

Affleck-Dine curvaton

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The Standard Model of particle physics does not explain the origin of the universe's baryon asymmetry or its primordial fluctuations. The Affleck-Dine mechanism is a well motivated scenario for generating the baryon asymmetry through the postinflationary dynamics of a complex scalar field with baryon number. The curvaton mechanism is a popular approach for producing curvature perturbations through the dynamics of a light spectator field that decays after inflation. We demonstrate that the same complex field can viably perform both roles without any modifications to the minimal realization of Affleck-Dine baryogenesis. This scenario can also accommodate appreciable levels of primordial non-Gaussianity, beyond those achievable with only a real-valued curvaton field, and may be observable with future cosmic microwave background experiments.

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I. INTRODUCTION

Our universe exhibits nearly scale invariant perturbations that seed cosmic structures and imprint temperature anisotropies onto the cosmic microwave background (CMB) [1]. Since these fluctuations are correlated on scales outside the cosmological horizon during recombination, they can be dynamically generated during a period of inflation when the comoving Hubble radius decreases with time and modes become frozen on superhorizon scales. When inflation ends and the horizon begins to grow, they eventually reestablish causal contact, but retain large scale correlations—see Ref. [2] for a review. Since the Standard Model (SM) does not explain the origin of these fluctuations, their existence is technically evidence of new physics.¹

Although it is often assumed that these fluctuations arise from the field driving inflation, this need not be the case. Curvaton models [4–6] posit an additional light scalar field that is a spectator during inflation, during which it acquires isocurvature perturbations. After reheating, the curvaton decays to radiation, and its isocurvature perturbations are converted to adiabatic. The result is generically larger primordial non-Gaussianities, as compared with models where a single inflaton generates perturbations during slow-roll evolution.

Generating the baryon asymmetry of the universe also necessitates new physics [7]. Although the SM does include baryon number violating and both C - and CP -violating interactions, the amounts of violation are insufficient to explain the observed asymmetry. Further, these processes must occur out-of-equilibrium. While the electroweak phase transition could in principle provide such conditions, extensions are necessary to make it strongly first order [8]. As it stands, the SM does not provide a complete framework to account for the observed baryon-to-photon ratio $\eta_B = n_B/n_\gamma = 6 \times 10^{-10}$ [1]—see Ref. [9] for a review.

A popular scenario for generating the baryon asymmetry is the Affleck-Dine mechanism, in which a complex scalar field with nonzero baryon number is initially displaced from the minimum of its potential in the early universe [10]. If the potential contains feeble baryon violating self-interactions, as the scalar evolves the comoving baryon

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¹In principle, the SM Higgs field could drive inflation and generate the curvature perturbation, but viable realizations require, e.g., nonminimal couplings to gravity [3].

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number begins to accumulate until the scalar decays to SM particles through baryon preserving operators to transfer this asymmetry to the SM plasma. Unlike most other mechanisms, Affleck-Dine stands out for generically predicting too large of an asymmetry unless the parameters that govern the baryon density have extremely small values.

Here we demonstrate that the same complex scalar field realizing Affleck-Dine baryogenesis can also serve as the curvaton without sacrificing any of the appealing features of either mechanism.² Specifically, we find that within this combined framework, the observed baryon asymmetry can be generated with the same model parameters that also yield the observed curvature perturbation. A distinctive feature of this scenario is larger levels of primordial non-Gaussianity, which can in principle be observed in future CMB experiments.

II. AFFLECK-DINE REVIEW

We begin by briefly reviewing the Affleck-Dine mechanism for generating the baryon asymmetry [10]. Consider extending the SM to include a complex scalar field χ with nonzero baryon number B_χ .³ It is useful to decompose χ into a modulus r and complex phase θ ,

$$\chi = \frac{r}{\sqrt{2}} e^{i\theta}. \quad (1)$$

Under global phase rotation, χ transforms as $\chi \rightarrow e^{i\alpha B_\chi} \chi$, with α an arbitrary constant. The baryon number density stored in the χ field is identified with the timelike component of the baryon Noether current

$$n_B = iB_\chi(\dot{\chi}^* \chi - \chi^* \dot{\chi}) = B_\chi r^2 \dot{\theta}, \quad (2)$$

which is only nonzero when the complex phase is dynamical in the early universe.

The potential for χ has both baryon preserving and (small) baryon violating self-interactions

$$V(\chi) = m_\chi^2 |\chi|^2 + \frac{\lambda}{2} |\chi|^4 + \frac{\lambda'}{4} (\chi^4 + \chi^{*4}), \quad (3)$$

where the last term explicitly breaks the baryon number, so we demand $\lambda' \ll \lambda$. In terms of r and θ , the action is

$$S = \int d^4x a^3 \left(\frac{1}{2} \dot{r}^2 + \frac{1}{2} r^2 \dot{\theta}^2 - V(r, \theta) \right), \quad (4)$$

²It has previously been shown that the inflaton can double as the Affleck-Dine field in certain inflationary contexts—see, e.g., Refs. [11–14]. Here our key point is that if the Affleck-Dine field is not the inflaton, but rather a light spectator during inflation, then it automatically plays the role of a curvaton without additional requirements.

³This is a well-motivated scenario that arises in many supersymmetric extensions to the SM [15].

where a is the scale factor. The radial equation of motion can be written

$$\ddot{r} + 3H\dot{r} - \dot{\theta}^2 r = -\frac{\partial V}{\partial r} \approx -m_\chi^2 r, \quad (5)$$

where $H = \dot{a}/a$ is the Hubble rate and for simplicity we have assumed $\chi \ll m_\chi/\sqrt{\lambda}$ such that the potential is always dominated by the quadratic interaction.⁴ Using Eq. (2), the equation of motion for θ can be written in terms of the baryon number density

$$\frac{1}{a^3} \frac{\partial}{\partial t} (a^3 n_B) = -\frac{\lambda' B_\chi}{2} r^4 \sin(4\theta). \quad (6)$$

If $\lambda' = 0$, the comoving baryon number is conserved, in analogy with angular momentum conservation in classical systems with radially symmetric potentials. Integrating Eq. (6) with the symmetric initial condition $n_B(t_i) = 0$, we obtain

$$n_B(t) = \frac{\lambda' B_\chi}{2a^3(t)} \int_{t_i}^t dt' a^3(t') r^4(t') \sin[4\theta(t')], \quad (7)$$

which is valid for $\tau_\chi > t > t_i$. Here, τ_χ is the χ lifetime and t_i is the initial time, which we identify with reheating.

To transfer the asymmetry, χ must also couple to SM fields through baryon number *preserving* interactions of the form

$$\mathcal{L}_{\text{int}} = c\chi \hat{\mathcal{O}}_{\text{SM}} + \text{H.c.}, \quad (8)$$

where $\hat{\mathcal{O}}_{\text{SM}}$ is a composite operator of SM fields with baryon number $B_{\hat{\mathcal{O}}} = -B_\chi$ and c is a constant coefficient.⁵ To be precise, $\hat{\mathcal{O}}_{\text{SM}}$ should carry the net $B - L$ number if the decays that serve to transfer the asymmetry occur above the electroweak scale, while electroweak sphalerons are still active. Any operator with a net baryon number suffices to transfer the asymmetry if the reheat temperature is below the weak scale.

Thus, viable baryogenesis can be achieved if χ decays to SM fields through the interaction in Eq. (8) before the asymmetry is depleted in the early universe. If there are no other decay channels for χ , the baryon-to-photon ratio can be written

⁴See, e.g., Refs. [16–19] for model-building efforts with a self-interacting curvaton.

⁵For additional flexibility, it is often advantageous to also introduce a new fermion ψ that does not carry baryon number, such that one can write a Yukawa term of the form $\mathcal{L}_{\text{int}} = y \bar{q} \psi \chi$, where q may be a composite operator of SM fields carrying baryon number $B_q = -B_\chi$.

$$\eta_B = \frac{\pi^2}{4\zeta(3)} \frac{\lambda' B_\chi}{T_i^3} \int_{t_i}^{\tau_\chi} dt' a^3(t') r^4(t') \sin[4\theta(t')], \quad (9)$$

where $T_i \equiv T_\gamma(t_i)$ is the initial photon temperature. In our later analysis, we take $B_\chi = 1$ without loss of generality.

III. CURVATON REVIEW

In the original curvaton scenario [4–6], a real scalar field σ is present as a light spectator during inflation. The energy density of the curvaton is always subdominant during inflation, such that it does not drive Hubble expansion; it is, however, responsible for eventually generating the dominant contribution to the curvature perturbation. The minimal realization posits a quadratic background potential $V(\sigma) = \frac{1}{2} m_\sigma^2 \sigma^2$ with $m_\sigma \ll H_i$, such that the curvaton acquires Gaussian isocurvature fluctuations of amplitude $\delta\sigma_i \sim H_i/(2\pi)$, which are the dominant influence on its evolution. Here $H = \dot{a}/a$ is the Hubble rate and the subscript i denotes a value during inflation.

When inflation ends, the inflaton decays to reheat the universe. The curvaton is initially Hubble damped, but begins to oscillate about its quadratic potential when $H \sim m_\sigma$. During this oscillatory phase, the curvaton redshifts like nonrelativistic matter $\rho_\sigma \propto a^{-3}$ and the energy density in σ grows relative to that in radiation. Finally, when $H \sim \tau_\sigma^{-1}$, where τ_σ is the curvaton lifetime, σ decays to SM radiation and its isocurvature perturbations are converted to adiabatic perturbations, presuming the decay products thermalize with the background plasma.⁶

To calculate ζ in this scenario, it is useful to employ the δN formalism [20–24], in which the curvaton curvature perturbation ζ_σ on a general hypersurface of uniform curvaton density is [25]

$$\zeta_\sigma(t, \vec{x}) = \delta N(t, \vec{x}) + \frac{1}{3} \log \left(\frac{\rho_\sigma(t, \vec{x})}{\bar{\rho}_\sigma(t)} \right), \quad (10)$$

where δN is the perturbed number of e -folds, ρ_σ is the curvaton energy density, and bars denote background quantities. On a uniform *total* density slice, we can equate δN with the total curvaton perturbation $\zeta = \delta N$ and write

$$\rho_\sigma(t, \vec{x}) = e^{3(\zeta_\sigma - \zeta)} \bar{\rho}_\sigma(t), \quad \rho_R(t, \vec{x}) \approx e^{-4\zeta} \bar{\rho}_R(t), \quad (11)$$

where ρ_R is the energy density of SM radiation and we have assumed $\zeta_\sigma \gg \zeta_R$. At the time of curvaton decay, energy

conservation yields an implicit relation for the curvature perturbation

$$e^{4\zeta} - (e^{3\zeta_\sigma} \Omega_\sigma) e^\zeta + (\Omega_\sigma - 1) = 0, \quad (12)$$

where $\Omega_i \equiv \bar{\rho}_i(\tau_\sigma)/\bar{\rho}_{\text{tot}}(\tau_\sigma)$, $\bar{\rho}_{\text{tot}} = \bar{\rho}_\sigma + \bar{\rho}_R$ is the total background energy density, and we have assumed that total density is uniform on the decay surface. In terms of the curvaton energy fraction at decay $\kappa = 3\Omega_{\sigma, \tau_\sigma}/(4 - \Omega_{\sigma, \tau_\sigma})$, the general solution is

$$\zeta = \log \left(\frac{\beta^{1/2} + \sqrt{\alpha\kappa\beta^{-1/2} - \beta}}{(3 + \kappa)^{1/3}} \right), \quad (13)$$

where we have defined parameters

$$\beta = \frac{1}{2} [\gamma^{1/3} + (\kappa - 1)(\kappa + 3)^{1/3} \gamma^{-1/3}], \quad (14)$$

$$\gamma = (\alpha\kappa)^2 + \sqrt{(\alpha\kappa)^4 + (3 + \kappa)(1 - \kappa)^3}, \quad (15)$$

and for a strictly quadratic potential ζ_σ can be related to the initial field perturbation as [25]

$$\alpha \equiv e^{3\zeta_\sigma} = \left(1 + \frac{\delta\sigma_i}{\bar{\sigma}_i} \right)^2, \quad (16)$$

where $\bar{\sigma}_i$ and $\delta\sigma_i$ are the curvaton background value and perturbation during inflation, respectively. Thus, ζ depends only on the initial curvaton perturbation and the curvaton energy fraction at the time of decay.

The expression in Eq. (13) is a nonlinear mapping between the Gaussian perturbation ζ_σ and the total curvature perturbation ζ , so this scenario generically yields larger non-Gaussianities than models where a minimal inflaton generates all perturbations. Taylor expanding in the Gaussian variable ζ_g to third order yields

$$\zeta \simeq \zeta_\sigma + \frac{3}{5} f_{\text{NL}} \zeta_\sigma^2 + \frac{9}{25} g_{\text{NL}} \zeta_\sigma^3, \quad (17)$$

where we have defined the nonlinearity parameters f_{NL} and g_{NL} , constrained by the Planck Collaboration as [26]

$$f_{\text{NL}} = 5.0 \pm 8.4, \quad g_{\text{NL}} = (-5.8 \pm 6.5) \times 10^4, \quad (18)$$

which are both consistent with zero. For a discussion of primordial non-Gaussianity in curvaton models, see Ref. [6]. We now turn to generalizing this argument in the case of a complex curvaton identified with the Affleck-Dine scalar.

⁶If dark matter has been produced by this point and is decoupled from the radiation, then curvaton decays induce a mismatch between photon and dark matter perturbations, which can conflict with CMB measurements. It is straightforward to avoid this fate (e.g., by producing dark matter thermally after the curvaton decays), so we ignore this possibility. Thus, throughout our analysis we assume that dark matter does not introduce additional isocurvature perturbations.

IV. COMBINED SCENARIO

We now promote the original Affleck-Dine field χ from Sec. II to also play the role of the curvaton. The key conceptual difference in this scenario is that $\chi = \chi_R + i\chi_I$ is a complex field with 2 degrees of freedom

$$r = \sqrt{2(\chi_R^2 + \chi_I^2)}, \quad \theta \equiv \tan^{-1}\left(\frac{\chi_I}{\chi_R}\right), \quad (19)$$

which are related through the equations of motion in Eqs. (5) and (6). In our numerical analysis, these equations are supplemented with decay terms, and in Fig. 1 we show the χ energy density as a function of time for a benchmark model.

A. Curvature perturbation

Since $m_\chi \ll H_i$, the field components $\chi_{R,I}$ acquire uncorrelated quantum fluctuations $\delta\chi_R$ and $\delta\chi_I$ during inflation, which evolve after inflation to set the values of the corresponding curvature perturbations ζ_{χ_R} and ζ_{χ_I} . The curvature perturbation from two independent curvaton fields has been studied in Ref. [27], and we use this formalism to describe our scenario in which the dynamics of χ_R and χ_I are related through the role they play in setting the complex phase in Eq. (19).

Repeating the same sequence of steps that leads to Eq. (12), for a complex curvaton we find

$$\Omega_\gamma e^{4(\zeta_\gamma - \zeta)} + \Omega_{\chi_R} e^{3(\zeta_{\chi_R} - \zeta)} + \Omega_{\chi_I} e^{3(\zeta_{\chi_I} - \zeta)} = 1, \quad (20)$$

where $\Omega_i = \bar{\rho}_i(\tau_\chi)/\bar{\rho}_{\text{tot}}(\tau_\chi)$ are the density parameters of each field evaluated at the time of χ decay. We numerically evolve this system to obtain the parameters Ω_i and then solve for the perturbation of each field to determine ζ using

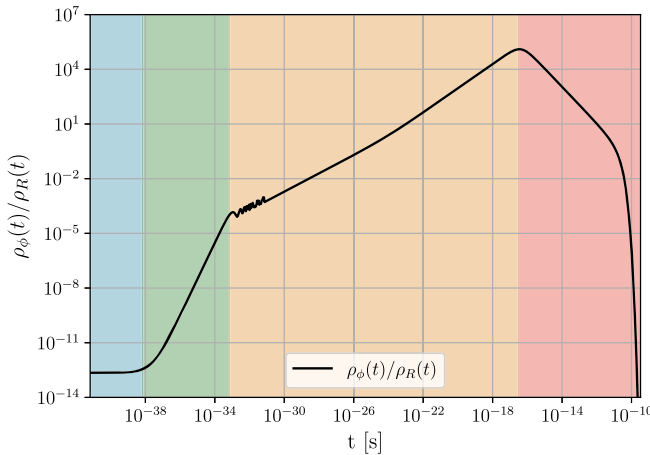


FIG. 1. The curvaton-radiation energy density ratio in various cosmic phases: the Hubble-damped regime (green), the quadratic oscillation regime (orange), and the χ decay regime (red). Here we assume the benchmark parameters from Model 1 in Table I.

Eq. (20). Taylor expanding Eq. (20) to first order in all perturbations, we find

$$4\Omega_\gamma(\zeta_\gamma - \zeta) + 3\Omega_{\chi_R}(\zeta_{\chi_R} - \zeta) + 3\Omega_{\chi_I}(\zeta_{\chi_I} - \zeta) = 0, \quad (21)$$

and the linear-order solution for the total curvature perturbation is then

$$\zeta = f_\gamma \zeta_\gamma + f_{\chi_R} \zeta_{\chi_R} + f_{\chi_I} \zeta_{\chi_I}, \quad (22)$$

where we have defined the parameters

$$f_\gamma = \frac{4\Omega_\gamma}{4\Omega_\gamma + 3\Omega_{\chi_R} + 3\Omega_{\chi_I}}, \quad (23)$$

$$f_{\chi_R} = \frac{3\Omega_{\chi_R}}{4\Omega_\gamma + 3\Omega_{\chi_R} + 3\Omega_{\chi_I}}, \quad (24)$$

$$f_{\chi_I} = \frac{3\Omega_{\chi_I}}{4\Omega_\gamma + 3\Omega_{\chi_R} + 3\Omega_{\chi_I}}. \quad (25)$$

Using Eq. (22), we can obtain the final curvature perturbation at the time of χ decay. Note, however, that these relations only hold at linear order, and calculating the primordial non-Gaussianity parameters—defined in Eq. (17)—requires a higher-order expansion. Thus, in Appendix A, we present the second-order relations we use to compute f_{NL} in our analysis.

B. Non-Gaussianity

In this subsection, we follow the treatment in Ref. [27] to extract f_{NL} in our scenario. To leading order, the primordial curvature perturbation power spectrum is

$$P_\zeta = f_{\chi_R}^2 P_{\chi_R} + f_{\chi_I}^2 P_{\chi_I}. \quad (26)$$

The power spectra of the curvaton components at the onset of curvaton oscillations $H \sim m_\chi$ are related by

$$P_{\chi_R} = \beta^2 P_{\chi_I}, \quad (27)$$

which holds at leading order. In the linear approximation, β is reduced to the ratio of the values of the curvaton fields at horizon crossing

$$\beta = \frac{\bar{\chi}_I}{\bar{\chi}_R} = \tan \theta, \quad (28)$$

where the bar denotes values during inflation. Since the Fourier components of χ_R and χ_I are uncorrelated, we have

$$\langle \zeta_{\chi_R}(\mathbf{k}) \zeta_{\chi_I}(\mathbf{k}') \rangle = 0. \quad (29)$$

Using also that the correlation function of an odd number of Gaussian perturbations vanishes, the three-point correlation function of ζ can be written

$$\begin{aligned}
& \langle \zeta(\mathbf{k}_1) \zeta(\mathbf{k}_2) \zeta(\mathbf{k}_3) \rangle \\
&= \frac{1}{2} f_{\chi_R}^2 \left(C - \frac{3}{2} A \right) \langle \zeta_{\chi_R}(\mathbf{k}_1) \zeta_{\chi_R}(\mathbf{k}_2) (\zeta_{\chi_R} * \zeta_{\chi_R})(\mathbf{k}_3) \rangle \\
&+ \frac{1}{2} f_{\chi_I}^2 \left(D - \frac{3}{2} B \right) \langle \zeta_{\chi_I}(\mathbf{k}_1) \zeta_{\chi_I}(\mathbf{k}_2) (\zeta_{\chi_I} * \zeta_{\chi_I})(\mathbf{k}_3) \rangle \\
&+ \frac{1}{2} f_{\chi_R} f_{\chi_I} E \langle \zeta_{\chi_R}(\mathbf{k}_1) \zeta_{\chi_I}(\mathbf{k}_2) (\zeta_{\chi_R} * \zeta_{\chi_I})(\mathbf{k}_3) \rangle \\
&+ (2 \text{ permutation terms}), \tag{30}
\end{aligned}$$

where the parameters A , B , C , D , and E are all functions of f_{χ_R} , f_{χ_I} , and f_γ defined in Appendix A, and an asterisk between functions signifies convolution

$$(\zeta_i * \zeta_j)(\mathbf{k}) = \int \frac{d^3 q}{(2\pi)^3} \zeta_i(\mathbf{q}) \zeta_j(\mathbf{k} - \mathbf{q}). \tag{31}$$

Rewriting Eq. (30) in terms of the power spectrum of ζ by using Eqs. (26) and (27), one finds

$$f_{\text{NL}} = \frac{5[\beta^4 f_{\chi_R}^2 (2C - 3A) + \beta^2 f_{\chi_R} f_{\chi_I} E + f_{\chi_I}^2 (2D - 3B)]}{6(\beta^2 f_{\chi_R}^2 + f_{\chi_I}^2)^2}, \tag{32}$$

which we use to obtain our numerical results.

C. Results

In Table I we present three benchmark realizations of our scenario with different curvaton masses. In each case, the baryon asymmetry and the primordial curvature perturbation both achieve their observed values. We also show the

TABLE I. Three benchmark models that predict acceptable values for ζ , the local nonlinearity parameter $f_{\text{NL}}^{\text{local}}$, η_B , the tensor-to-scalar ratio r_T , and α_{II} . For clarity, we have chosen rounded values of the input parameters, but simple order-unity variations of these choices can accommodate the measured values of $\eta_B = 6.1 \times 10^{-10}$, $\zeta = 4.6 \times 10^{-5}$, and $r_T < 0.036$ [30].

	Model 1	Model 2	Model 3
m_χ/GeV	10^9	10^{10}	10^8
Γ_χ/GeV	10^{-13}	10^{-13}	10^{-15}
H_i/M_{pl}	1×10^{-6}	1×10^{-7}	1×10^{-8}
r_i/H_i	1×10^4	5×10^3	1×10^4
$\theta_i/(\pi/8)$	1.1	1.1	1.1
λ	10^{-20}	10^{-15}	10^{-15}
λ'	10^{-24}	5×10^{-17}	10^{-19}
$\zeta[10^{-5}]$	-2.3	4.6	2.1
$f_{\text{NL}}^{\text{local}}$	-0.88	-0.88	-0.71
$\eta_B[10^{-10}]$	4.9	3.9	1.1
r_T	1.1×10^{-3}	3.0×10^{-6}	1.4×10^{-7}
α_{II}	1.1×10^{-3}	7.4×10^{-4}	8.7×10^{-4}

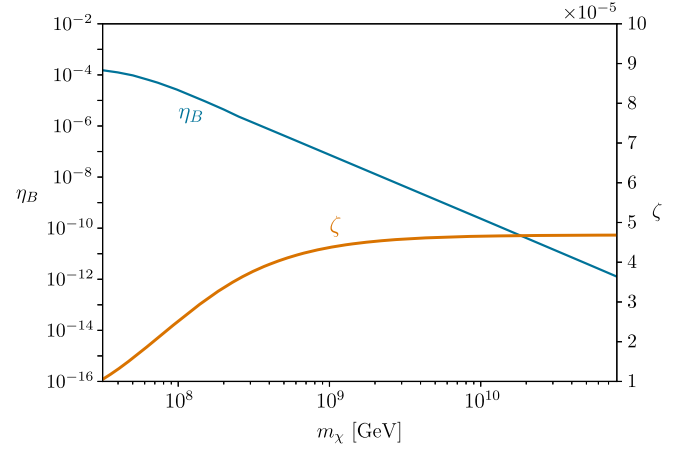


FIG. 2. The baryon-to-photon ratio η_B and curvature perturbation ζ as a function of curvaton mass m_χ . For all points along the η_B curve, ζ is chosen to have the observed value $\zeta \simeq 2 \times 10^{-5}$ and $\lambda' = 10^{-15}$. For all points along the ζ curve, $\eta_B = 6 \times 10^{-10}$ to match the observed asymmetry, λ' is chosen accordingly, and $H_i = 3 \times 10^{-5} M_{\text{pl}}$. Unless otherwise specified, all other parameters match the values from Model 2 in Table I.

degree of non-Gaussianity predicted in each scenario and find values of f_{NL} within the forecast sensitivities of CMB-S4 [28]. However, a complete analysis of the observational consequences is beyond the scope of this work. We have checked also that all benchmark points satisfy the bounds from baryon isocurvature (see Supplemental Material [29] for detail).

In Fig. 2, we show how η_B and ζ vary as a function of the curvaton mass, with other parameters held fixed. Here the blue curve represents η_B variation with ζ held fixed to its observed value and the red curve represents ζ variation with mass as η_B is fixed to its observed value.

V. DISCUSSION

In this article, we have shown that a complex curvaton can also generate the baryon asymmetry of the universe through the Affleck-Dine mechanism. We have found that this unified construction economically accommodates viable values for both the curvature perturbation and the baryon asymmetry. We emphasize that this is not a new model; the field content and interactions are identical to that of a minimal Affleck-Dine scenario. The key point is that this field serves as a curvaton without requiring any additional modifications. Indeed, if the Affleck-Dine field is light compared to the Hubble rate during inflation, it will necessarily acquire fluctuations and contribute to the total curvature perturbation.

In our treatment, we calculate the baryon asymmetry by solving the classical equations of motion that govern the homogeneous background value of χ . The real and imaginary parts of the curvaton acquire independent fluctuations during inflation, and we track their evolution using the δN

formalism to calculate the total curvature perturbation at the time of χ decay. Unlike in the minimal curvaton model, the additional degree of freedom allows for enhanced primordial non-Gaussianity.

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DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: SECOND-ORDER CALCULATION

In this section, we extend the calculation in the text to account for second-order effects required for non-Gaussianities. For each species $i = \chi_R, \chi_I, \gamma$, we take the Ansatz

$$\zeta_i = \zeta_i^{(1)} + \frac{1}{2}\zeta_i^{(2)}, \quad (\text{A1})$$

where the (1) and (2) label the first- and second-order perturbations, respectively. Expanding Eq. (20) to second order yields

$$\begin{aligned} & 4\Omega_\gamma \left[\frac{1}{2}(\zeta_\gamma^{(2)} - \zeta_\gamma^{(2)}) + 2(\zeta_\gamma^{(1)} - \zeta_\gamma^{(1)})^2 \right] \\ & + 3\Omega_{\chi_R} \left[\frac{1}{2}(\zeta_{\chi_R}^{(2)} - \zeta_{\chi_R}^{(2)}) + \frac{3}{2}(\zeta_{\chi_R}^{(1)} - \zeta_{\chi_R}^{(1)})^2 \right] \\ & + 3\Omega_{\chi_I} \left[\frac{1}{2}(\zeta_{\chi_I}^{(2)} - \zeta_{\chi_I}^{(2)}) + \frac{3}{2}(\zeta_{\chi_I}^{(1)} - \zeta_{\chi_I}^{(1)})^2 \right] = 0. \end{aligned} \quad (\text{A2})$$

Solving for the second-order perturbation, we find

$$\begin{aligned} \zeta^{(2)} = & f_\gamma [\zeta_\gamma^{(2)} + 4(\zeta_\gamma^{(1)} - \zeta_\gamma^{(1)})^2] \\ & + f_{\chi_R} [\zeta_{\chi_R}^{(2)} + 3(\zeta_{\chi_R}^{(1)} - \zeta_{\chi_R}^{(1)})^2] \\ & + f_{\chi_I} [\zeta_{\chi_I}^{(2)} + 3(\zeta_{\chi_I}^{(1)} - \zeta_{\chi_I}^{(1)})^2], \end{aligned} \quad (\text{A3})$$

where we have defined f_i as in Eq. (23). The expression above can be simplified using the form of the linear-order curvature perturbation $\zeta^{(1)}$ in Eq. (22). Further, we assume that the initial radiation is homogeneous and that the dominant contribution to the curvature perturbation is coming from the curvaton, allowing us to neglect the terms involving ζ_γ . This simplifies the result considerably, and we write

$$\zeta^{(2)} = A\zeta_{\chi_R}^{(2)} + B\zeta_{\chi_I}^{(2)} + C[\zeta_{\chi_R}^{(1)}]^2 + D[\zeta_{\chi_I}^{(1)}]^2 + E\zeta_{\chi_R}^{(1)}\zeta_{\chi_I}^{(1)}, \quad (\text{A4})$$

where

$$A = f_{\chi_R}, \quad (\text{A5})$$

$$B = f_{\chi_I}, \quad (\text{A6})$$

$$C = 3f_{\chi_R}(1 - f_{\chi_R})^2 + 3f_{\chi_R}^2 f_{\chi_I}, \quad (\text{A7})$$

$$D = 3f_{\chi_I}(1 - f_{\chi_I})^2 + 3f_{\chi_I}^2 f_{\chi_R}, \quad (\text{A8})$$

$$E = -6f_{\chi_R}f_{\chi_I}(2 - f_{\chi_R} - f_{\chi_I}). \quad (\text{A9})$$

We further assume that each curvaton field evolves linearly between inflationary horizon exit and the beginning of oscillations, which implies $g_a'' = g_b'' = 0$ in Eqs. (17) and (18) of Ref. [27]. This yields the relations

$$\zeta_{\chi_R}^{(2)} = -\frac{3}{2}[\zeta_{\chi_R}^{(1)}]^2, \quad \zeta_{\chi_I}^{(2)} = -\frac{3}{2}[\zeta_{\chi_I}^{(1)}]^2. \quad (\text{A10})$$

Thus, the second-order contribution to the curvature perturbation $\zeta = \zeta^{(1)} + \frac{1}{2}\zeta^{(2)}$ is

$$\zeta^{(2)} = \left(C - \frac{3}{2}A \right) [\zeta_{\chi_R}^{(1)}]^2 + \left(D - \frac{3}{2}B \right) [\zeta_{\chi_I}^{(1)}]^2 + E\zeta_{\chi_R}^{(1)}\zeta_{\chi_I}^{(1)}, \quad (\text{A11})$$

which we use in our results throughout the paper.

APPENDIX B: BARYON ISOCURVATURE

Multifield inflationary models featuring light new degrees of freedom result in isocurvature perturbations on top of the adiabatic spectrum of fluctuations predicted from the inflaton, and so generically have to contend with bounds on the allowed amount of isocurvature coming from CMB observations [31]. In the minimal (real) curvaton

scenario, the perturbations of the curvaton are initially isocurvature in nature, but become adiabatic when the curvaton decays to reheat the universe sometime after inflation. So long as the energy density of the curvaton is sufficiently large at decay, and so long as dark matter, baryons, and radiation either are among these decay products or produced later, the isocurvature bounds are trivially avoided [4,6,32].

Because we consider here a complex curvaton, however, we must account for potential isocurvature. Provided the curvaton comes to dominate the energy density prior to decay, the fluctuations in the radial direction δr will establish the adiabatic mode after decay. Fluctuations in the angular direction $\delta\theta$ are in general orthogonal, and so constitute isocurvature. These angular fluctuations lead to isocurvature fluctuations in the baryon density, since the baryon number is necessarily produced via dynamics *before* curvaton decay. To avoid this possibility and the stringent constraints on isocurvature, one either needs to tune the initial conditions such that this is vanishing or compensate the baryon isocurvature with equal and opposite isocurvature in the dark matter. We will describe the first option, which requires less fine-tuning overall.

To be more concrete, in our scenario, a complex curvaton $\chi = re^{i\theta}/\sqrt{2}$ dominates the universe as a non-relativistic species before decaying. At the time of its decay, the field is nonrelativistic and so has the approximate energy density $\rho_\chi \approx m_\chi^2 r^2$, with only subleading corrections from the quartic terms that depend on λ and λ' . The gauge invariant curvature perturbation can then be written

$$\zeta = \zeta_\chi = -H \frac{\delta\rho_\chi}{\dot{\rho}_\chi} = \frac{1}{3} \frac{\delta\rho_\chi}{\rho_\chi} = \frac{2}{3} \frac{\delta r}{r}, \quad (\text{B1})$$

where we have used the equation of motion, $\dot{\rho}_\chi = -3H\rho_\chi$, to simplify the denominator and neglected the baryon-violating contribution to the energy density because $\lambda'\chi^4 \ll m^2 r^2$ at decay. After χ decay, $\rho_\chi \approx \rho_R$ and $\dot{\rho}_\chi \approx \dot{\rho}_R$, so the photon perturbation is

$$\zeta_\gamma = \zeta_\chi = \frac{2}{3} \frac{\delta r}{r}. \quad (\text{B2})$$

Since the baryon number is created before χ decays and satisfies $n_B = r^2 \dot{\theta}$, the perturbation can be written

$$\delta n_B = 2r\dot{\theta}\delta r + r^2\delta\dot{\theta}, \quad (\text{B3})$$

and $\rho_B = m_p n_B$, which we can write as the fractional fluctuation

$$\frac{\delta\rho_B}{\rho_B} = \frac{2\delta r}{r} + \frac{\delta\dot{\theta}}{\dot{\theta}}. \quad (\text{B4})$$

We can then relate the fluctuation in the angular derivative to the field fluctuation $\delta\theta$ using the equation of motion, $\dot{n}_B \approx V'$, integrated over a Hubble time,

$$\dot{\theta} \approx \frac{V'}{r^2 H}, \quad \delta\dot{\theta} \approx \frac{V''}{r^2 H} \delta\theta, \quad (\text{B5})$$

so we can write $\delta\dot{\theta}/\dot{\theta} = (V''/V')\delta\theta$, where $' \equiv \partial_\theta$. Thus, using the definition of the baryon curvature perturbation

$$\zeta_B = -H \frac{\delta\rho_B}{\dot{\rho}_B} = \frac{1}{3} \frac{\delta\rho_B}{\rho_B} = \frac{1}{3} \left(\frac{2\delta r}{r} + \frac{V''}{V'} \delta\theta \right), \quad (\text{B6})$$

so combining this result with Eq. (B2), the baryon-photon isocurvature can be written

$$\mathcal{S}_{B\gamma} = 3(\zeta_B - \zeta_\gamma) = \frac{1}{3} \left(\frac{V''}{V'} \right) \delta\theta, \quad (\text{B7})$$

which only depends on the angular fluctuation since the quantity (V''/V') depends only on the phase θ .

The fluctuations in the angular mode are set during inflation as

$$\sqrt{\langle \delta\theta^2 \rangle} \simeq \frac{\gamma}{2\pi} \frac{H_i}{r_i}, \quad (\text{B8})$$

where γ is an $\mathcal{O}(1)$ factor. Provided that the dark matter and radiation arise from the decays of the curvaton radial mode, $\delta n_B/n_B \simeq \delta\eta/\eta$ sets the entropy perturbation at early times. This is then conserved until horizon reentry, at which point it contributes to the temperature perturbation as [12]

$$\left. \frac{\delta T}{T} \right|_{\text{iso}} = -\frac{6}{15} \frac{\Omega_B}{\Omega_m} \frac{\delta\eta}{\eta}. \quad (\text{B9})$$

Substituting in the form of $\delta\eta/\eta$ above,

$$\left\langle \left(\frac{\delta T}{T} \right)^2 \right\rangle_{\text{iso}} \simeq \frac{144\gamma^2}{225\pi^2} \frac{\Omega_B^2}{\Omega_m^2} \frac{H_i^2}{r_i^2} \cot^2(4\theta_i). \quad (\text{B10})$$

Meanwhile, during inflation the radial mode of the curvaton also picks up fluctuations of amplitude $\delta r \sim H_i/2\pi$ such that the dimensionless curvature power spectrum is schematically $\mathcal{P}_\zeta \sim \kappa^2 (H_i/2\pi r_*)^2$, with r_* the value at horizon exit. The adiabatic contribution to the temperature anisotropy is [33]

$$\left\langle \left(\frac{\delta T}{T} \right)^2 \right\rangle_{\text{adi}} \sim \frac{1}{25} \mathcal{P}_\zeta, \quad (\text{B11})$$

with the factor $1/25$ coming from the Sachs-Wolfe effect. We parametrize the fractional contribution of isocurvature to the total CMB temperature anisotropy by

$$\alpha_{II} \equiv \frac{\langle (\delta T/T)^2 \rangle_{\text{iso}}}{\langle (\delta T/T)^2 \rangle_{\text{adi}} + \langle (\delta T/T)^2 \rangle_{\text{iso}}}. \quad (\text{B12})$$

At 95% CL, *Planck* TT, TE, EE + lowE + lensing data constrain [31]

$$\alpha_{II} < 1.55 \times 10^{-2}. \quad (\text{B13})$$

We have checked that all of our benchmark points of Table I satisfy this condition.

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