Minimal approach to baryogenesis via the Affleck-Dine mechanism and inflaton mass terms

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We present a minimal approach to the generation of the baryon (B) asymmetry of the Universe, in which the asymmetry is generated in a complex inflaton condensate via B -violating quadratic inflaton potential terms and the Affleck-Dine mechanism. We show that the B -violating quadratic mass terms create an oscillating asymmetry in the complex inflaton condensate at late times. The final asymmetry transferred to the Standard Model sector at reheating is naturally reduced to the magnitude of the observed B asymmetry by the effect of averaging over the B oscillations. This approach to baryogenesis can easily be realized in a wide range of inflation models.

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I. INTRODUCTION

The Affleck-Dine (AD) mechanism [1,2] provides a remarkably simple and elegant explanation for the baryon (B) asymmetry of the Universe. A complex scalar with a $U(1)$ global symmetry, corresponding to conserved baryon number, evolves into a coherently oscillating condensate. B violating terms in the potential act on the field, pushing it into an elliptical trajectory in the complex field plane, which is equivalent to a B asymmetry in the scalar field.

The conventional AD mechanism is based on a complex scalar field Φ with a potential which at late times is dominated by a $|\Phi|^2$ mass term. Higher-order operators that violate baryon number cause the real (ϕ_1) and imaginary (ϕ_2) parts of Φ to evolve differently when the $|\Phi|^2$ term comes to dominate the potential, pushing the trajectory into an ellipse in the (ϕ_1, ϕ_2) plane. The higher-order operators become less important as the magnitude of Φ decreases due to expansion, effectively switching off the B violation and leaving a conserved baryon asymmetry in the complex field at late times.

Here we will present a new and unconventional implementation of AD baryogenesis, in which B -violating Φ^2 terms in the potential of a complex inflaton Φ generate the asymmetry.¹ (Applications of the conventional AD

mechanism to a complex inflaton have been considered in [3–7].) We will show that these terms generate a B asymmetry in the Φ condensate which oscillates about zero. When the condensate asymmetry is transferred to the Standard Model (SM) sector by Φ decay, a net asymmetry is left in the SM sector. The oscillating baryon asymmetry initially generated in the Φ condensate is typically much larger than that required to explain the observed baryon-to-entropy ratio. The asymmetry transferred to the SM is subsequently suppressed by averaging over the condensate asymmetry oscillations, reducing the asymmetry to the observed value.² The resulting model is dynamically quite different from existing inflaton based AD baryogenesis models, with the inflaton asymmetry being generated at late times during inflaton oscillations rather than during or shortly after inflation.

The paper is organized as follows: In Sec. II we discuss the generation of the asymmetry via quadratic B -violating potential terms. In Sec. III we consider possible washout of the asymmetry via inflaton exchange operators. In Sec. IV we discuss the validity of the classical calculation of the asymmetry. In Sec. V we present our conclusions.

II. AD BARYOGENESIS VIA QUADRATIC POTENTIAL TERMS

We will consider a renormalizable B symmetric inflaton potential together with B -violating Φ^2 terms,

²AD baryogenesis via mass terms has previously been considered in the context of a different class of model in [8]. The analysis of [8] assumes that the averaging over of asymmetry oscillations washes out the final asymmetry. We will show that although the asymmetry is suppressed, it is significantly nonzero. This suppression plays an important role in the model described here.

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¹The same model can also be used to generate a lepton asymmetry which is subsequently processed via sphalerons into a baryon asymmetry.

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$$V(\Phi) = m_\Phi^2 |\Phi|^2 + \lambda_\Phi |\Phi|^4 - (A\Phi^2 + \text{H.c.}), \quad (1)$$

where A is real and positive. Such potentials are naturally compatible with inflation models which are nonminimally coupled to gravity [9]. More generally, they represent the leading-order terms of an inflaton potential during post-inflation evolution.³ Φ is initially coherently oscillating, with the potential dominated by the $|\Phi|^4$ term and with no asymmetry in the field. In terms of $\Phi = (\phi_1 + i\phi_2)/\sqrt{2}$, the potential becomes

$$V(\Phi) = \frac{1}{2}(m_\Phi^2 - 2A)\phi_1^2 + \frac{1}{2}(m_\Phi^2 + 2A)\phi_2^2 + \frac{\lambda_\Phi}{4}(\phi_1^2 + \phi_2^2)^2. \quad (2)$$

The field equations are

$$\ddot{\phi}_1 + 3H\dot{\phi}_1 = -m_1^2\phi_1 - \lambda_\Phi(\phi_1^2 + \phi_2^2)\phi_1, \quad (3)$$

and

$$\ddot{\phi}_2 + 3H\dot{\phi}_2 = -m_2^2\phi_2 - \lambda_\Phi(\phi_1^2 + \phi_2^2)\phi_2, \quad (4)$$

where

$$m_1^2 = m_\Phi^2 - 2A; \quad m_2^2 = m_\Phi^2 + 2A. \quad (5)$$

In the limit $\lambda_\Phi \rightarrow 0$ the equations for ϕ_1 and ϕ_2 are decoupled from each other, with coherently oscillating solutions for ϕ_1 and ϕ_2 which have angular frequencies m_1 and m_2 , respectively.

We first derive an analytical expression for the asymmetry using a threshold approximation, which we then compare to a complete numerical solution. In the threshold approximation we consider the potential to be approximated by

$$V(\Phi) = \lambda_\Phi |\Phi|^4; \quad \phi > \phi_*, \\ V(\Phi) = m_\Phi^2 |\Phi|^2 - (A\Phi^2 + \text{H.c.}); \quad \phi < \phi_*, \quad (6)$$

where $\phi_* = m_\Phi/\sqrt{\lambda_\Phi}$ is the value of ϕ at which $V'(\phi)$ becomes dominated by the $|\Phi|^4$ term (here $\Phi = \phi e^{i\theta}/\sqrt{2}$ and we have set $A = 0$ when determining ϕ_*). The potential is initially strongly dominated by the $|\Phi|^4$ term, with $\phi_i \gg \phi_*$, and the field is initially at rest with initial values $(\phi_{1,i}, \phi_{2,i})$. Assuming rapid coherent oscillations, the field amplitude will initially evolve as $\phi \propto 1/a$ when $\phi > \phi_*$. Therefore the field amplitudes at a_* are

$$\phi_{1,*} = \left(\frac{a_i}{a_*}\right)\phi_{1,i} = \left(\frac{\phi_*}{\phi_i}\right)\phi_{1,i}; \\ \phi_{2,*} = \left(\frac{a_i}{a_*}\right)\phi_{2,i} = \left(\frac{\phi_*}{\phi_i}\right)\phi_{2,i}, \quad (7)$$

where $\phi_i = (\phi_{1,i}^2 + \phi_{2,i}^2)^{1/2}$. The field evolves purely due to the mass squared terms once $a > a_*$. We assume that $m_{1,2} \gg H$, so that we can neglect the effect of expansion on the rapid $\phi_{1,2}$ oscillations and simply factor in the effect of expansion by damping the oscillation amplitude. The solution for ϕ_1 and ϕ_2 is then

$$\phi_1 = \phi_{1,*} \left(\frac{a_*}{a}\right)^{3/2} \cos(m_1(t - t_*)); \\ \phi_2 = \phi_{2,*} \left(\frac{a_*}{a}\right)^{3/2} \cos(m_2(t - t_*)). \quad (8)$$

The baryon asymmetry in the Φ condensate is

$$n(t) = i(\Phi^\dagger \dot{\Phi} - \dot{\Phi}^\dagger \Phi) = \dot{\phi}_1 \phi_2 - \dot{\phi}_2 \phi_1. \quad (9)$$

Therefore,

$$n(t) = \phi_{1,*} \phi_{2,*} \left(\frac{a_*}{a}\right)^3 [m_2 \sin(m_2(t - t_*)) \cos(m_1(t - t_*)) \\ - m_1 \sin(m_1(t - t_*)) \cos(m_2(t - t_*))]. \quad (10)$$

We will assume that $2A \ll m_\Phi^2$. In this limit, to leading order in A/m_Φ^2 , the condensate baryon asymmetry becomes

$$n(t) = \phi_{1,*} \phi_{2,*} \left(\frac{a_*}{a}\right)^3 \left[m_\Phi \sin\left(\frac{2A(t - t_*)}{m_\Phi}\right) \right. \\ \left. + \frac{A}{m_\Phi} \sin(2m_\Phi(t - t_*)) \right]. \quad (11)$$

During averaging over the $\phi_{1,2}$ coherent field oscillations, we can consider the scale factor to be constant since $H \ll m_\Phi$. The second term in Eq. (11) then averages to zero. The condensate asymmetry at $t > t_*$, in terms of the initial field values, is then

$$n(t) = \phi_{1,i} \phi_{2,i} \left(\frac{\phi_i}{\phi_*}\right) \left(\frac{a_i}{a}\right)^3 m_\Phi \sin\left(\frac{2A(t - t_*)}{m_\Phi}\right). \quad (12)$$

Thus the baryon asymmetry in the Φ condensate oscillates about zero with period $T_{\text{asy}} = \pi m_\Phi/A$.

It is useful to define a comoving asymmetry $n_c(t) \equiv (a(t)/a_i)^3 n(t)$, which is constant when there is no production or decay of the asymmetry. For the threshold model at $t > t_*$

³While the inflaton is the natural candidate for the field responsible for reheating, we note that the model can apply to any coherently oscillating complex scalar that is responsible for reheating.

$$n_c(t) = \phi_{1,i}\phi_{2,i}\left(\frac{\phi_i}{\phi_*}\right)m_\Phi \sin\left(\frac{2A(t-t_*)}{m_\Phi}\right), \quad (13)$$

with $n_c(t) = 0$ at $t < t_*$. The Φ condensate asymmetry is assumed to transfer to a conserved SM baryon asymmetry via B -conserving Φ decays to SM particles.^{4,5} The condensate will decay away completely after a time $t_R \approx \Gamma_\Phi^{-1}$, where R denotes reheating, with continuous production of SM baryon asymmetry due to decay of the condensate asymmetry from t_* to t_R . Neglecting any reduction of the Φ field due to decays at $t < t_R$, the comoving baryon asymmetry transferred to the SM sector, which we denote by $\hat{n}_c(t)$, is

$$\hat{n}_c(t) = \int_{t_i}^t \Gamma_\Phi n_c(t) dt. \quad (14)$$

Thus the comoving baryon asymmetry transferred out of the Φ condensate as a function of t is

$$\hat{n}_c(t) = \frac{\Gamma_\Phi \phi_{1,i} \phi_{2,i} m_\Phi^2}{2A} \left(\frac{\phi_i}{\phi_*}\right) \left[1 - \cos\left(\frac{2A(t-t_*)}{m_\Phi}\right)\right]. \quad (15)$$

$\hat{n}_c(t)$ increases linearly with $t - t_*$ until $t - t_* \approx \pi m_\Phi / 4A$. On longer timescales, $\hat{n}_c(t)$ oscillates between a maximum value and zero with period T_{asy} . The maximum possible asymmetry is obtained when $A = A_{\text{max}} = \pi m_\Phi \Gamma_\Phi / 2$.

The Φ condensate decays away completely once $t - t_* \gtrsim \Gamma_\Phi^{-1}$. To take into account the B -conserving decay of the condensate asymmetry, we include in Eq. (14) an exponential decay factor,

$$\hat{n}_c(t) = \int_{t_i}^t \Gamma_\Phi n_c(t) e^{-\Gamma_\Phi(t-t_*)} dt. \quad (16)$$

The total comoving asymmetry transferred to the SM sector as $t \rightarrow \infty$ is then

$$\hat{n}_{c,\text{tot}} = \frac{\Gamma_\Phi \phi_{1,i} \phi_{2,i} m_\Phi^2}{2A} \left(\frac{\phi_i}{\phi_*}\right) \left(1 + \left(\frac{\Gamma_\Phi m_\Phi}{2A}\right)^2\right)^{-1}. \quad (17)$$

The transferred asymmetry is proportional to A until $A > \Gamma_\Phi m_\Phi / 2$, in which case $\tau_\Phi > T_{\text{asy}}$ and the transferred asymmetry decreases as A^{-1} and τ_Φ^{-1} , where $\tau_\Phi = \Gamma_\Phi^{-1}$ is the lifetime of the Φ scalars. This can be understood as due to the effect of averaging condensate oscillations over the time taken for the condensate to decay. When $\tau_\Phi \gg T_{\text{asy}}$, the

asymmetry in the condensate will undergo many oscillations from positive to negative values during the decay of the condensate. Therefore, the asymmetry produced during a positive half cycle will almost cancel against that produced during the following negative half cycle, up to the effect of the small decrease in the condensate asymmetry amplitude due to the decay of the condensate during $\Delta t \sim T_{\text{asy}}$. Therefore only a small net asymmetry is produced during each condensate oscillation cycle as compared to the case with $T_{\text{asy}} \gtrsim \tau_\Phi$, where there is no averaging over oscillations.

We first consider the case where the lifetime of Φ is much longer than T_{asy} , such that $2A/m_\Phi \Gamma_\Phi \gg 1$. $\hat{n}_{c,\text{tot}}$ can then be expressed as

$$\hat{n}_{c,\text{tot}} = \frac{\Gamma_\Phi \phi_i^2 m_\Phi^2 \sin(2\theta)}{4A} \left(\frac{\phi_i}{\phi_*}\right), \quad (18)$$

where θ is the initial phase of Φ . The total baryon asymmetry transferred to the SM, $\hat{n}_{\text{tot}}$, is then

$$\hat{n}_{\text{tot}} = \left(\frac{a_i}{a_R}\right)^3 \hat{n}_{c,\text{tot}} = \frac{3M_{\text{Pl}}^2 \Gamma_\Phi^3 \sin(2\theta)}{2A}, \quad (19)$$

where we have used $a \propto H^{-2/3}$ when $a > a_R$ and $a \propto 1/\phi$ when $a < a_R$ to obtain the final expression. This can also be expressed in terms of the baryon-to-entropy ratio, n_B/s . Using $s = 4k_T^2 T^3$ and $\Gamma_\Phi = H_R = k_{T_R} T_R^2 / M_{\text{Pl}}$, where T_R is the reheating temperature and $k_T = (\pi^2 g(T)/90)^{1/2}$, the baryon-to-entropy ratio is

$$\begin{aligned} \frac{n_B}{s} &\equiv \frac{\hat{n}_{\text{tot}}}{s} = \frac{3 k_{T_R} T_R^3 \sin(2\theta)}{8 A M_{\text{Pl}}} \\ &= 5.2 \times 10^{-21} \frac{m_\Phi^2}{A} \left(\frac{T_R}{10^8 \text{ GeV}}\right)^3 \left(\frac{10^{13} \text{ GeV}}{m_\Phi}\right)^2 \sin(2\theta), \end{aligned} \quad (20)$$

where we have normalized the expression to some representative values⁶ of T_R and m_Φ . The observed baryon-to-entropy ratio is $(n_B/s)_{\text{obs}} = 0.861 \pm 0.005 \times 10^{-10}$. In order to account for the observed asymmetry, we require that

$$\frac{A^{1/2}}{m_\Phi} = 7.8 \times 10^{-6} \sin^{1/2}(2\theta) \left(\frac{10^{13} \text{ GeV}}{m_\Phi}\right) \left(\frac{T_R}{10^8 \text{ GeV}}\right)^{3/2}. \quad (21)$$

The maximum possible asymmetry, which corresponds to $A = \Gamma_\Phi m_\Phi / 2$ in Eq. (17), is

⁶ $T_R = 10^8 \text{ GeV}$ is within the range of reheating temperatures that may be detectable in the spectrum of primordial gravitational waves [12].

⁴A specific implementation of the model to baryogenesis from AD leptogenesis via a decaying inflaton will be presented in a future work [10]. Here we focus on the general features of inflaton mass term AD baryogenesis.

⁵It is also possible for the inflaton to decay via gravity mediated modes [11]. The importance of this process will depend upon the coupling of the inflaton to the Ricci curvature in a given inflation model.

FIG. 1. The condensate asymmetry (left) and transferred asymmetry (right) for the case $m_\Phi = 10^{16}$ GeV, $\lambda_\Phi = 0.1$ and $T_R = 10^8$ GeV. The threshold asymmetry and the numerical results for $A^{1/2}/m_\Phi = 0.001, 0.005, 0.007, 0.01$ and 0.05 are shown. [The numerical result for $A^{1/2}/m_\Phi = 0.001$ coincides with the threshold result, in agreement with Eq. (25).]

$$\begin{aligned} \frac{n_{B,\max}}{s} &= \frac{3T_R \sin(2\theta)}{8m_\Phi} \\ &= 3.8 \times 10^{-6} \left(\frac{T_R}{10^8 \text{ GeV}} \right) \left(\frac{10^{13} \text{ GeV}}{m_\Phi} \right) \sin(2\theta). \end{aligned} \quad (22)$$

This can easily be much larger than the observed baryon asymmetry. Therefore, the suppression of the asymmetry by averaging over oscillations plays an important role in this model.

In the case where Φ decays before any condensate asymmetry oscillations can occur, corresponding to $\Gamma_\Phi m_\Phi / 2A \gg 1$ in Eq. (17), the total transferred asymmetry obtains an additional factor $(2A/\Gamma_\Phi m_\Phi)^2$ compared to Eq. (20). Therefore,

$$\frac{n_B}{s} = \frac{3 A M_{\text{Pl}} \sin(2\theta)}{2 k_{T_R} T_R m_\Phi^2} \quad (23)$$

and we find that the required value of $A^{1/2}/m_\Phi$ is

$$\frac{A^{1/2}}{m_\Phi} = 8.9 \times 10^{-11} \left(\frac{T_R}{10^8 \text{ GeV}} \right)^{1/2} \left(\frac{1}{\sin(2\theta)} \right)^{1/2}. \quad (24)$$

This is typically much smaller than in the case with asymmetry oscillations, due to the lack of additional suppression of the baryon asymmetry from averaging over condensate oscillations.

The threshold asymmetry is a good approximation if the B -violating mass terms do not cause the field to significantly evolve until the potential is $|\Phi|^2$ dominated. The condition for this to be true, which we have confirmed in our numerical solutions, is that the mass of the angular field perturbations about the minimum of the potential as a function of θ , $m_{\delta\theta} = 2A^{1/2}$, is less than H when $\phi = \phi_*$. This is satisfied if

$$\begin{aligned} \frac{A^{1/2}}{m_\Phi} &\lesssim \frac{A_{\text{th}}^{1/2}}{m_\Phi} = \frac{m_\Phi}{4\sqrt{\lambda_\Phi} M_{\text{Pl}}} \\ &\equiv 1.0 \times 10^{-6} \lambda_\Phi^{-1/2} \left(\frac{m_\Phi}{10^{13} \text{ GeV}} \right). \end{aligned} \quad (25)$$

We finally compare the threshold approximation to the complete numerical solution for the case $\Gamma_\Phi(t - t_*) \ll 1$.⁷ As an example, we show in Fig. 1 the numerical results for the case $m_\Phi = 10^{16}$ GeV and $\lambda_\Phi = 0.1$ for a range of values of $A^{1/2}/m_\Phi$. The analytical approximation in the figure shown on the left is given by Eq. (13) and in the figure shown on the right by Eq. (15), with Γ_Φ corresponding to $T_R = 10^8$ GeV. For this case, the upper limit for the threshold approximation to be valid is $A_{\text{th}}^{1/2}/m_\Phi \approx 3 \times 10^{-3}$. We find that the threshold approximation is in perfect agreement with the numerical solution for both the condensate and transferred asymmetries when Eq. (25) is satisfied. For larger $A^{1/2}/m_\Phi$, the evolution during the $|\Phi|^4$ dominated era modifies the asymmetries. The amplitude of the transferred asymmetry $A\hat{n}_c$ rapidly decreases with increasing $A > A_{\text{th}}$ down to an approximately constant value, which is suppressed relative to the threshold value of $A\hat{n}_c$ by a factor that numerically is approximately $m_\Phi/10^{17}$ GeV. The transferred asymmetry $A\hat{n}_c$ oscillates between zero and a maximum when the threshold approximation is valid, but for larger $A^{1/2}/m_\Phi$ it oscillates about zero. However, since the transferred asymmetry is the total asymmetry transferred to the SM sector as a function of time after averaging over condensate asymmetry oscillations, the oscillation of the transferred asymmetry about

⁷Further details of the numerical analysis will be presented in [10].

zero has no impact on the typical magnitude of the baryon asymmetry transferred to the SM.

III. BARYON WASHOUT DUE TO INFLATON EXCHANGE

In application to a specific model, the possible washout of the asymmetry must be considered. The interaction which allows the decay of the inflaton will generally result in a B -violating operator via Φ exchange. Dimensionally, the rate of B -violating scattering processes at reheating due to ϕ_1 and ϕ_2 exchange is

$$\Gamma_{\Delta B} \sim \frac{\lambda_\psi^2 A^2 T_R^5}{m_\Phi^8}, \quad (26)$$

where λ_ψ is the coupling responsible for Φ decay and A is necessary in the scattering amplitude in order to have B violation. Washout due to Φ exchange will be negligible if $\Gamma_{\Delta B} < H(T_R)$, which is satisfied if

$$\begin{aligned} \lambda_\psi &\lesssim \frac{m_\Phi^4}{M_{\text{Pl}}^{1/2} T_R^{3/2} A} \\ &= 6 \times 10^4 \left(\frac{m_\Phi^2}{A} \right) \left(\frac{m_\Phi}{10^{13} \text{ GeV}} \right)^2 \left(\frac{10^8 \text{ GeV}}{T_R} \right)^{3/2}. \end{aligned} \quad (27)$$

The inflaton decay rate is $\Gamma_\Phi \approx \lambda_\psi^2 m_\Phi / 4\pi$, therefore the reheating temperature from $H(T_R) = \Gamma_\Phi$ is $T_R \approx \lambda_\psi (m_\Phi M_{\text{Pl}})^{1/2}$. Thus, Eq. (27) is satisfied if

$$T_R \lesssim \left(\frac{m_\Phi^2}{A} \right)^{2/5} m_\Phi, \quad (28)$$

where $A < m_\Phi^2$. Therefore washout due to B -violating Φ exchange is negligible if $T_R \lesssim m_\Phi$ and so it is unlikely to present a serious obstacle to this class of model. A complete analysis of washout will depend upon the specific model for the decay of the inflaton and the transfer of the baryon asymmetry.

IV. VALIDITY OF THE CLASSICAL ANALYSIS OF THE BARYON ASYMMETRY

Throughout our analysis we have assumed that classical fields can be used to calculate the baryon asymmetry. When the potential is dominated by quadratic terms, the ϕ_1 and ϕ_2 fields evolve as independent noninteracting coherently oscillating scalars. In general, a classical oscillating scalar field corresponds to a quantum coherent state in the limit where the occupation number of the state is large compared to one [13,14]. The condition for this to be true is that $\phi_i > m_\Phi$ ($i = 1, 2$). However, this is typically not satisfied at inflaton decay in the present model. Nevertheless, the classical calculation of the baryon asymmetry remains correct. This is because it is the coherent

state corresponding to the classical field that is important for AD baryogenesis.

By construction, the expectation value of the field operator $\hat{\phi}_i$ in the coherent state $|\phi_i(t)\rangle$ is equal to the classical field $\phi_{i,\text{cl}}(t)$

$$\langle \phi_i(t) | \hat{\phi}_i | \phi_i(t) \rangle = \phi_{i,\text{cl}}(t). \quad (29)$$

We have included a time dependence in the coherent state to take into account the dilution of the number density by expansion. Since the scalar fields ϕ_1 and ϕ_2 are independent fields, the coherent state of the complex field is a product of the coherent states for ϕ_1 and ϕ_2 , $|\Phi(t)\rangle = |\phi_1(t)\rangle |\phi_2(t)\rangle$. Therefore, with the baryon number density operator given by $\hat{n} = \hat{\phi}_1 \hat{\phi}_2 - \hat{\phi}_2 \hat{\phi}_1$, the expectation value of the baryon asymmetry in the coherent state is given by

$$\begin{aligned} \langle \Phi(t) | \hat{n} | \Phi(t) \rangle &= \langle \Phi(t) | \hat{\phi}_1 \hat{\phi}_2 - \hat{\phi}_2 \hat{\phi}_1 | \Phi(t) \rangle \\ &= \dot{\phi}_{1,\text{cl}} \phi_{2,\text{cl}} - \dot{\phi}_{1,\text{cl}} \phi_{1,\text{cl}} \equiv n_{\text{cl}}. \end{aligned} \quad (30)$$

Therefore the expectation value of the baryon number density operator is equal to the baryon number density n_{cl} calculated using the classical fields. When $\phi_i < m_\Phi$, the variance of the field in the coherent state will become large compared to the squared classical field. Therefore, there will be large quantum fluctuations of the fields about their expectation values and so the field cannot be considered classical. However, the correlation length of the quantum fluctuations cannot be larger than the horizon at inflaton decay. Since the volume that evolves into the presently observed Universe will be very much larger than the horizon volume at inflaton decay, the observed baryon asymmetry will be given by its spatial average value and so will equal the expectation value of the baryon asymmetry. Therefore the baryon asymmetry will equal its classical value even when $\phi_i < m_\Phi$. This shows that it is the coherent state describing the scalar field, rather than its classical nature, that is essential for AD baryogenesis.

In reaching this conclusion, we have assumed that the mean asymmetry transferred from the condensate by decay is equal to the mean asymmetry in the coherent state of the condensate, and that there is no additional washout effect due to the decay process. Condensate decay in this model occurs when the occupation number is less than one, therefore the conventional classical analysis based on production of particles due to a time-dependent classical field is no longer valid. While there is no obvious reason to expect an additional source of washout due to the decay process when the coherent state is no longer in the classical limit, this should be confirmed by a full quantum field theory analysis.

V. CONCLUSIONS

We have presented a new minimal approach to baryogenesis which is based on B -violating mass terms for the inflaton. The resulting model requires only the addition of B -violating mass terms to an existing inflaton potential and therefore can easily be realized in a wide range of inflation models. The asymmetry is generated at late times during inflaton oscillations, in contrast to existing inflaton-based AD baryogenesis models which generate the asymmetry during or shortly after inflation. The model also provides exact analytical expressions for the resulting baryon asymmetry.

In this analysis we have not addressed the question of baryon isocurvature perturbations. We note that these can

easily be controlled by including a $\Phi^4 + \Phi^{\dagger 4}$ term in the potential which is significant during inflation and becomes negligible after inflation, while leaving open the possibility of observable isocurvature perturbations. A detailed implementation of the mechanism to baryogenesis from AD leptogenesis via inflaton decay, including a discussion of isocurvature perturbations, will be presented in a future work [10]. The model also raises new questions regarding the AD mechanism in the limit where the classical approximation is no longer valid, which requires a dedicated analysis.

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