} \frac{\lambda \Gamma_\phi^{1/2} M_{\rm Pl}^{n-3/2}}{m^3} \sin(n\theta_i). \quad (53)$$

For this η to match with the observed value $\eta_{\rm obs} \approx 6 \times 10^{-10}$, we require $\lambda \Gamma_\phi^{1/2} M_{\rm Pl}^{n-3/2}/m^3$ to take on a particular value, namely,

$$\left.\frac{\lambda \Gamma_\phi^{1/2} M_{\rm Pl}^{n-3/2}}{m^3}\right|_{\rm req} \approx 7 \times 10^{-11} c_n^{-1} \times \left(\frac{g_*^{3/4}}{30} \beta b_\phi |\sin(n\theta_i)|\right)^{-1}, \quad (54)$$

where the subscript "req" indicates that this is the "required" value for agreement with η_{obs}. Inside the parenthesis is a term that should be $\mathcal{O}(1)$, since we expect

$$|\sin(n\theta_i)| \sim 1, b_\phi \sim 1, \beta \sim 1, g_* \sim 10^2 \quad (55)$$

(provided θ_i is not near a special value, as mentioned earlier). As a result, the prefactor of $\sim 7 \times 10^{-11} c_n^{-1}$ is most important here.

V. CONSTRAINTS FROM INFLATION

A. Quadratic inflation

During inflation we assume that the potential is dominated by the symmetric $m^2\phi^*\phi$ term, and so to first approximation the motion is radial. We can thus keep the phase of our complex field fixed, $\phi = \rho e^{i\theta_i}/\sqrt{2}$, and write the potential and the so-called slow-roll parameters as

$$V = \frac{1}{2} m^2 \rho^2, \quad (56)$$

$$\epsilon_{sr} \equiv \frac{M_{\rm Pl}^2}{2}\left(\frac{V'}{V}\right)^2 = \frac{2M_{\rm Pl}^2}{\rho^2}, \quad (57)$$

$$\eta_{sr} \equiv M_{\rm Pl}^2 \frac{V''}{V} = \frac{2M_{\rm Pl}^2}{\rho^2}. \quad (58)$$

Also, the number of e-foldings is given by

$$N_e = \frac{1}{M_{\rm Pl}} \int_{\rho_{\rm end}}^{\rho_i} \frac{d\rho}{\sqrt{2\varepsilon}} = \frac{\rho_i^2 - \rho_{\rm end}^2}{4M_{\rm Pl}^2}. \quad (59)$$

Hence, the modes that describe the Universe on large scales are emitted when the field value is

$$\rho_i \approx 2\sqrt{N_e} M_{\rm Pl} \quad (60)$$

and inflation ends for $\rho_{\rm end} \sim M_{\rm Pl}$ (when $\epsilon_{sr} \sim 1$).

For any simple single-field slow-roll model of inflation, the squared amplitude of density fluctuations is predicted to be

$$\Delta_R^2 = \frac{V_i}{24\pi^2 M_{\rm Pl}^4 \epsilon_{sr}}, \quad (61)$$

where V_i is the potential energy when modes leave the horizon. The observed value is $\Delta_{R,\rm obs}^2 \approx 2.45 \times 10^{-9}$ from WMAP and Planck data [6,7]. For the simple quadratic model, we have

$$\Delta_R^2 \approx \frac{N_e^2 m^2}{6\pi^2 M_{\rm Pl}^2}. \quad (62)$$

Hence, in order to have the correct amplitude of density fluctuations, the mass of the field must be

$$m \approx \frac{\sqrt{6}\pi\Delta_{R,\mathrm{obs}}M_{\mathrm{Pl}}}{N_e} \approx 1.5\times 10^{13}\ \mathrm{GeV}\left(\frac{60}{N_e}\right). \quad (63)$$

We have indicated that $N_e = 60$ is a good reference value, but its precise value can be smaller if the reheat temperature is low. The precise relationship between the number of e-foldings of inflation and the scale of interest k is given by [38]

$$N_e = 62 - \ln\frac{k}{aH} + \frac{1}{4}\ln\frac{V_i}{(10^{16}\ \mathrm{GeV})^4} + \frac{1}{4}\ln\frac{V_i}{V_{\mathrm{end}}} - \frac{1}{12}\ln\frac{V_{\mathrm{end}}}{\varepsilon_{\mathrm{reh}}}, \quad (64)$$

where V_{end} is the potential energy at the end of inflation and $\varepsilon_{\mathrm{reh}}$ is the energy density at the start of reheating. For typical high-scale inflation models, such as quadratic inflation, $V_i \sim (10^{16}\ \mathrm{GeV})^4$. For efficient reheating this leads to $N_e \sim 60$, while for inefficient reheating, say $H_{\mathrm{reh}} \sim 1$ eV, we have $N_e \sim 50$, for a typical CMB scale k. For $N_e \sim 50$–60 the spectral index $n_s = 1 - 6\varepsilon_{sr} + 2\eta_{sr} \approx 1 - 2/N_e$ is in excellent agreement with observation $n_s \approx 0.96$ [6,7].

On the other hand, a small problem with these quadratic models is the tensor-to-scalar ratio, which is $r = 16\varepsilon_{sr} \approx 8/N_e$ and is somewhat large and just outside the 2 sigma region of recent Planck CMB constraints [7]. However, it is important to note that the details of the inflationary potential are not crucial to our baryogenesis mechanism. We can easily construct other potentials which flatten at large field values, decreasing the tensor-to-scalar ratio, and still provide the same qualitative mechanism for baryogenesis as outlined in this work. For the present purposes, it is enough to use the above mass scale m and number of e-foldings N_e as characteristic values (with m having the greatest variability between different inflationary models).

B. Implications for baryogenesis

An important constraint is that the symmetry-breaking term in the potential $\lambda(\phi^n + \phi^{*n})$ be subdominant during inflation. Since this contribution to the potential goes negative at large field values, we obviously need it to be small during inflation. Writing ϕ in terms of polar variables, the constraint is

$$\frac{\lambda}{2^{n/2-1}}\rho_i^n\cos(n\theta_i) \ll \frac{1}{2}m^2\rho_i^2. \quad (65)$$

Then using $\rho_i = 2\sqrt{N_e}M_{\mathrm{Pl}}$, and demanding this to be true for all θ_i, gives the following upper bound on λ:

$$\lambda \ll \lambda_0 \equiv \frac{m^2}{2^{n/2}N_e^{n/2-1}M_{\mathrm{Pl}}^{n-2}}. \quad (66)$$

For typical values of mass and duration of inflation (e.g., $m \sim 10^{13}$ GeV and $N_e \sim 60$), this provides an important bound on λ for each n.

We now use the threshold value λ_0 to rewrite the condition for the correct baryon-to-photon ratio (54) as a condition on the decay rate:

$$\Gamma_{\phi,\mathrm{req}} \approx 10^{-7}\ \mathrm{eV}\times 2^{n+1}N_e^{n-2}c_n^{-2}\left(\frac{\lambda_0}{\lambda}\right) \times\left(\frac{m}{10^{13}\ \mathrm{GeV}}\right)^2\left(\frac{g_*^{3/4}}{30}\beta b_\phi|\sin(n\theta_i)|\right)^{-2}. \quad (67)$$

To provide concrete quantitative results for the required decay rate, we need to choose some characteristic values for parameters. For instance, if we assume that the coupling λ is a factor of 10 smaller than its inflationary upper bound λ_0, and take $\beta b_\phi|\sin(n\theta_i)| \approx 1$, $m \approx 1.5\times 10^{13}$ GeV, $N_e \approx 55$, $g_* \approx 10^2$, and we insert the c_n from Eq. (26) [or Eq. (41)], then for each of the different n (from 3 to 10), we find that the required decay rate is

$$\begin{aligned} n = 3 &\Rightarrow \Gamma_{\phi,\mathrm{req}} \approx 4\times 10^{-5}\ \mathrm{eV},\\ n = 4 &\Rightarrow \Gamma_{\phi,\mathrm{req}} \approx 2\times 10^{-3}\ \mathrm{eV},\\ n = 5 &\Rightarrow \Gamma_{\phi,\mathrm{req}} \approx 10^{-1}\ \mathrm{eV},\\ n = 6 &\Rightarrow \Gamma_{\phi,\mathrm{req}} \approx 6\ \mathrm{eV},\\ n = 7 &\Rightarrow \Gamma_{\phi,\mathrm{req}} \approx 2\times 10^{2}\ \mathrm{eV},\\ n = 8 &\Rightarrow \Gamma_{\phi,\mathrm{req}} \approx 9\times 10^{3}\ \mathrm{eV}\\ n = 9 &\Rightarrow \Gamma_{\phi,\mathrm{req}} \approx 3\times 10^{5}\ \mathrm{eV},\\ n = 10 &\Rightarrow \Gamma_{\phi,\mathrm{req}} \approx 10^{7}\ \mathrm{eV}. \end{aligned} \quad (68)$$

It is important to compute the associated reheat temperature using Eq. (51). For $n = 3$, corresponding to the lowest value of $\Gamma_{\phi,\mathrm{req}}$, the reheat temperature is still substantial: $T_r \sim 10^2$ GeV. For higher n, $\Gamma_{\phi,\mathrm{req}}$ increases, so too does T_r; e.g., for $n = 10$, $T_r \sim 10^8$ GeV—all much higher than ~MeV, the characteristic temperature of the Universe at the onset of big bang nucleosynthesis.

Having established a consistent cosmology of inflation, reheating, and big bang nucleosynthesis, we would now like to go further and explore if such decay rates are plausible in various particle physics frameworks.

VI. PARTICLE PHYSICS MODELS

In the following, we introduce two different types of particle physics models. First, we consider ϕ to be a gauge singlet, and, second, we consider ϕ to carry color. For each model we estimate the decay rates and compare to the required decay rates $\Gamma_{\phi,\mathrm{req}}$ to obtain η_{obs}.

Our upcoming estimates of the decay rates will involve a perturbative analysis. At the end of inflation when the field value is large, a nonperturbative regime involving

parametric resonance can sometimes take place. However, its existence and/or effectiveness is highly model dependent. In particular there can be self-resonance and/or resonance into other fields. The possibility of self-resonance could arise from the nonlinear $\sim\lambda(\phi^n+\phi^{*n})$ term in the action. However, it can be shown that since this term is assumed to be subdominant during inflation, it is even more subdominant after inflation, and the corresponding resonance is entirely inefficient and negligible. So this justifies our perturbative treatment while ignoring nonperturbative self-resonance. The possibility of resonance into other fields is interesting and is possible future work that we discuss later in the discussion section. In this case its existence depends on the value of the couplings to other fields, such as the Higgs. It is radiatively stable and self-consistent to assume these couplings are sufficiently small that this too is ignorable. We assume, for simplicity, that this is the case in the present work.

A. High-dimension operators

In the simplest case, one can take ϕ to be a gauge singlet. In this case there are many operators which could couple ϕ to Standard Model degrees of freedom. This is problematic because there is no argument based on degree of freedom counting (namely, gauge redundancy) as to why the particle number contained in ϕ should only decay into quarks, etc., generating baryons. For instance, one could include operators, such as $\sim\phi H^\dagger H$ or $\phi\bar{f}f$, which would violate baryon-number conservation, and ϕ would just decay to Higgs or fermions, etc.

However, a natural way around this problem of ϕ decaying into nonbaryonic particles is to suppose that the global $U(1)$ symmetry is almost an exact symmetry of nature (or at least in the ϕ sector). This assumption is consistent with the requirement that the $U(1)$ breaking is small for a consistent model of inflation, which is technically natural, although the symmetry is not expected to be exact.

The reason for this is the following. We begin by imaging we have some $U(1)$ global symmetry. However, we know that global symmetries cannot be exact in quantum gravity (one way to see this is from the "black-hole-no-hair theorem" which forbids a black hole to carry any baryon number, and therefore the $U(1)$ baryon symmetry cannot be exact). So a complete theory should allow for at least a weak breaking of the symmetry. Symmetries are often broken by some high-dimension operator, representing the breaking due to some microscopic, perhaps Planckian, physics. Another way to argue this is to just impose a $\mathbb{Z}_n$ symmetry and assume n is greater than 4, which is a reasonable model building assumption. For high n, the breaking parameter will need to satisfy $\lambda\lesssim(\text{few}/\sqrt{G})^{4-n}$ to be consistent with inflation. This is compatible with quantum gravity expectations.

This suggests the intriguing possibility of using a high value of n in the breaking term. For instance, we can imagine that the $U(1)$ symmetry breaking occurs at dimension $n\geq 8$ operators. If this is the case, then all low-dimension operators, such as $\sim\phi H^\dagger H$, which break the $U(1)$ or discrete $\mathbb{Z}_n$ symmetries, would be forbidden.

For definiteness, let us take $n=8$ and suppose that the $U(1)$ symmetry is intact for all dimensions 7 or lower (which is protected against radiative corrections by a $\mathbb{Z}_8$ symmetry). Since ϕ carries baryon number, then up to dimension 7, it could only decay into quarks. It is well known that at dimension 6 the Standard Model degrees of freedom allow for gauge singlet operators carrying baryon number of the form $qqql$. In fact there are five types of such operators [39], but the full details do not concern us here. To construct a $U(1)$ invariant operator, we multiply the $qqql$ operator by ϕ (ϕ^*) and build the $U(1)$ symmetric dimension-7 operator

$$\Delta\mathcal{L}\sim\frac{c}{\Lambda^3}\phi^* qqql+\text{H.c.}, \tag{69}$$

where we are suppressing indices for brevity. Here we have introduced an energy scale Λ that sets the scale of new physics (and the cutoff on the field theory), and c is some dimensionless coupling. In this case we have $b_\phi=1$, as this operator causes ϕ to decay into three quarks (and a lepton).

The decay rate associated with this operator is roughly

$$\Gamma_\phi(\phi\to q+q+q+l)\sim\frac{c^2}{8\pi}\frac{m^7}{\Lambda^6}. \tag{70}$$

We now compare this to the required decay rates from Eq. (68). For the case of $n=8$, we would need this to be $\Gamma_\phi\sim 10^4$ eV. For $m\approx 1.5\times 10^{13}$ GeV and $c=\mathcal{O}(1)$, we find that the model has the required decay rate for

$$\Lambda\sim 10^{16}\text{ GeV}, \tag{71}$$

which is intriguingly around the GUT scale. Similarly, if we push the breaking off to even higher operators, such as $n=10$ or $n=12$, then the model requires slightly smaller Λ, approaching $\Lambda\sim 10^{15}$ GeV, which is still close to the GUT scale.

Since we are using high-dimension operators, we are studying only an effective field theory. As such, we need to check that the parameters are within the regime of validity of the theory. There are two basic constraints: The first is that the physics that is being integrated out is below the scale of quantum gravity E_{QG}; i.e., it would be rather unrealistic if we were appealing to physics beyond the quantum gravity scale to UV complete the theory. The second is that the inflationary regime, which probes scales of order Hubble, is below the cutoff of the field theory. Together these two constraints imply

$$H_i \ll \Lambda \ll E_{QG}. \tag{72}$$

Since $H_i \sim 10^{13,14}$ GeV for high-scale inflation models, and $E_{QG} \sim 10^{18,19}$ GeV, then we require values for Λ around $\Lambda \sim 10^{15,16}$ GeV for the effective field theory to make sense. Remarkably, this is precisely the value found above to obtain $\eta_{\rm obs}$.

To summarize, if the visible sector has an almost exact $U(1)$ symmetry that is broken only by high-dimension operators ($n \geq 8$), then the leading operator, allowed by symmetry, that mediates ϕ decay is dimension 7, and the decay products are quarks and leptons. After imposing constraints from inflation on the breaking parameter and by taking the dimension-7 operator to be controlled by the GUT scale, which is within the regime of validity of the effective field theory, we can obtain the observed matter/antimatter asymmetry η_{obs}.

B. Colored inflaton

Another possibility is to allow the inflaton to carry color and to transform under the fundamental representation of $SU(3)_c$. We give ϕ a color index, $i = r, w, b$, and allow for "up" ϕ_u and "down" ϕ_d versions and different generations labelled by g. In addition to the usual symmetric mass term, etc., we can construct $U(1)$-violating terms in the potential that respect the $SU(3)_c$ symmetry. For instance, at dimension $n = 3$, we can introduce the breaking term

$$V_b(\phi, \phi^*) = \lambda^{gg'g''} \varepsilon_{ii'i''} \phi^i_{ug} \phi^{i'}_{dg'} \phi^{i''}_{dg''} + \text{H.c.}, \tag{73}$$

where ε_{ijk} is the totally antisymmetric tensor, and we have summed over color indices and different generations. We need one up and two down for a gauge singlet under electric charge. We also need at least two different generations, or otherwise the antisymmetric tensor will cause the term to vanish. For simplicity we have written down the leading $U(1)$-violating operator, dimension $n = 3$. As in the previous subsection, we could imagine that the symmetry is only broken by some higher-dimension operators; such a generalization is straightforward.

To compute the generated particle number, we should repeat the analysis from Secs. II and III, but now with various indices to track. We expect the basic qualitative conclusions from those sections to be the same, so we will not go through those details here. Moreover, corrections from gluons should be small at these high energies due to asymptotic freedom [40,41].

Since ϕ carries color, we can readily build operators that mediate ϕ decay into quarks, while respecting the global symmetry, such as the following dimension-4 operator

$$\Delta\mathcal{L} \sim y \phi^{i*} q^i \bar{f} + \text{H.c.}, \tag{74}$$

where f is some color neutral fermion and y is a type of Yukawa coupling. In this case we have $b_\phi = 1/3$, as this operator causes ϕ to decay into one quark (and another fermion). This decay rate is roughly

$$\Gamma_\phi(\phi \to q + \bar{f}) \sim \frac{y^2}{8\pi} m. \tag{75}$$

For high-scale inflation, such as quadratic inflation that we discussed earlier, the inflaton mass is large $m \sim 10^{13}$ GeV, so one would require an extremely small value of y to obtain decay rates comparable to the required values we computed earlier in Eq. (68) (of course that was only for the gauge singlet case, but we expect similar values within an order of magnitude). For example, in order to have $\Gamma_\phi \sim$ eV, we would need $y \sim 10^{-10}$. In generic models of particle physics, such a small dimensionless coupling would not normally be radiatively stable. In certain settings, such as supersymmetry (which would provide extra motivation for the existence of such colored scalars, or "squarks"), one could examine if some nonrenormalization theorem may help to stabilize y at such small values.

Perhaps a more promising possibility would be to push the $U(1)$ symmetry breaking parameter to values much smaller than is required by the inflationary constraint. This is, as mentioned earlier, technically natural. Recall from Eq. (67) that the required decay rate to obtain the correct value of η is inversely proportional to λ. So by making λ extremely small, we can raise the values in Eq. (68) to much higher values (recall those values were based on $\lambda \sim \lambda_0/10$). A consistent model would then only require a moderately small value for y.

VII. OBSERVATIONAL CONSEQUENCES

A. Large-scale dipole

As in the original Affleck–Dine model, an interesting feature of this mechanism is that the dynamics respect C (and CP) symmetry, but it is broken spontaneously by the state of the Universe. In the present work, this is due to the initial value of the inflaton field in the complex plane. It is often thought that any initial conditions are wiped out by inflation, but the vev of the inflaton is evidently not. Instead this initial condition has a very important consequence for the final matter/antimatter in the Universe; it determines if the imbalance favors matter or antimatter. Recall from Eq. (53) that η depends on the initial angle of the inflaton:

$$\eta \propto -\sin(n\theta_i). \tag{76}$$

Since we expect the initial angle to be randomly distributed in the domain $\theta_i \in [0, 2\pi]$, then we expect η to be a random variable, with 50% chance positive (matter domination; see the blue curve in Fig. 1) and 50% chance negative (antimatter domination; see the red curve in Fig. 1).

If inflation only lasted for the minimum number of e-foldings, say, 50–60, then the initial θ_i angle can be inhomogeneous on the scale of order today's horizon. This

translates directly into an inhomogeneity in η on the corresponding scale. Recent Planck results have bounded the curvature of the Universe to be small [7],

$$|\Omega_k| \lesssim 10^{-2}, \tag{77}$$

which indicates that inflation did last at least slightly longer than the absolute minimum required. But there exist various arguments regarding fine-tuning of the inflaton potential to suggest that it may not have lasted much more than this [42]. If this is the case, then a small spatial variation would arise in η. Since this would be a superhorizon mode, it would appear as a dipole across the Universe. Its magnitude would be correlated with Ω_k and potentially detectable. In fact it could conceivably be related to existing CMB anomalies [43].

Another possibility is that our pocket Universe arose from bubble nucleation in an eternally inflating false vacuum. In this case there exist arguments based on Coleman De Luccia [44] tunneling to indicate that the resulting bubble would be completely homogeneous (plus small quantum fluctuations). This would forbid the formation of such a dipole. On the other hand, this process would produce an infinite number of different bubbles, each with a different value of θ_i; a multiverse of different baryon densities.

B. Isocurvature fluctuations

Quantum fluctuations from inflation provide an excellent candidate for the origin of density fluctuations in the Universe. In simple single-field models, only a curvature ("adiabatic") fluctuation is generated, due to fluctuations in the inflaton. Earlier, in Eq. (61), we mentioned the amplitude of these fluctuations Δ_R^2.

For multifield inflationary models, an isocurvature ("entropic") fluctuation is also generated [45]. This is due to quantum fluctuations in the field orthogonal to the classical field trajectory, which leaves the total density unchanged. Since we are studying complex (two-field) inflation, there will be isocurvature fluctuations in ϕ, and these will generate isocurvature fluctuations in the baryon density.

The orthogonal fluctuations are in the angular variable θ, with mean value θ_i and $\langle|\delta\phi|^2\rangle_{\rm iso} = \rho_i^2\langle\delta\theta^2\rangle/2$. During inflation, this orthogonal fluctuation is a light (Goldstone) field with an approximately scale-invariant spectrum of fluctuations of amplitude [46]

$$\langle\delta\theta^2\rangle \approx \frac{\gamma^2}{4\pi^2}\frac{H_i^2}{\rho_i^2}, \tag{78}$$

where H_i is the Hubble parameter during inflation, ρ_i is the corresponding radial field value, and γ is an $\mathcal{O}(1)$ factor that accounts for Brownian motion during inflation.

In the radiation era, after all annihilations have occurred, but well before equality, the baryon density is related to θ by $n_b \propto -\sin(n\theta)$. Working to first order, we can relate the early-time baryon fluctuations to the θ fluctuations as

$$\left.\frac{\delta n_b}{n_b}\right|_{\rm early} = n\delta\theta\cot(n\theta_i). \tag{79}$$

Note that near special values of the initial angle $\theta_i = \frac{p\pi}{n}|p \in \mathbb{Z}$, this leading-order approximation breaks down. This first-order approximation is valid whenever there is a non-negligible average baryon number, which is necessary for baryogenesis.

We would now like to determine the corresponding isocurvature fluctuations in the CMB temperature. By definition, the isocurvature fluctuation satisfies

$$0 = \delta\rho_{\rm iso} = m_b\delta n_b + m_{\rm cdm}\delta n_{\rm cdm} + \delta\rho_\gamma, \tag{80}$$

where we have included baryons, dark matter, and radiation. We assume that for all species s, other than baryons, the fluctuations are adiabatic, i.e., $\delta n_s/n_s - \delta n_\gamma/n_\gamma = 0$ for $s \neq b$. If we divide Eq. (80) by $\rho_b = m_b n_b$, it is easy to show that this implies $\delta n_b/n_b \gg \delta\rho_\gamma/\rho_\gamma (= 4\delta T/T = (4/3)\delta n_\gamma/n_\gamma)$ at early times. Hence, the entropy perturbation is

$$\frac{\delta\eta}{\eta} = \frac{\delta n_b}{n_b} - \frac{\delta n_\gamma}{n_\gamma} \approx \left.\frac{\delta n_b}{n_b}\right|_{\rm early}. \tag{81}$$

Using this result and using the fact that entropy perturbations are approximately conserved outside the horizon, we find that the isocurvature contribution to the temperature at late time reentry is

$$\left(\frac{\delta T}{T}\right)_{\rm iso} = -\frac{6}{15}\frac{\Omega_b}{\Omega_m}\frac{\delta\eta}{\eta}. \tag{82}$$

We have included a factor of 6/5 due to the Sachs– Wolfe effect [47] (the details for which we shall not go into here). Putting together Eqs. (78), (79), (81), (82) leads to the following result for the squared isocurvature fluctuations:

$$\left\langle\left(\frac{\delta T}{T}\right)^2\right\rangle_{\rm iso} \approx \frac{9\gamma^2}{225\pi^2}\frac{\Omega_b^2}{\Omega_m^2}\frac{n^2H_i^2}{\rho_i^2}\cot^2(n\theta_i). \tag{83}$$

We now compute the adiabatic temperature fluctuation. Recall from Eq. (61) the formula for the adiabatic density fluctuations. To convert this to an adiabatic temperature fluctuation requires averaging over spherical harmonics, etc. The net result is roughly a factor of 1/20 difference, i.e.,

$$\left\langle\left(\frac{\delta T}{T}\right)^2\right\rangle_{\rm adi} \approx \frac{1}{20}\frac{H_i^2}{8\pi^2M_{\rm Pl}^2\epsilon_{sr}}, \tag{84}$$

(using $V_i \approx 3H_i^2 M_{\rm Pl}^2$). Since the isocurvature fluctuation is small, we can use this adiabatic fluctuation as an approximation for the total fluctuations.

The ratio of the isocurvature fluctuations to the total fluctuations is called α_{II}. Putting together Eqs. (83), (84), we obtain

$$\alpha_{II} \approx \frac{32\gamma^2}{5}\frac{\Omega_b^2}{\Omega_m^2}\frac{n^2 M_{\rm Pl}^2 \epsilon_{sr}}{\rho_i^2}\cot^2(n\theta_i). \tag{85}$$

Planck data reveal that the baryon-to-matter ratio is $\Omega_b/\Omega_m \approx 0.16$. Let us take $\gamma \sim 2$ as a representative value and $\cot(n\theta_i) \sim 1$, which gives an isocurvature fraction,

$$\alpha_{II} \sim 0.7\frac{n^2 M_{\rm Pl}^2 \epsilon_{sr}}{\rho_i^2}. \tag{86}$$

Note that this result holds for any simple inflationary potential.

If we now specialize to the case of quadratic inflation, we have $\epsilon_{sr} \approx 1/(2N_e)$ and $\rho_i \approx 2\sqrt{N_e} M_{\rm Pl}$. Then setting $N_e \approx 55$, we have our prediction for the isocurvature fraction:

$$\alpha_{II} \sim 3 \times 10^{-5} n^2. \tag{87}$$

Recent Planck results have provided an upper bound on cold dark matter isocurvature fluctuations of [7]

$$\alpha_{II} < 3.9 \times 10^{-2}, \qquad 95\% \text{ confidence}, \tag{88}$$

and we shall use this as a rough bound on baryon isocurvature fluctuations. This suggests that only ridiculously large values of n ($n > 36$) are ruled out, but such values are unrealistic anyhow. For the lowest value of n, namely, $n = 3$, we predict $\alpha_{II} \sim 3 \times 10^{-4}$, i.e., two orders of magnitude below the current bound. On the other hand, in Sec. VI A we explained that moderately high values of n are especially interesting. For instance, if we take $n = 8, 10, 12$, then our prediction is $\alpha_{II} \sim 3 \times 10^{-3}$, i.e., only 1 order of magnitude below the current bound. This is quite exciting as it is potentially detectable in the next generation of data.

VIII. DISCUSSION AND CONCLUSIONS

In this work we have proposed a way to directly unify early Universe inflation and baryogenesis, with motivation from the Affleck–Dine mechanism. We developed in great detail the basic proposal summarized in our accompanying paper [15].

A. Inflationary models

As a concrete example, we studied the simplest inflation model; a quadratic ("chaotic") inflation potential. Other potential functions, such as hilltop models, cosine potentials ("natural inflation"), or other non-polynomial potentials which are concave down during inflation, are marginally preferred by recent CMB data [7]. Our methods are directly applicable to these cases and can be adopted straightforwardly. Our idea is simply to use a symmetric potential for inflation, under a $U(1)$ global symmetry, and then introduce a sufficiently weak breaking that does not spoil the flatness of the potential. This will generate a particle number during the latter stage of inflation. Other inflationary models, which go far beyond this minimal inflationary setup, such as models dominated by higher-derivative kinetic terms, appear to be disfavored for over predicting non-Gaussianity, etc.

B. Particle physics models

We proposed two interesting particle physics models of this idea. The first model was to promote the $U(1)$ breaking to the level of a good symmetry until high-dimension operators $n \geq 8$. This allows ϕ to be a gauge singlet and then decay to quarks (and leptons) through dimension-7 operators while still conserving baryon number. If the decay is controlled by the scale $\Lambda \sim 10^{15-16}$ GeV, then we obtain good agreement with the observed baryon-to-photon ratio. Importantly, this value for Λ is precisely in the regime of validity of the field theory, between the Hubble scale and the Planck scale. The second model was to promote ϕ to a colored scalar and allow decay to quarks readily through lower-dimension operators. This requires a very small breaking parameter, which is technically natural, or a very small coupling to quarks, which deserves further exploration in contexts such as supersymmetry.

C. Inflaton constraints

An important parameter in the analysis is the inflaton mass m. As can be seen in Eq. (67), the required decay rate to obtain the observed asymmetry scales as $\Gamma_{\phi,\rm req} \propto m^2$. So if m is much smaller than the $m \sim 10^{13}$ GeV used as a reference value in this work, then the required decay rates become much smaller. If decay occurs through high-dimension operators, then this is easily achieved, while if decay occurs through low-dimension operators, then this becomes more difficult. In simple models of inflation, the mass of the inflaton tends to be related, within an order of magnitude or so, to the Hubble parameter. So this intertwines parameters of high-energy particle physics and the energy scale of inflation in an interesting way (complimentary to the difficulty at low energies [48]). In particular, this means that decay through low-dimension operators, as would be allowed by a colored inflaton, tends to favor high-scale inflation. In turn this favors appreciable tensor modes, which are being actively searched for.

D. Large-scale dipole

A distinguishing feature of these models, compared to other more common forms of baryogenesis, is that the dynamics respect the C and CP symmetry. Instead it is broken spontaneously by the initial state of the inflaton in the complex plane. For an initially inhomogeneous inflaton field, different regions of the Universe will acquire different baryon-to-photon ratios in the late Universe. This is an exciting property of the theory. This would allow for a large-scale dipole in the baryon density in the Universe and could even be relevant to CMB anomalies [43]. This is analogous to, and may be correlated with, a large-scale dipole in the dark-matter-to-photon density in the Universe. Indeed the latter can occur if the dark matter is comprised of axions with a large Peccei–Quinn scale [49,50]. So if one were to observe a dipole in one or both of these densities, it would provide tremendous clues about the early Universe and fundamental physics. Alternatively, these effects would be small if there were many e-foldings of inflation or if our pocket Universe arose from bubble nucleation.

E. Isocurvature—prediction

We also found that these models predict a baryonic isocurvature fluctuation at a level consistent with current observational bounds, and summarized in Eqs. (85), (86) for any symmetric potential V_s. For the quadratic inflation case, we found an isocurvature fraction, 1 or 2 orders of magnitude below the current bounds, potentially detectable in the next generation of experiments. For other choices of the symmetric potential V_s, the slow-roll parameter ϵ_{sr} will be different. This parameter is related to the tensor-to-scalar ratio by $r = 16\epsilon_{sr}$. For simple models of inflation, $\rho_i \sim \text{few}\, M_{\rm Pl}$, so all parameters are essentially fixed, leaving a relationship between the isocurvature fraction α_{II} and the tensor-to-scalar ratio r. An observational confirmation of this relationship would make our proposal quite compelling.

F. Isocurvature—comparison

Let us compare the aforementioned isocurvature fluctuation to the usual Affleck–Dine scenario where ϕ is not the inflaton. In that case, there is no obvious reason why the vev of the field during inflation $\rho_i = \sqrt{2}|\phi_i|$ must be larger than $M_{\rm Pl}$. Instead if ρ_i is somewhat smaller, say of order the GUT scale, then the isocurvature fraction would be very large (at least for reasonably high-scale inflation) and already ruled out [51]. By contrast, the current proposal of identifying ϕ with the inflaton naturally explains why ρ_i is of the order of or slightly larger than $M_{\rm Pl}$, which is especially interesting.

G. Small scales

Another important subject is the possibility of inhomogeneity on very small scales. In this paper we focused on the homogeneous mode of the inflaton (plus large-scale inhomogeneities in Sec. VII). A potentially important consideration is the possibility of preheating after inflation [52]. Under certain conditions, nonlinear dynamics will lead to explosive production of high k modes, and the breakup of ϕ, possibly into objects such as Q balls [53,54] and oscillons [55–57]. If this is efficient, it could throw the field up the potential in different ways in different patches of the Universe. This means that the effective θ_i could be different in different patches of the universes on small scales. After decay, this would lead to patches of baryons and antibaryons [58], which would presumably annihilate during thermalization, reducing the final value of η. On the other hand, this explosive process may raise the reheat temperature, raising the final value of η. So there are potentially competing effects. Such a process should only be important if the potential is strongly nonlinear, which is not the case for quadratic inflation, or if there is significant couplings to other bosonic fields, such as $\sim\phi^*\phi H^\dagger H$. These considerations will be important for some inflationary models and are a topic of future work.

ACKNOWLEDGMENTS

We would like to thank Mustafa Amin, Michael Dine, Alan Guth, Jesse Thaler, Tanmay Vachaspati, and Frank Wilczek for discussion, and we would like to acknowledge support by the Center for Theoretical Physics at MIT. This work is supported by the U.S. Department of Energy under cooperative research agreement Contract No. DE-FG02-05ER41360. J. K. is supported by a NSERC PDF fellowship.

APPENDIX: ANALYTICAL VERSION OF AFFLECK–DINE

In Ref. [16] Affleck and Dine studied a simpler problem. They studied a scalar field in a (noninflationary) background, with a known Hubble parameter, say $H = 2/3t$ for matter era or $H = 1/2t$ for radiation era, with some initial starting value for the field, say ρ_i. To solve this problem to order λ, the authors used a perturbative technique, where they solved for the field at zeroth order in λ and then used this as a source term to solve for the field at first order in λ, where the latter solution involved tracking the full complex field. This led to an approximation for the asymmetry A in terms of some constants that they obtained numerically.

Here we mention that by using the techniques of Secs. II A, II B we can solve this problem much more rapidly (only needing the zeroth-order solution and solving only a single ordinary differential equation, rather than a coupled system), and we obtain the coefficients analytically.

At zeroth order the equation of motion for the field (after factorizing for $e^{i\theta_i}$) is

$$\ddot{\rho}_0 + 3H\dot{\rho}_0 + m^2\rho_0 = 0, \tag{A1}$$

with H specified by the background (here we will not rescale the field by $M_{\rm Pl}$ to make it dimensionless, since $M_{\rm Pl}$ is not relevant in this computation). In the Affleck– Dine case, one recognizes that the field is frozen at early times when $H \gg m$. So we impose initial conditions: $\dot{\rho} = 0$ at early times. It is then easy to solve this differential equation. For the matter and radiation cases, the solution is

$$\rho_0(t) = \frac{\sin(mt)}{mt}\rho_i, \qquad \text{Matter} \tag{A2}$$

$$\rho_0(t) = 2^{1/4}\Gamma(5/4)\frac{J_{1/4}(mt)}{(mt)^{1/4}}\rho_i, \qquad \text{Radiation,} \tag{A3}$$

where Γ is the (complete) gamma function. This allows us to readily perform the integral that appears in ΔN_ϕ [Eq. (12)] by taking the limits of integration $t_i = 0$ and $t_f = \infty$. We write the scale factor in each case as

$$a = a_0\left(\frac{t}{t_0}\right)^{2/3}, \qquad \text{Matter} \tag{A4}$$

$$a = a_0\left(\frac{t}{t_0}\right)^{1/2}, \qquad \text{Radiation} \tag{A5}$$

in terms of some arbitrary reference constants a_0, t_0. The integral that appears in ΔN_ϕ

$$I \equiv \int_0^\infty dt a(t)^3\rho_0(t)^n, \tag{A6}$$

can be easily performed in each of the eras. We find

$$I = \frac{a_0^3\rho_i^n}{t_0^2 m^3}b_n, \qquad \text{Matter} \tag{A7}$$

$$I = \frac{a_0^3\rho_i^n}{t_0^{3/2}m^{5/2}}d_n, \qquad \text{Radiation,} \tag{A8}$$

where

$$b_n \equiv \int_0^\infty d\tau\tau^{2-n}\sin(\tau)^n, \tag{A9}$$

$$d_n \equiv 2^{n/4}\Gamma(5/4)^n\int_0^\infty d\tau\tau^{3/2-n/4}J_{1/4}(\tau)^n. \tag{A10}$$

The first few values of these constants are

$$b_3 = \frac{\pi}{4}, \qquad b_4 = \frac{\pi}{4}, \qquad b_5 = \frac{5\pi}{32},$$
$$b_6 = \frac{\pi}{8}, \qquad b_7 = \frac{77\pi}{768}, \qquad b_8 = \frac{\pi}{12}, \tag{A11}$$

$$d_3 = \frac{2\Gamma(5/4)^3}{3^{1/4}\sqrt{\pi}\Gamma(3/4)} \approx 0.521,$$
$$d_4 = \frac{4\Gamma(5/4)^5}{\sqrt{\pi}\Gamma(3/4)^3} \approx 0.750. \tag{A12}$$

To compute A we need to divide by the energy density ϵ_0 at late times. It is simple to show that at late times we have

$$\epsilon_0 = \frac{\rho_i^2}{2t^2}, \qquad \text{Matter} \tag{A13}$$

$$\epsilon_0 = \frac{\sqrt{2}\Gamma(5/4)^2m^2\rho_i^2}{\pi(mt)^{3/2}}, \qquad \text{Radiation.} \tag{A14}$$

Now recall the definition of the asymmetry parameter A from Eqs. (12), (16), (17). Putting the pieces together in the Affleck–Dine regime, we obtain

$$A = -\tilde{b}_n\frac{\lambda\rho_i^{n-2}}{m^2}\sin(n\theta_i), \qquad \text{Matter} \tag{A15}$$

$$A = -\tilde{d}_n\frac{\lambda\rho_i^{n-2}}{m^2}\sin(n\theta_i), \qquad \text{Radiation,} \tag{A16}$$

where

$$\tilde{b}_n \equiv \frac{n}{2^{n/2-2}}b_n, \tag{A17}$$

$$\tilde{d}_n \equiv \frac{\pi n}{2^{(n-1)/2}\Gamma(5/4)^2}d_n. \tag{A18}$$

This is in rough agreement with Ref. [16], where the authors computed the constant prefactors for $n = 4$ numerically using their alternate technique. Their numerics is found to be in small error ~10%–20% from the correct values obtained analytically here.

[1] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012).
[2] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **716**, 30 (2012).
[3] A. D. Sakharov, Pis'ma Zh. Eksp. Teor. Fiz. **5**, 32 (1967) [JETP Lett. **5**, 24 (1967)]; Usp. Fiz. Nauk **161**, 61 (1991) [Sov. Phys. Usp. **34**, 392 (1991)].
[4] A. Guth, Phys. Rev. D **23**, 347 (1981).
[5] A. D. Linde, Phys. Lett. **108B**, 389 (1982).
[6] G. Hinshaw *et al.* (WMAP Collaboration), Astrophys. J. Suppl. Ser. **208**, 19 (2013).
[7] P. A. R. Ade *et al.* (Planck Collaboration), arXiv:1303.5082.
[8] A. Riotto, arXiv:hep-ph/9807454.
[9] J. M. Cline, arXiv:hep-ph/0609145.
[10] W. Buchmuller, arXiv:0710.5857.
[11] M. Trodden, Rev. Mod. Phys. **71**, 1463 (1999).
[12] A. Menon and D. E. Morrissey, Phys. Rev. D **79**, 115020 (2009).
[13] J. M. Cline, arXiv:hep-ph/0201286.
[14] N. Rius and V. Sanz, Nucl. Phys. **B570**, 155 (2000).
[15] M. P. Hertzberg and J. Karouby, arXiv:1309.0007.
[16] I. Affleck and M. Dine, Nucl. Phys. **B249**, 361 (1985).
[17] R. Allahverdi and A. Mazumdar, New J. Phys. **14**, 125013 (2012).
[18] M. Dine, L. Randall, and S. D. Thomas, Nucl. Phys. **B458**, 291 (1996).
[19] M. Dine, L. Randall, and S. D. Thomas, Phys. Rev. Lett. **75**, 398 (1995).
[20] Y.-Y. Charng, D.-S. Lee, C. N. Leung, and K.-W. Ng, Phys. Rev. D **80**, 063519 (2009).
[21] K. Enqvist and J. McDonald, Phys. Rev. Lett. **83**, 2510 (1999).
[22] A. Mazumdar and A. Perez-Lorenzana, Phys. Rev. D **65**, 107301 (2002).
[23] K. Koyama and J. Soda, Phys. Rev. Lett. **82**, 2632 (1999).
[24] R. Allahverdi, M. Drees, and A. Mazumdar, Phys. Rev. D **65**, 065010 (2002).
[25] O. Seto, Phys. Rev. D **73**, 043509 (2006).
[26] S. Kasuya and M. Kawasaki, Phys. Rev. D **74**, 063507 (2006).
[27] B. Dutta and K. Sinha, Phys. Rev. D **82**, 095003 (2010).
[28] D. Marsh, J. High Energy Phys. 05 (2012) 041.
[29] M. A. G. Garcia and K. A. Olive, J. Cosmol. Astropart. Phys. 09 (2013) 007.
[30] K. Enqvist and J. McDonald, Phys. Rev. Lett. **81**, 3071 (1998).
[31] D. Delepine, C. Martnez, and L. A. Urena-Lopez, Phys. Rev. Lett. **98**, 161302 (2007).
[32] M. Bastero-Gil, A. Berera, R. O. Ramos, and J. G. Rosa, Phys. Lett. B **712**, 425 (2012).
[33] R. Kitano, H. Murayama, and M. Ratz, Phys. Lett. B **669**, 145 (2008).
[34] H. Murayama, H. Suzuki, T. Yanagida, and J. 'i. Yokoyama, Phys. Rev. D **50**, R2356 (1994).
[35] H. Murayama, H. Suzuki, T. Yanagida, and J. 'i. Yokoyama, Phys. Rev. Lett. **70**, 1912 (1993).
[36] A. D. Linde, Phys. Lett. **129B**, 177 (1983)
[37] L. Kofman, A. D. Linde, and A. A. Starobinsky, Phys. Rev. D **56**, 3258 (1997).
[38] A. R. Liddle and D. H. Lyth, *Cosmological Inflation and Large Scale Structure* (Cambridge University Press, Cambridge, England, 2000).
[39] B. Grzadkowski, M. Iskrzynski, M. Misiak, and J. Rosiek, J. High Energy Phys. 10 (2010) 085.
[40] D. J. Gross and F. Wilczek, Phys. Rev. Lett. **30**, 1343 (1973).
[41] H. D. Politzer, Phys. Rev. Lett. **30**, 1346 (1973).
[42] B. Freivogel, M. Kleban, M. Rodriguez Martinez, and L. Susskind, J. High Energy Phys. 03 (2006) 039.
[43] P. A. R. Ade *et al.* (Planck Collaboration), arXiv:1303.5083.
[44] S. R. Coleman and F. De Luccia, Phys. Rev. D **21**, 3305 (1980).
[45] N. Bartolo, S. Matarrese, and A. Riotto, Phys. Rev. D **64**, 123504 (2001).
[46] A. D. Linde, Prog. Theor. Phys. Suppl. **163**, 295 (2006).
[47] R. K. Sachs and A. M. Wolfe, Astrophys. J. **147**, 73 (1967).
[48] M. P. Hertzberg, J. Cosmol. Astropart. Phys. 08 (2012) 008.
[49] M. Tegmark, A. Aguirre, M. Rees, and F. Wilczek, Phys. Rev. D **73**, 023505 (2006).
[50] M. P. Hertzberg, M. Tegmark, and F. Wilczek, Phys. Rev. D **78**, 083507 (2008).
[51] S. Kasuya, M. Kawasaki, and F. Takahashi, J. Cosmol. Astropart. Phys. 10 (2008) 017.
[52] L. Kofman, arXiv:astro-ph/9802221.
[53] S. Coleman, Nucl. Phys. **B262**, 2 (1985).
[54] K. Enqvist and J. McDonald, Phys. Lett. B **425**, 309 (1998).
[55] E. J. Copeland, M. Gleiser, and H.-R. Muller, Phys. Rev. D **52**, 1920 (1995).
[56] M. A. Amin, R. Easther, and H. Finkel, J. Cosmol. Astropart. Phys. 12 (2010) 001.
[57] M. A. Amin, R. Easther, H. Finkel, R. Flauger, and M. P. Hertzberg, Phys. Rev. Lett. **108**, 241302 (2012).
[58] A. D. Dolgov, M. Kawasaki, and N. Kevlishvili, Nucl. Phys. **B807**, 229 (2009).