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Dark Matter and Baryon Asymmetry from the very Dawn of Universe

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Abstract. We discuss a possibility of universally producing dark matter and baryon charge at inflation. For this purpose, we introduce a complex scalar field with the mass exceeding the Hubble rate during the last e-folds of inflation. We assume that the phase of the complex scalar is linearly coupled to the inflaton. This interaction explicitly breaking $U(1)$-symmetry leads to the production of a non-zero Noether charge. The latter serves as a source of dark matter abundance, or baryon asymmetry, if the complex scalar carries the baryon charge.

1 Introduction and Summary

Inflation is the hypothetical stage in the early Universe introduced to address hot Big Bang problems: horizon, flatness, entropy problems, the overclosure of the Universe by heavy relics [1–5]. Furthermore, inflaton perturbations serve as the source of initial inhomogeneities [6], which are necessary for the existence of the observed large scale structures. On the flipside, inflation is too efficient in a sense that any pre-existing dark matter (DM) or baryon asymmetry (BA) are washed out by the end of the (quasi)-de Sitter expansion. Of course, there is no shortage of DM or BA generation mechanism starting from preheating. However, those mechanisms are typically different by nature and operate at different times. This leads to a problem of coincidence of the observed DM and baryon abundance.

In the present manuscript, closely following Ref. [7] we discuss the counterexample to the statement that no DM and BA is produced during inflation. The key idea is to linearly couple the field responsible for the generation of DM or the baryon charge, to a function of the inflaton. During inflation this coupling induces an almost constant force dragging the field to the non-zero value. The simplest realization of the mechanism involves a complex scalar field Ψ with the phase φ coupled to the trace of the inflaton energy-momentum tensor. That interaction explicitly breaks the $U(1)$-symmetry and thus leads to the production of the non-zero Noether charge density. If the field Ψ carries the non-zero baryon charge, the Noether charge density can be reprocessed into BA *a la* the Affleck–Dine mechanism [8]. Alternatively, if the field Ψ is sterile, the Noether charge density is associated with the number density of DM particles. See Sec. 3.

The mechanism considered here works for the super-heavy fields Ψ with the masses $M \gtrsim H$. Another motivation to introduce so heavy fields is the presence of isocurvature modes,

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which are naturally large for light fields during inflation. On the other hand, provided that $M \gtrsim H$, the isocurvature modes are exponentially suppressed, and the resulting perturbations are adiabatic (Sec. 4)—in a comfortable agreement with the Planck data [9]. Note, however, that super-heavy fields are required, if the complex scalar is minimally coupled to gravity. Instead, if there is a non-minimal coupling to gravity, the constraint on the mass M can be avoided. We consider this option in Sec. 4.

2 Generation of dark matter/baryon asymmetry during inflation

The full action of the model we are interested in, is given by

$$S = S_{EH} + S_{\Psi} + S_{int} + S_{infl} , \tag{1}$$

Here S_{EH} is the Einstein–Hilbert action,

$$S_{EH} = \frac{M_{Pl}^2}{16\pi} \int d^4x \sqrt{-g} R ,$$

where M_{Pl} is the Planck mass; S_{Ψ} is the action of the complex scalar field,

$$S_{\Psi} = \int d^4x \sqrt{-g} \left[\frac{1}{2} |\partial_\mu \Psi|^2 - V(|\Psi|) \right] , \tag{2}$$

where $V(|\Psi|)$ is the potential. For the purpose of modeling DM, we set the self-interactions to zero, so that

$$V(|\Psi|) = \frac{M^2 |\Psi|^2}{2} ,$$

where M is the mass of the complex scalar field. In Eq. (1), S_{inf} is the action for the inflaton, ϕ, which we assume to be a canonical real scalar field minimally coupled to gravity and slowly rolling the potential U, i.e.,

$$S_{inf} = \int d^4x \sqrt{-g} \left(\frac{1}{2} (\partial_\mu \phi)^2 - U(\phi) \right) .$$

We assume the following interaction between the inflaton and the phase φ of the complex scalar field $\Psi = \lambda e^{i\varphi}$,

$$S_{int} = \int d^4x \sqrt{-g} \cdot \beta \cdot \varphi \cdot T_{infl} . \tag{3}$$

Here T_{infl} is the inflaton energy-momentum tensor, $T_{infl} \approx 4U$.

For $\beta = 0$ (no interaction with the inflaton) the model (2) possesses a global $U(1)$-symmetry. Hence, there is a Noether current $J_\mu \equiv \frac{\delta S}{\delta \partial^\mu \varphi} = \lambda^2 \partial_\mu \varphi$, which is conserved,

$$\nabla_\mu J^\mu = \frac{1}{\sqrt{-g}} \partial_\mu \left(\sqrt{-g} J^\mu \right) = 0 ,$$

Consequently, the Noether charge density

$$Q \equiv J_0 = \lambda^2 \dot{\varphi} , \tag{4}$$

is conserved in the physical volume. This means that $Q \propto 1/a^3$. Thus, for $\beta = 0$, the Noether charge density relaxes to zero exponentially fast in the inflationary Universe. No DM or BA is generated in that case.

Now we switch to the situation of interest, when $\beta \neq 0$. In that case, the $U(1)$-symmetry is explicitly violated. This leads to the production of the Noether charge density,

$$\frac{1}{a^3}\frac{d}{dt}\left(Qa^3\right) = \beta T_{infl} \ . \tag{5}$$

Integrating the latter, we obtain

$$Q(t) \approx \frac{4\beta}{a^3(t)}\int_{t_{in}}^{t} U(t')a^3(t')dt' + \frac{C}{a^3(t)} \ . \tag{6}$$

The second term on the r.h.s. governed by the integration constant C is what we would have in the absence of the coupling to the inflaton. This term decreases exponentially fast in the expanding Universe. On the other hand, the first term on the r.h.s. gives the contribution, which remains constant during inflation. With the logarithmic accuracy, we have

$$Q \simeq \frac{4\beta U}{3H} \ . \tag{7}$$

We see that the Noether charge density generated is constant in the exact de Sitter space-time approximation. Its small variation with time is due to non-zero slow roll parameters.

Next we vary the action (1) with respect to the amplitude λ,

$$\Box\lambda - \lambda\dot{\varphi}^2 + M^2\lambda = 0 \ . \tag{8}$$

We use Eq. (4) to express $\dot{\varphi}$, and then substitute it back into Eq. (8),

$$\ddot{\lambda} + 3H\dot{\lambda} - \frac{Q^2}{\lambda^3} + M^2\lambda = 0 \ .$$

Hence, the field λ lives in the effective potential

$$V_{eff} = \frac{M^2\lambda^2}{2} + \frac{Q^2}{2\lambda^2} \ . \tag{9}$$

The minimum of the potential (9) is located at

$$\bar{\lambda} = \sqrt{\frac{Q}{M}} \simeq \sqrt{\frac{4\beta U(t)}{3MH(t)}} \ , \tag{10}$$

where we made use of the estimate (7).

Now let us concretize the relevant range of the mass M values. For small masses $M \ll H$, the amplitude λ is in the slow roll regime during inflation and may never reach the minimum (10). We are interested in the opposite regime involving very large masses

$$M \gtrsim H \ .$$

In that case, the amplitude λ rolls fast the slope of the effective potential (9) and relaxes to its minimum (10) within a few Hubble times. See Fig. 1.

Given that $M \gtrsim H$, right after inflation the complex scalar behaves as the pressureless fluid: its energy density decays as $\rho_{DM}(t) = M^2\lambda^2(t) \propto 1/a^3(t)$. At the onset of oscillations, which start after the end of inflation at the time $t = t_{end}$, the amplitude λ is given by Eq. (10), and hence the DM energy density is estimated by

$$\rho_{DM}(t_{end}) \simeq M^2\bar{\lambda}^2 \simeq MQ(t_{end}) \ . \tag{11}$$

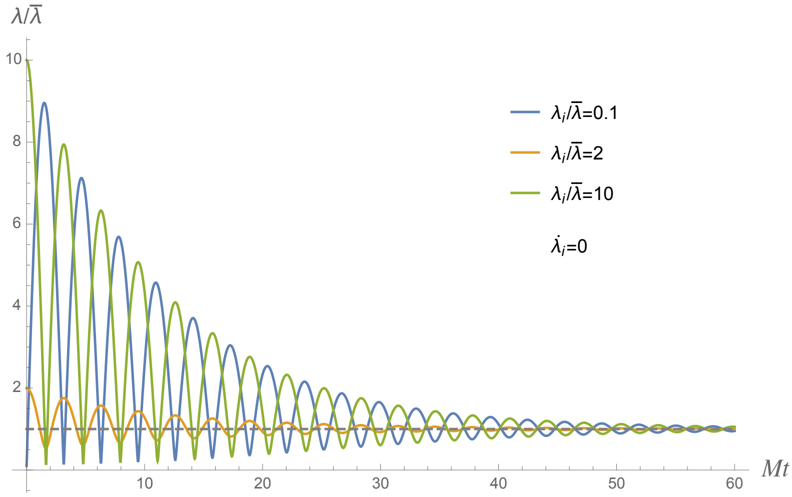

Figure 1. The amplitude λ relaxes to the minimum point $\bar{\lambda}$ (dashed line) of the effective potential (9) within a few Hubble times, independently of the initial value. We have chosen $\frac{M}{H} = 20$, so that the time range shown corresponds to 3 Hubble times, and assumed the exact de Sitter space-time for simplicity. The figure is taken from Ref. [7].

We see that it is fully determined by the Noether charge density and the mass M of the scalar, and is independent of the pre-inflationary value of the field Ψ.

Our further discussion forks depending on the nature of the field Ψ. First assume that it is sterile, and thus serves as a DM candidate. In that case, the energy density (11) should correspond to the observed abundance of DM. Namely,

$$\rho_{DM}(t_{eq}) \simeq \frac{\beta M U(t_{end})}{H(t_{end})} \cdot \frac{a_{end}^3}{a_{eq}^3} \simeq \rho_{rad}(t_{eq}) \; ,$$

where the subscript $'eq'$ denotes the equality between the matter and the radiation. Assume that the radiation-dominated stage takes place immediately after inflation, so that $\rho_{rad}(t_{end}) \simeq U(t_{end})$. Then one can write

$$\rho_{rad}(t_{eq}) \simeq \rho_{rad}(t_{end}) \cdot \frac{a_{end}^4}{a_{eq}^4} \simeq U(t_{end}) \cdot \frac{a_{end}^4}{a_{eq}^4} \; .$$

The constant β is then estimated by

$$\beta \simeq \frac{H(t_{end})}{M} \cdot \frac{a_{end}}{a_{eq}} \simeq \frac{H(t_{end})}{M} \cdot \frac{T_{eq}}{T_{end}} \; .$$

The Hubble rate is estimated by $H(t_{end}) \sim \frac{T_{end}^2}{M_{Pl}}$. We get for the coupling constant β,

$$\beta \simeq \frac{T_{end}}{M_{Pl}} \cdot \frac{T_{eq}}{M} \; . \tag{12}$$

For the choice $T_{end} \sim 10^{16}$ GeV (the maximal possible temperature in the post-inflationary Universe) and $M \sim H \sim 10^{-5} M_{Pl}$, one obtains $\beta \sim 10^{-26}$.

Alternatively, assigning the non-zero baryonic charge to the field Ψ, one can generate BA Δ_B from its decays into quarks,

$$\Delta_B \simeq \frac{Q}{s} \ ,$$

where s is the entropy density. Compared to the DM case, production of BA generically requires much larger coupling constants β. To estimate the constant β, we again assume the immediate conversion of the inflaton energy into radiation. We get

$$\Delta_B \simeq \frac{\beta U(t_{end}) T_{end}}{H(t_{end})\rho_{rad}(t_{end})} \simeq \frac{\beta T_{end}}{H(t_{end})} \simeq \beta \cdot \sqrt{\frac{M_{Pl}}{H(t_{end})}} \ .$$

Here we use that $s = \frac{4\rho_{rad}}{3T}$ and $T \simeq \sqrt{M_{Pl}H}$. The observed BA of the Universe is $\Delta_B = 0.87 \cdot 10^{-10}$. We obtain

$$\beta \simeq 10^{-10}\sqrt{\frac{H(t_{end})}{M_{Pl}}} \ . \tag{13}$$

Assuming the high scale inflation with the Hubble rate $H \simeq 10^{14}$ GeV, we end up with the coupling constant β as large as $\beta \simeq 10^{-12}$.

3 Perturbations

Now let us discuss complex field perturbations in the inflationary Universe. As usual, we split the perturbations into adiabatic and isocurvature ones. The former are sourced by the inflaton ϕ fluctuations and obey the standard relation,

$$\frac{\delta\phi}{\dot{\phi}} = \frac{\delta\lambda_{ad}}{\dot{\lambda}} = \frac{\delta\varphi_{ad}}{\dot{\varphi}} \ , \tag{14}$$

c.f. Ref. [10]. In Ref. [7] we have checked that this relation holds beyond the slow roll approximation, independently of the choice of the inflaton potential U and for the arbitrary mass M of the complex field.

Isocurvature perturbations are the ones, which the complex field has on its own. Namely, we switch off metric fluctuations and neglect inflaton perturbations (this is possible in the spectator approximation). The relevant system of equations describing the complex field perturbations is given by

$$\delta\ddot{\lambda}_{iso} + 3H\delta\dot{\lambda}_{iso} - \frac{1}{a^2}\partial_i\partial_i\delta\lambda_{iso} - \delta\lambda_{iso}\dot{\varphi}^2 - 2\lambda\dot{\varphi}\delta\dot{\varphi}_{iso} + M^2\delta\lambda_{iso} = 0 \ , \tag{15}$$

and

$$\frac{1}{a^3}\frac{d}{dt}\Big(a^3\lambda^2\delta\dot{\varphi}_{iso}\Big) - \frac{1}{a^2}\lambda^2\partial_i\partial_i\delta\varphi_{iso} + \frac{1}{a^3}\frac{d}{dt}\Big(a^3\dot{\varphi}\delta\lambda^2_{iso}\Big) = 0 \ . \tag{16}$$

One can neglect spatial derivatives behind the horizon. Then, Eq. (16) is simplified to

$$\frac{d}{dt}\left(a^3\delta[\lambda^2\dot{\varphi}]_{iso}\right) = 0 \ .$$

Integrating this out, we obtain

$$\delta(\lambda^2\dot{\varphi})_{iso} = \frac{C}{a^3} \ ,$$

where C is the constant of integration. We see that the isocurvature perturbation of the Noether charge density $\lambda^2\dot{\varphi}$ decays fast, and one gets

$$\frac{\delta\dot{\varphi}_{iso}}{\dot{\varphi}} = -\frac{2\delta\lambda_{iso}}{\lambda} \, . \tag{17}$$

Expressing the perturbation $\delta\dot{\varphi}$ from the latter and plugging into Eq. (15), we obtain

$$\delta\ddot{\lambda}_{iso} + 3H\delta\dot{\lambda}_{iso} + M^2_{eff}\delta\lambda_{iso} = 0 \, ,$$

where

$$M^2_{eff} = M^2 + 3\dot{\varphi}^2 \, . \tag{18}$$

We are interested in the masses $M \gtrsim H$ and hence $M_{eff} \gtrsim H$. For those M_{eff}, the perturbations $\delta\lambda_{iso}$ follow rapid oscillations with the amplitude, which drops to zero exponentially fast. To the contrast, the adiabatic perturbations $\delta\lambda_{ad}$ remain constant behind the horizon. Thus, for $M \gtrsim H$ we end up with the adiabatic perturbations of the complex field by the end of inflation. Note that the results of the present section are generic and stay valid for fairly arbitrary steep potentials V [11].

4 Extensions

We could see in the previous sections that generating DM requires a tiny coupling constant β. We do not know, if this is a technically natural assumption or not. The issue can be avoided in the slight modifications of the base scenario considered in the manuscript.

First, let us modify the model by introducing the non-minimal coupling of the field Ψ to gravity,

$$S_{non-min} = \int d^4x \sqrt{-g} \left(\frac{\xi\lambda^2}{2} R \right) \, ,$$

where ξ is some dimensionless parameter. During inflation, the Ricci scalar $R \approx -12H^2$ is nearly constant. Thus, the non-minimal coupling to gravity mimics the mass term of the complex field. The equation of motion (8) for the amplitude λ now takes the form,

$$\ddot{\lambda} - \lambda\dot{\varphi}^2 + 3H\dot{\lambda} + M^2_{eff}\lambda = 0 \, ,$$

where the effective mass squared M^2_{eff} is the sum of two contributions: one coming from the standard mass term and another due to the non-minimal coupling, i.e.,

$$M^2_{eff} \approx M^2 + 12\xi H^2 \, . \tag{19}$$

For $\xi \gtrsim 1$, one has $M_{eff} \gtrsim H$. Therefore, the discussion of the previous sections is applied, even though the intrinsic mass M can be very small, i.e., $M \ll H$. Namely, the amplitude λ relaxes to the minimum of the effective potential (9), where one should just replace M by M_{eff}, within a few Hubble times. Furthermore, isocurvature perturbations decay fast in the super-horizon regime. On the other hand, the value of the coupling constant β is fixed from the post-inflationary evolution of the complex scalar, when $R \to 0$ and hence $M_{eff} \to M$.

There is another way to relax the constraint on the coupling constant β. In the manuscript, we assumed that the field Ψ is stable in the DM case. Instead, one may consider the situation, when it decays into some light stable particles with the mass $m \ll M$, e.g., sterile neutrinos, at some point. In that case, the coupling constant β is larger by the factor $\frac{M}{m}$ compared to that given in Eq. (12).

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