



Letter

Peccei-Quinn genesis

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ABSTRACT

We propose aogenesis mechanism that unifies the origin of QCD axion dark matter and the baryon asymmetry of the Universe in the framework of Peccei-Quinn pole inflation. The model integrates the Peccei-Quinn symmetry with the seesaw mechanism for neutrino masses. This allows for spontaneous leptogenesis, which generates the required $B - L$ asymmetry around the seesaw scale. The necessary initial axion kinetic misalignment is naturally sourced by a PQ field driving pole inflation. Analysis within the KSVZ axion model demonstrates that achieving simultaneous correct DM abundance and baryon asymmetry limits the axion decay constant to be smaller than about 10^9 GeV. This framework offers a unified solution to four fundamental problems: the strong CP problem, neutrino mass, matter-antimatter asymmetry, and inflation.

1. Introduction

The QCD axion is a Goldstone boson of a spontaneously broken Peccei-Quinn (PQ) symmetry $U(1)_{PQ}$ that has a color anomaly. It has been introduced to solve the strong CP problem [1–4] as its vacuum expectation value, and thus the effective θ term settles down to zero due to the QCD potential. As a consequence, the axion obtains a tiny mass $m_a \propto 1/f_a$ where f_a is the axion decay constant restricted to the range $f_a \sim 10^{8-12}$ GeV by various observations. For the KSVZ axion [3], which couples almost only to protons, the strongest lower bound can be obtained from SN1987A: $f_a > 4 \times 10^8$ GeV [5]. The upper limit is set by the axion dark matter abundance, which is produced by the standard misalignment mechanism [6]. On the other hand, one can rely on the kinetic misalignment [7] as the origin of the dark matter density even for a lower value of f_a .

The axion kinetic misalignment can also be the source of the observed baryon asymmetry through the mechanism of spontaneous baryogenesis [8], as it can provide a CPT violating background while the electroweak sphaleron interaction, that violates $B + L$ symmetry, is in equilibrium [9]. The final baryon asymmetry is determined at the electroweak phase transition below which the sphaleron interaction decouples. However, it conflicts with the dark matter abundance from standard QCD axions.

To resolve this difficulty, we consider a simple extension of the PQ symmetry to include the seesaw mechanism that explains the observed neutrino masses and mixing [10], which involves lepton number violation. This enables spontaneous leptogenesis that generates $B - L$ asym-

metry around the seesaw scale [11,12]. Adjusting this scale, we can realize theogenesis of baryon asymmetry and dark matter of the Universe.

Another key ingredient of spontaneous leptogenesis is the origin of the initial kinetic misalignment, which could be related to inflation dynamics [13]. In this regard, it appears natural to explore a minimal setup where the PQ field, breaking the PQ symmetry, drives inflation. Its radial degree is supposed to play a role of the inflaton and the axial degree (axion) develops the initial kinetic motion at the end of inflation through a higher-dimensional operator violating the PQ symmetry explicitly [14]. We focus on the KSVZ model for solving the strong CP problem and accommodate neutrino masses by Type I seesaw mechanism. Then, we will examine theogenesis condition and analyze the relevant inflation and post-inflation dynamics in the framework of pole inflation. As a result, we find out the allowed parameter regions which turns out to be quite limited. On the other hand, the DFSZ framework [4] is not considered, since reheating through PQ field decay into Higgs doublets is too inefficient to achieveogenesis. Recent developments of spontaneous leptogenesis with the axion or axion-like field in connection to neutrino masses can be referred to Refs. [15–17].

2. Cogenesis requirements

2.1. Spontaneous leptogenesis and dark matter abundance

Consider a field Φ that carries the PQ charge +1 and breaks $U(1)_{PQ}$ spontaneously. After breaking the symmetry, it is represented by Φ

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$= \frac{v_a}{\sqrt{2}} e^{i\theta}$ where $\theta \equiv a/v_a$. Given the kinetic misalignment $\dot{\theta}$ at a temperature T , the PQ number normalized by the entropy density is a conserved quantity:

$$Y_\theta \equiv \frac{n_\theta}{s} = \frac{v_a^2 \dot{\theta}}{\frac{2\pi^2}{45} g_* T^3} \quad (1)$$

In the background of $\dot{\theta}$, neutrino Yukawa interactions, $y_\nu L N H$ with an RHN N , may be in equilibrium to generate a non-vanishing $B-L$ asymmetry: $Y_{B-L} \equiv \frac{n_{B-L}}{s}$ where $n_{B-L} = \frac{1}{6} \mu_{B-L} T^2$ and the chemical potential $\mu_{B-L} \propto \dot{\theta}$ is determined by the equilibrium conditions of a given system. Thus, one can parameterize the resulting baryon asymmetry Y_B by

$$Y_{B-L} = c_{B-L} Y_\theta \left(\frac{T_{B-L}}{v_a} \right)^2 \quad (2)$$

where $c_{B-L} \equiv \mu_{B-L}/6\dot{\theta}$ is a model-dependent constant, and T_{B-L} is the temperature at which the final $B-L$ asymmetry is determined. Then, the final baryon asymmetry is $Y_B = c_{\text{sph}} Y_{B-L}$ where c_{sph} is the usual $(B-L)$ -to- B conversion factor which is given by $c_{\text{sph}} = \frac{28}{79}$ when all the Yukawa interactions are in equilibrium (namely, there is no flavor dependence).

The neutrino Yukawa interactions violating lepton number in our case are the decay and inverse decay of N , and $T_{B-L} \equiv M_N/z_N$ depends on their strength. In the so-called strong washout regime, where the decay rate $\Gamma_D = y_\nu^2 M_N/8\pi$ of N is higher than the Hubble rate at temperature $T = M_N$: $K \equiv \Gamma_D/H(M_N) \gg 1$ and thus the inverse decay is active to erase the lepton asymmetry even at lower temperature $T < M_N$ [18], Y_{B-L} is determined by equilibrium conditions which is assumed in the traditional treatment and T_{B-L} is set by the temperature when the inverse decay goes out of equilibrium. In the opposite regime $K \lesssim 1$, the effect of CP-violating RHN decay has to be considered properly to get Y_{B-L} . The efficiency factor $\kappa \equiv Y_{B-L}/\tilde{Y}_{B-L}$, with $\tilde{Y}_{B-L} \equiv c_{B-L} \dot{\theta}_1 T_1^2/6s(T_1)$ and $T_1 = M_N$, was computed as a function of K in [12]. The parameter $z_N = M_N/T_{B-L}$ is related to κ through $z_N = 1/\sqrt{\kappa}$ where the value of κ can be read from the y -axis of Fig. A.1 of [12]. For our presentation, we consider two benchmark cases: $K = 4$, which yields relatively larger efficiency, and $K = 50$, corresponding to the effective neutrino mass that parametrizes the decay rate, taken to be at the atmospheric neutrino mass scale: $\tilde{m}_\nu \equiv y_\nu^2 v_{EW}^2/M_N = 8\pi\Gamma_D v_{EW}^2/M_N^2 = 0.05$ eV [18]. We remark that the effective neutrino mass defined in terms of one right handed neutrino is not equivalent to a physical neutrino mass obtained in the three generation picture, and can be larger than the actual neutrino mass scale [18]. The corresponding efficiency factors are $\kappa \approx 0.03$ and 0.01 , respectively. In our model, a nonzero initial abundance of the RH neutrinos can be non-thermally produced at reheating from the inflaton decays, although being suppressed for $y_N \ll y_Q$. But, these two benchmark points with $K = 4, 50$ are less sensitive to the initial abundance of the RH neutrinos at reheating.

The kinetic misalignment can also be a source of the axion dark matter density [7,9]. The initial energy density of the axion dominated by the kinetic energy redshifts quickly to the QCD potential energy. Then, the axion trapped by the potential barrier starts oscillation around its minimum which can happen later than the conventional oscillation temperature. In this case, our Y_θ becomes the main source of dark matter density if the following condition is met [7,9]:

$$2m_a |Y_\theta| \approx 0.44 \text{ eV}. \quad (3)$$

Following the convention of the anomaly coupling of the axion a to the gluon fields:

$$-\mathcal{L}_{aGG} = \frac{1}{32\pi^2} \frac{a}{f_a} G\tilde{G}, \quad (4)$$

we get the relation $v_a = N_{\text{DW}} f_a$ where N_{DW} is the QCD anomaly coefficient which we will also denote as c_S . Then, the axion mass is determined to be [19]

$$m_a \approx 5.7 \text{ meV} \left(\frac{10^9 \text{ GeV}}{f_a} \right). \quad (5)$$

Combining Eqs. (2) and (3) with $Y_B = c_{\text{sph}} Y_{B-L}$, one finds a successful cogenesis with $Y_B \approx 0.87 \times 10^{-10}$:

$$2.25 \frac{N_{\text{DW}}^2}{|c_B|} \left(\frac{f_a}{10^9 \text{ GeV}} \right) \left(\frac{10^3 \text{ GeV}}{M_N/z_N} \right)^2 = 1, \quad (6)$$

where $c_B \equiv c_{\text{sph}} c_{B-L}$. In the following subsections, we will examine the KSVZ axion model that incorporates the seesaw mechanism for neutrino masses and determine the coefficient c_B appearing in the baryon asymmetry in (2).

2.2. KSVZ axion and seesaw mechanism

The KSVZ axion is implemented by the existence of extra quarks (Q, Q^c) . In our scheme, RHNs couple to the PQ field Φ and thus we have

$$-\mathcal{L}_{\text{KSVZ}} = y_Q \Phi Q Q^c + \frac{1}{2} y_N \Phi N N + h.c. \quad (7)$$

where $\Phi = \frac{v_a}{\sqrt{2}} e^{ia/v_a}$ is the PQ field (with the PQ charge +1) after the PQ symmetry breaking, and $Q Q^c$ is a Dirac fermion bilinear with the mass $M_Q = \frac{y_Q}{\sqrt{2}} v_a$ and N gets the Majorana mass $M_N = \frac{y_N}{\sqrt{2}} v_a$ after rotating away the phase field e^{ia/v_a} . Note that we use the chiral notation for the fermions. At the energy scale below M_Q , the Standard Model (SM) Lagrangian includes

$$\begin{aligned} -\mathcal{L}_{\text{Yukawa}} &= y_u q u^c H + y_d q d^c \tilde{H} + y_e l e^c \tilde{H} + y_\nu l N H + h.c. \\ -\mathcal{L}_{\text{anomaly}} &= \frac{c_S}{32\pi^2} \frac{a}{v_a} G\tilde{G} + \frac{c_W}{32\pi^2} \frac{a}{v_a} W\tilde{W} \end{aligned} \quad (8)$$

where the chiral notation is also used for the quark (lepton) doublets $q(l)$ and singlets $u^c, d^c(e^c)$, and the Higgs doublet and its conjugate are denoted by H and $\tilde{H}$. Here we suppressed the family indices which will be recovered whenever necessary. In the anomaly term, we keep only the axion couplings to the $SU(3)_c$ and $SU(2)_L$ gauge bosons that have the anomaly coefficients c_S and c_W , respectively. In the KSVZ model, $c_S = N_{\text{DW}}$ and $c_W = 0$, are given by the total number of heavy quarks and the number of heavy quark doublets, respectively.

Assuming all the Yukawa and sphaleron interactions are in equilibrium, one can generalize the calculation in Ref. [20] to include the seesaw sector. As the first generation quark Yukawa interactions are much weaker than the others as well as the strong and weak sphaleron interactions: $\Gamma_{y_{u1,d1}} \ll \Gamma_{y_{u2,3,d2,3}}, \Gamma_{SS,WS}$, the equilibrium conditions are satisfied by the following chemical potential relations involving only quarks:

$$\begin{aligned} \mu_{q_i} + \mu_{u_i^c} + \mu_H &= c_S \frac{m_d^2}{m_u^2 + m_d^2} \dot{\theta} \delta_{i1}, \\ \mu_{q_i} + \mu_{d_i^c} - \mu_H &= c_S \frac{m_u^2}{m_u^2 + m_d^2} \dot{\theta} \delta_{i1}, \\ \sum_i (2\mu_{q_i} + \mu_{u_i^c} + \mu_{d_i^c}) &= c_S \dot{\theta}. \end{aligned} \quad (9)$$

On the other hand, the equilibrium condition involving leptons requires

$$\begin{aligned} \mu_{l_i} + \mu_{e_i^c} - \mu_H &= 0, \\ \mu_{l_i} + \mu_H &= x_l \dot{\theta}, \\ \sum_i (3\mu_{q_i} + \mu_{l_i}) &= c_W \dot{\theta}, \end{aligned} \quad (10)$$

with no flavor dependence: $\mu_{l_i} = \mu_l$, etc. Note that the second equation comes from the shift in the chemical potential $\mu_l \rightarrow \mu_l - x_l \dot{\theta}$ as leptons carry PQ charges, and thus their charge operators couple to $\dot{\theta}$.

Therefore, we find the chemical potential for the $B-L$ charge ($q_{B-L} = \frac{1}{6} \mu_{B-L} T^2$):

$$\mu_{B-L} = \frac{1}{33} \left(28c_W - c_S \frac{57m_u^2 - 15m_d^2}{m_u^2 + m_d^2} - 153x_l \right) \dot{\theta} \quad (11)$$

where x_l is the PQ charge of the lepton doublets l_i ($x_l = \frac{1}{2}$ in our case). This results in

$$c_B = \frac{28}{79} \frac{1}{198} \left(28c_W - c_S \frac{57m_u^2 - 15m_d^2}{m_u^2 + m_d^2} - 153x_l \right). \quad (12)$$

Taking $c_S = N_{DW} = 1$, $c_W = 0$, $x_l = 1/2$, and $m_u/m_d = 0.462$ [21], we find $c_B = -0.13$.

Let us finally comment on the coupling of the PQ field to the Higgs doublet: $\lambda_{\Phi H} |\Phi|^2 |H|^2$. To avoid fine-tuning in the Higgs sector, one may require $\lambda_{\Phi H} < \langle H \rangle^2 / \langle \Phi \rangle^2$, which is smaller than about 10^{-14} . This coupling is too small to provide a sufficient reheating in PQ inflation scenarios, which will be discussed in the following section.

Given $c_B = -0.13$, we can determine the value of M_N or the Yukawa coupling $y_N = \sqrt{2} M_N / v_a$ forogenesis (6):

$$y_N = 5.9 \times 10^{-6} z_N \sqrt{\frac{10^9 \text{GeV}}{f_a}}. \quad (13)$$

As noted in Section 2.1, we have $z_N \approx 5.8$ and 10 for $K = 4$ and 50 , respectively.

3. PQ pole inflation and axion kinetic misalignment

Let us begin by recapitulating the main features of the PQ inflation at the pole that we use for our purpose. In pole inflation, the inflaton's kinetic term becomes singular at a specific point, which effectively stretches the field space. When the inflaton field approaches the pole, the redefinition of the field to make its kinetic term canonical results in an effectively flat potential [22,23]. Thus, a wide range of different underlying potentials, when subjected to the pole in the kinetic term, can drive inflation and lead to very similar predictions for the observable cosmological parameters such as the spectral index n_s and the tensor-to-scalar ratio r .

In the PQ mechanism, $U(1)_{\text{PQ}}$ is broken by a complex (PQ) field Φ typically by a quartic potential

$$V_{\text{PQC}} = \lambda_{\Phi} \left(|\Phi|^2 - \frac{v_a^2}{2} \right)^2, \quad (14)$$

which we take as an underlying inflaton potential where the radial mode of Φ acts as the inflaton. The axial mode of Φ is the axion whose kinetic misalignment is supposed to generate the baryon asymmetry of the universe and then turns into dark matter. For this purpose, we introduce a higher-dimensional operator that softly violates the PQ symmetry [14]:

$$V_{\text{PQV}} = \lambda_n \frac{\Phi^n}{M_P^{n-4}} + \text{h.c.} \quad (15)$$

which is supposed not to perturb the features of inflation driven by (14), and the axion solution to the strong CP problem [24].

3.1. PQ pole inflation

We consider the Jordan frame Lagrangian for the PQ field Φ with a non-conformal coupling

$$\frac{\mathcal{L}_J}{\sqrt{-g_J}} = -\frac{M_P^2}{2} \Omega(\Phi) R + |\partial_{\mu} \Phi|^2 - \Omega^2 V_E. \quad (16)$$

where $V_E = V_{\text{PQC}} + V_{\text{PQV}}$ and the non-minimal coupling function is taken to be

$$\Omega(\Phi) = 1 - 2 \frac{|\Phi|^2}{\Lambda^2}. \quad (17)$$

Here, the new cutoff-scale Λ in the non-minimal coupling is related to a parameter $\zeta \equiv M_P^2 / \Lambda^2$, which was taken to a conformal coupling, $\zeta = \frac{1}{6}$, in Ref. [14].

We can go to the Einstein frame by making a Weyl transformation: $g_{J,\mu\nu} = g_{E,\mu\nu} / \Omega$:

$$\begin{aligned} \frac{\mathcal{L}_E}{\sqrt{-g_E}} &= -\frac{M_P^2}{2} R + \frac{1}{\Omega} |\partial_{\mu} \Phi|^2 + \frac{3}{4} M_P^2 \frac{(\partial_{\mu} \Omega)^2}{\Omega^2} - V_E \\ &= -\frac{M_P^2}{2} R + \frac{1 + \zeta(6\zeta - 1) \frac{\rho^2}{M_P^2}}{2(1 - \zeta \frac{\rho^2}{M_P^2})^2} (\partial_{\mu} \rho)^2 + \frac{\rho^2 (\partial_{\mu} \theta)^2}{2(1 - \zeta \frac{\rho^2}{M_P^2})} - V_E \end{aligned} \quad (18)$$

where we used the notation $\Phi = \frac{1}{\sqrt{2}} \rho e^{i\theta}$ in the second line. The potential $V_E = V_{\text{PQC}} + V_{\text{PQV}}$ is explicitly given by

$$\begin{aligned} V_{\text{PQC}} &= \frac{\lambda_{\Phi}}{4} (\rho^2 - v_a^2)^2, \\ V_{\text{PQV}} &= \frac{|\lambda_n|}{\sqrt{2^{n-2}}} \frac{\rho^n}{M_P^{n-4}} \cos(n\theta + \delta). \end{aligned} \quad (19)$$

Here, δ is the phase of a complex parameter λ_n .

One can see that a pole appears at $\rho = M_P / \sqrt{\zeta} = \Lambda$ around which inflation is occurring. To get the canonically normalized inflaton field ϕ , we make the field redefinition $\rho = \rho(\phi)$ satisfying

$$\frac{d\rho}{d\phi} = \frac{1}{\sqrt{\Xi(\rho)}} \equiv \frac{(1 - \zeta \frac{\rho^2}{M_P^2})}{\sqrt{1 + \zeta(6\zeta - 1) \frac{\rho^2}{M_P^2}}}, \quad (20)$$

where ϕ is the canonical normalized real scalar. In the limit $\rho \rightarrow \frac{M_P}{\sqrt{\zeta}}$ during inflation, we can approximately get the canonical normalized field for $\zeta \gtrsim 1$ as

$$\begin{aligned} \frac{d\rho}{d\phi} &\simeq \frac{(1 - \zeta \frac{\rho^2}{M_P^2})}{\sqrt{6\zeta}}, \\ \rho &\simeq \frac{M_P}{\sqrt{\zeta}} \tanh\left(\frac{\phi}{\sqrt{6} M_P}\right). \end{aligned} \quad (21)$$

We rewrite the PQ field Lagrangian in terms of ϕ and θ , where ρ is understood to be a function of ϕ , $\rho(\phi)$:

$$\mathcal{L}_{\text{PQ}} = \frac{1}{2} (\partial_{\mu} \phi)^2 + \frac{1}{2} \frac{\rho^2}{\Omega} (\partial_{\mu} \theta)^2 - V_E(\rho, \theta) \quad (22)$$

where $\Omega = 1 - \zeta \frac{\rho^2}{M_P^2}$. Then, dynamics of the inflaton ϕ and the axion θ is governed by

$$\begin{aligned} \ddot{\phi} + 3H\dot{\phi} - \frac{\rho}{\Omega^2} \frac{d\rho}{d\phi} \dot{\theta}^2 + \frac{\partial V_E}{\partial \rho} \frac{d\rho}{d\phi} &= 0, \\ \ddot{\theta} + 3H\dot{\theta} + \frac{2}{\rho\Omega} \frac{d\rho}{d\phi} \dot{\phi} \dot{\theta} + \frac{\Omega}{\rho^2} \frac{\partial V_E}{\partial \theta} &= 0 \end{aligned} \quad (23)$$

where the Hubble parameter is determined by

$$H^2 = \frac{1}{3M_P^2} \left(\frac{1}{2} \dot{\phi}^2 + \frac{1}{2} \frac{\rho^2}{\Omega} \dot{\theta}^2 + V_E \right).$$

During inflation around the pole, the inflaton energy density and the slow-roll parameters are assumed to be dominated by ϕ , that is, $V_E \simeq V_{\text{PQC}} \gg V_{\text{PQV}}$, $\partial V_{\text{PQC}} / \partial \phi \gg \partial V_{\text{PQV}} / \partial \phi$, as well as

$$\epsilon_{\phi} = \frac{M_P^2}{2V_E^2} \left(\frac{\partial V_E}{\partial \rho} \frac{d\rho}{d\phi} \right)^2 \gg \epsilon_{\theta} = \frac{M_P^2}{2V_E^2} \frac{\Omega}{\rho^2} \left(\frac{\partial V_E}{\partial \theta} \right)^2. \quad (24)$$

Then, the inflation energy V_I and the slow-roll parameter $\epsilon_* \simeq \epsilon_{\phi*}$ at the horizon exit are

$$V_I \simeq \frac{1}{4} \frac{\lambda_{\Phi}}{\zeta^2} M_P^4, \quad \epsilon_* \simeq \frac{3}{4N^2} \quad (25)$$

where N is the number of e-folding. Thus, the CMB normalization tells us that

$$A_s = \frac{1}{24\pi^2 \epsilon_*} \frac{V_I}{M_P^4} = 2.1 \times 10^{-9} \Rightarrow \frac{\lambda_{\Phi}}{\zeta^2} \simeq 4 \times 10^{-10} \quad (26)$$

for $N = 60$. Inflation ends when $\epsilon \simeq \epsilon_{\phi} = 1$, corresponding to $\rho = \rho_e$:

$$\rho_e \simeq \frac{0.732}{\sqrt{\zeta}} M_P = 0.732 \Lambda, \quad (27)$$

in the regime $\zeta \gg 1$. Applying the slow-roll condition until the end of inflation, the evolution Eq. (23) simplifies to

$$3H\dot{\phi} + \frac{\partial V_E}{\partial \phi} \approx 0, \quad 3H\dot{\theta} + \frac{\Omega}{\rho^2} \frac{\partial V_E}{\partial \theta} \approx 0, \quad (28)$$

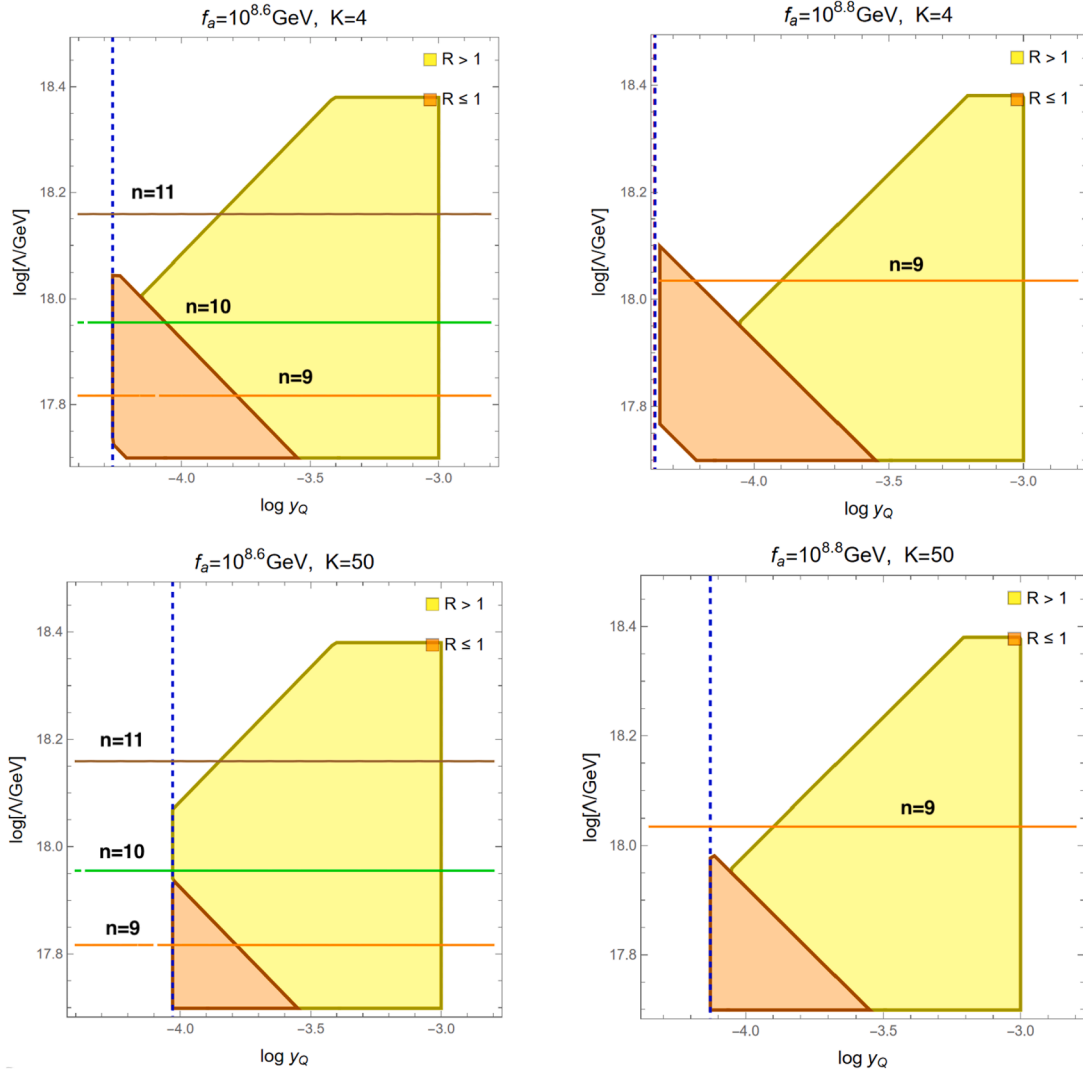


Fig. 1. Region of successfulogenesis in the parameter space (y_Q, Λ) . The blue-dashed line shows $y_Q = y_N$. For $y_Q > y_N$, the decays of the inflaton into extra heavy quarks are dominant processes for reheating. The boundary between the orange and yellow shaded regions corresponds to $R = 1$ or $y_Q = \sqrt{\lambda_\Phi}$. The upper-left boundary of the yellow region is given by the reheating condition for $R > 1$.

from which we find $H \approx \sqrt{V_{\text{PQC}}/3M_P^2}$, and the axion number density at the end of inflation is given by

$$n_{\theta,e} = \frac{\rho_e^2}{\Omega_e} \dot{\theta}_e \approx -\frac{1}{3H} \left. \frac{\partial V_{\text{PQV}}}{\partial \theta} \right|_{\rho=\rho_e}, \quad (29)$$

with V_{PQV} given in (19). Thus, n_θ is determined by the coupling parameter λ_n and the dimensionality n of V_{PQV} , which are supposed to satisfy the inflation conditions in and above (24). They are also bounded by the condition that the quality of the axion solution to the strong CP problem should not be spoiled [24], which reads, in terms of the mass-squared induced by V_{PQV} at $\rho = v_a$,

$$\delta m_a^2 \equiv n^2 \frac{|\lambda_n|}{\sqrt{2^{n-2}}} \frac{v_a^{n-2}}{M_P^{n-4}} \lesssim 10^{-10} \frac{n}{N_{\text{DW}}} m_a^2, \quad (30)$$

where m_a is given by (5). We find that this condition is much milder than the inflation conditions in the parameter space allowing the cogenesis. For the purposes of our analysis, we normalize all sinusoidal functions of θ to unit amplitude, and consider the case of $v_a = f_a$ ($N_{\text{DW}} = 1$).

3.2. Post-inflation evolution: Kinetic misalignment and reheating

After the end of inflation, the axion number $a^3 n_\theta$ in a volume a^3 is approximately conserved since the potential V_{PQV} becomes negligible as the inflaton field ϕ (or ρ) continues to roll down the potential V_{PQC} . That is, n_θ at the scale a is

$$n_\theta = n_{\theta,e} \left(\frac{a_e}{a} \right)^3. \quad (31)$$

The inflaton also starts a damped oscillation around the minimum of the potential. During this epoch, the oscillation amplitude of ϕ and the Hubble parameter decrease due to the Hubble friction. The factor, $\Xi(\rho) \equiv (d\rho/d\phi)^{-2}$ in the kinetic term driving the pole inflation, converges to 1 when $\rho \ll \rho_c \equiv \frac{M_P}{\sqrt{6}\zeta}$ for which ρ becomes a canonically normalized field, $\rho \approx \phi$. In this era, the potential V_{PQC} scales like radiation: $V_{\text{PQC}} \propto 1/a^4$ [25]. For the case of our interest, ρ_c is comparable to ρ_e , and thus the universe approaches quickly to the radiation-like era when reheating occurs.¹ To quantify the impact of this intermediate era, we solve numerically the equation of motion in terms of the non-canonical

¹ When the inflaton self-coupling is dominant for reheating as in our case, the inflaton fragmentation can be relevant when the couplings between the inflaton

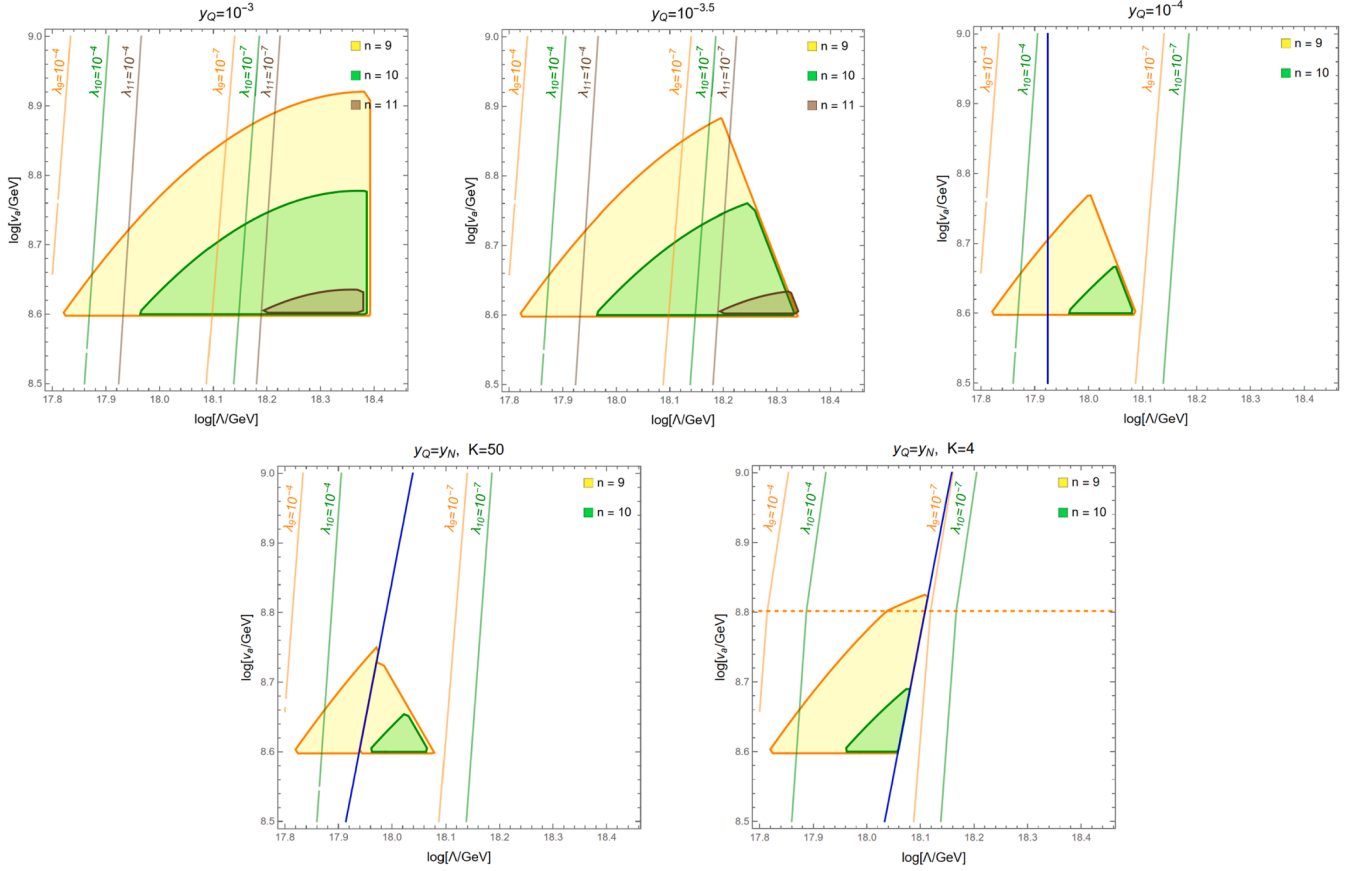


Fig. 2. Parameter space for the successfulogenesis for different values of y_Q in terms of the cutoff scale Λ vs. the PQ breaking scale f_a . Allowed regions are indicated by orange, green, and brown, denoting PQ-violating potentials with dimensionalities $n = 9, 10$, and 11 , respectively. The required values for the coupling λ_n are also depicted by almost vertical lines. The left (right) side of the blue line is for $R < 1$ (> 1).

field ρ :

$$\ddot{\rho} + 3H\dot{\rho} + \frac{1}{2} \frac{d\Xi}{d\rho} \dot{\rho}^2 + \frac{1}{\Xi} \frac{dV_{\text{PQC}}}{d\rho} \approx 0, \quad H^2 = \frac{1}{3M_P^2} \left[\frac{1}{2} \Xi \dot{\rho}^2 + V_{\text{PQC}} \right], \quad (32)$$

driven from (18), (20). From this, we obtain the numerical value of the quantity c_N as shown in detail in the appendix:

$$c_N \equiv \frac{V_{\text{PQC}}(\rho_c) a_c^4}{V_{\text{PQC}}(\rho_e) a_e^4} \approx 0.37\zeta + 1.79, \quad (33)$$

which allows us to parametrize the axion number density at the scale factor $a > a_c$ (or $\rho = \phi < \rho_c$), as

$$n_\theta(\phi) = n_{\theta,e} c_N^{-\frac{3}{4}} \left(\frac{V_{\text{PQC}}(\phi)}{V_{\text{PQC}}(\rho_e)} \right)^{\frac{3}{4}} \quad (34)$$

where $V_{\text{PQC}}(\phi)$ is the inflaton energy density at the scale factor a corresponding to $\rho = \phi$. Assuming that reheating is complete for $V_{\text{PQC}}(\phi_{\text{RH}}) = g_* \pi^2 T_{\text{RH}}^4 / 30$, we obtain the PQ number abundance Y_θ at the reheating temperature T_{RH} by

$$Y_\theta \equiv \frac{n_\theta(T_{\text{RH}})}{s(T_{\text{RH}})} \approx 2.5 \times 10^7 c_N^{-\frac{3}{4}} \frac{n_{\theta,e}}{M_P^3}. \quad (35)$$

In our scenario for the KSVZ axion, the reheating process can occur due to the inflaton decays into the extra heavy quarks or RHNs,

and the SM are not large enough [26]. However, we assume that the inflaton condensate remains homogeneous during reheating and focus on the perturbative decays of the inflaton in our work. It would be worthwhile to perform a more dedicated analysis of the inflaton fragmentation in a future work.

$\phi \rightarrow QQ^c$ or NN , through the Yukawa couplings in Eq. (7), as far as they are allowed kinematically. We assume that reheating is dominated by the decay process $\phi \rightarrow QQ^c$, whereas the mass of RHN is already constrained by the cogenesis condition as in Eq. (13), setting a lower bound of the reheating temperature. We note that the inflaton effective mass is $m_\phi(t) \equiv \sqrt{V''(\phi_0(t))} = \sqrt{3}\lambda_\phi\phi_0(t)$ at a given scale (time) $a(t)$ for the quartic inflaton potential, whereas the masses of the fermions are $m_f(t) = y_f\phi(t)/\sqrt{2}$ for $f = Q$ and N . Here, we note that $\phi(t) = \phi_0(t)\mathcal{P}(t)$ where $\mathcal{P}(t)$ is a periodic function that crosses zero [14]. While the envelope of the inflaton field, $\phi_0(t)$ (“a slowly varying envelope”), red-shifts as $\phi_0(t) \propto 1/a(t)$, the highly oscillating modes of the inflaton with the frequency $E_n(t) \approx 0.5nm_\phi(t)$ can decay into fermions f if $E_n(t) > 2m_f(t)$. Furthermore, even the zero mode of the inflaton field can decay into fermions as the inflaton field approaches zero. Then, the time averaged decay rate of the inflaton over one oscillation is expressed in terms of an “effective coupling” $y_{f,\text{eff}}$ [25]:

$$\Gamma_\phi(t) = \frac{N_f}{8\pi} \left(\frac{y_{f,\text{eff}}}{\sqrt{2}} \right)^2 m_\phi(t), \quad y_{f,\text{eff}}^2 \approx 0.5\alpha(\mathcal{R})y_f^2, \quad \mathcal{R} \equiv \left(\frac{4m_f}{m_\phi} \right)^2 \Big|_{\phi(t)=\phi_0} \quad (36)$$

where N_f is the number of degrees of freedom. In the two limits of $\mathcal{R}$, we have

$$\alpha(\mathcal{R}) \approx \begin{cases} 1 & \mathcal{R} \ll 1 \\ 0.5\mathcal{R}^{-\frac{1}{2}} & \mathcal{R} \gg 1 \end{cases} \quad (37)$$

For our analysis, we extend the first expression to $\mathcal{R} = 1$, and the second expression is reliable for $\mathcal{R} \geq 3$. In the intermediate region $1 < \mathcal{R} < 3$, we took $\alpha(\mathcal{R}) = \text{Min}(1, 0.5\mathcal{R}^{-\frac{1}{2}})$. Combining (37) with our approximate

relation $\rho_{\text{RH}}(T_{\text{RH}}) = \rho_\phi(\phi_{\text{RH}}) \simeq \frac{\lambda_\phi}{4} \phi_{\text{RH}}^4$, we get

$$T_{\text{RH}} \simeq 1.5 \times 10^6 \text{ GeV} \left(\frac{y_{f,\text{eff}}}{10^{-4}} \right)^2 \left(\frac{\lambda_\phi}{10^{-11}} \right)^{1/4}, \quad (38)$$

which agrees well with a rigorous analysis [25].

4. Parameter space for the PQ cogenesis

In Fig. 1, we show regions of the successful cogenesis in the parameter space of the Yukawa coupling of the extra heavy quark y_Q and the cutoff scale Λ for two different values of $f_a = 10^{8.6} \text{ GeV}$ and $10^{8.8} \text{ GeV}$. The orange and yellow regions represent $\mathcal{R} < 1$ and $\mathcal{R} > 1$, respectively. The maximum value of y_Q is taken to be 10^{-3} above which the one-loop correction to λ_ϕ becomes bigger than the tree-level value. On the other hand, we assume that the inflaton decays into extra heavy quarks are dominant processes for reheating, so the lower limit on $y_Q = y_N$ is indicated in the blue-dashed line. The region above the upper-left boundary of the yellow region, which satisfies $\mathcal{R} > 1$, is excluded, because reheating cannot occur as the inflaton settles down to the true vacuum, $\phi \simeq f_a$, and the decay channel is kinematically closed for $y_Q > \sqrt{\lambda_\phi}$. We also note that the upper boundary of the yellow region is given by $\Lambda < M_P$.

Moreover, in Fig. 1, we also indicate the regions allowed by the successful cogenesis in orange, green, and brown horizontal lines, for the dimensions of the PQ-violating potential, $n = 9, 10$, and 11 , respectively. The correct axion relic density, the cogenesis condition as well as inflation condition, $\frac{dV_{\text{PQC}}}{d\phi} > \frac{dV_{\text{PQV}}}{d\phi}$, are satisfied above the horizontal line for each case. Thus, as the dimension of the PQ-violating potential increases (or from bottom to top), the consistent parameter space for the Yukawa coupling of the extra heavy quark decreases. We note that there is no consistent parameter space in the right plot of Fig. 1 for $n = 10, 11$, in contrast to the left plot of Fig. 1.

The maximum reheating temperature and the corresponding field value ϕ_{RH} are given by

$$T_{\text{RH}} \approx 1.2 \times 10^6 \left(\frac{M_P}{\Lambda} \right)^3 \text{ GeV}, \quad \phi_{\text{RH}} \approx 6.7 \times 10^8 \left(\frac{M_P}{\Lambda} \right)^2 \text{ GeV}, \quad (39)$$

respectively. We note that $\phi_{\text{RH}} \ll \rho_c = \frac{M_P}{\sqrt{6\zeta}}$ in the region of our interest $\Lambda > 10^{17} \text{ GeV}$.

Our final result is presented in Fig. 2 showing the regions of cogenesis in the parameter space of (Λ, f_a) . The allowed regions are found only for PQ-violating potentials V_{PQV} with dimensionalities $n = 9, 10$, and 11 corresponding to the areas in orange, green, and brown, respectively. The upper-left boundaries are determined by the inflation condition, $\frac{dV_{\text{PQC}}}{d\phi} > \frac{dV_{\text{PQV}}}{d\phi}$. Inflaton decays to extra heavy quarks govern the reheating process, depending on the Yukawa coupling of the extra heavy quark, which is restricted to $y_N \leq y_Q \leq 10^{-3}$. The blue line indicates $\mathcal{R} = 1$; the region to its left (right) corresponds to $\mathcal{R} < 1$ (> 1). The upper-right boundaries are set by the condition whether the reheating can occur. We also find that the cogenesis requires the values of the PQV couplings to be in the ranges of $10^{-6.6} < \lambda_9 < 10^{-4.5}$ and $10^{-6.1} < \lambda_{10} < 10^{-5}$, for all the choices of y_Q in Fig. 2, and $10^{-9} \lesssim \lambda_{11} \lesssim 10^{-7}$ for $y_Q = 10^{-3}, 10^{-3.5}$ in the upper panel of Fig. 2. We note that for $n < 8$, the axion quality in (30) constrains the PQ-violating potential more strongly so the resulting PQ number density is not sufficient for a successful cogenesis. On the other hand, for $n > 11$, the PQ-violating potential is constrained by the inflaton condition similarly as in the case with $n < 11$, but it is more suppressed at the end of inflation due to a sub-Planckian inflaton field value, resulting in a smaller PQ number density.

5. Conclusion

In this work, we have investigated a minimal extension of the KSVZ axion model to simultaneously explain the dark matter abundance and the baryon asymmetry of the Universe. To resolve the tension between

the standard spontaneous baryogenesis mechanism and the overproduction of axion dark matter, we incorporated the seesaw mechanism into the Peccei-Quinn sector. This setup allows for spontaneous leptogenesis, where a $B - L$ asymmetry is generated near the seesaw scale and subsequently converted into the observed baryon asymmetry via electroweak sphalerons. We also addressed the origin of the initial kinetic misalignment required for this scenario by identifying the PQ field as the driver of inflation. Within the framework of pole inflation, the radial component of the PQ field plays the role of the inflaton, while the angular degree of freedom (the axion) acquires a large initial velocity through higher-dimensional PQ-breaking operators at the end of inflation. Our analysis demonstrates that the cogenesis of baryon asymmetry and dark matter is indeed viable in this framework. Although the allowed region is found to be quite limited: $f_a = (4 - 9) \times 10^8 \text{ GeV}$, our results suggest that the QCD axion, coupled with the seesaw mechanism and pole inflation, remains a compelling candidate for solving the strong CP problem while providing a unified origin for the matter content of the Universe.

Data availability

No data was used for the research described in the article.

Declaration of competing interest

The authors declare the following financial interests/personal relationships which may be considered as potential competing interests: Hyun Min Lee reports financial support was provided by National Research Foundation of Korea. Jun-Ho Sing reports financial support was provided by National Research Foundation of Korea. If there are other authors, they declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

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Appendix A.

In this appendix, we consider the evolution of ρ_ϕ right after inflation. To that, we solve the equation of motion of the non-canonical ρ field numerically. From the Lagrangian in (18), we get

$$\ddot{\rho} = - \left(\frac{1}{2} \frac{d\Xi}{d\rho} \frac{\rho^2}{\Xi(\rho)} + \frac{1}{\Xi(\rho)} \frac{dV_{\text{PQC}}}{d\rho} + 3H\dot{\rho} \right) \quad (A.1)$$

where

$$\Xi(\rho) \equiv \frac{1 + \zeta(6\zeta - 1) \frac{\rho^2}{M_P^2}}{\left(1 - \zeta \frac{\rho^2}{M_P^2} \right)^2}. \quad (A.2)$$

At the end of inflation, the Hubble parameter is defined by

$$H^2 = \frac{1}{3M_P^2} \left[\frac{1}{2} \Xi(\rho) \dot{\rho}^2 + V_{\text{PQC}}(\rho) \right] \quad (A.3)$$

where $H^2 \simeq V_{\text{PQC}}/3M_P^2$. Together $\rho_e \approx 0.732/\sqrt{\zeta}$ and (28), we determine $\dot{\rho}_e$ by using

$$\dot{\rho} \approx - \frac{\partial V_{\text{PQC}}}{\partial \rho} \left(\frac{\partial \rho}{\partial \phi} \right)^2 \frac{1}{3H}, \quad (A.4)$$

which sets the initial condition for the numerical evolution in (A.1). In Fig. A.1, we show the asymptotic values of c_N defined in (33) for ζ

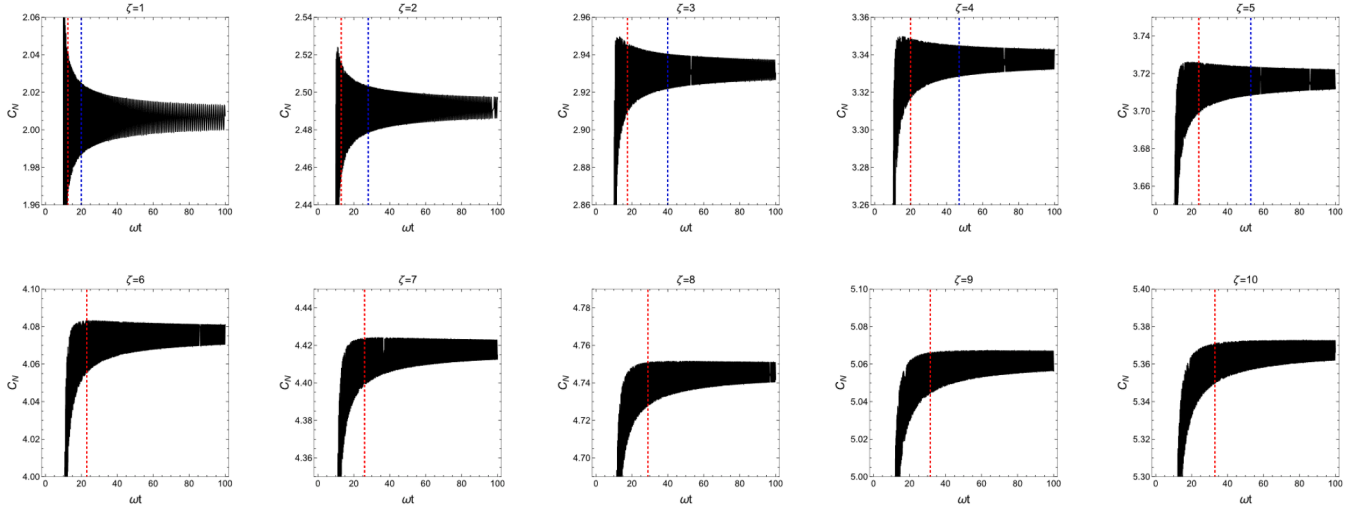


Fig. A.1. Time evolution of c_N for $\zeta = 1, 2, \dots, 10$, appearing in Eq. (35). The plots show that c_N rapidly converges to a constant value for a given ζ . Here, $\omega t = 10$ with $\omega = 10^{11}$ GeV, and t_e is the time at the end of inflation. The red and blues lines indicate the times when $\rho = \rho_c/10$ and $\rho = \rho_c/20$, respectively.

$= 1, 2, \dots, 10$, when $\rho \ll \rho_c = \frac{M_P}{\sqrt{6}\zeta}$. These reference times when $\rho = \rho_c/10, \rho_c/20$, indicate the onset of the regime where the non-canonical factor becomes $\Xi(\rho) \simeq 1$ and the background evolution is radiation-like. We find that the asymptotic value of c_N is well approximated by a linear function of ζ ,

$$c_N \simeq 0.37\zeta + 1.79, \quad (\text{A.5})$$

from which we get the numerical results for the PQ number within $\mathcal{O}(10\%)$.

References

- [1] R.D. Peccei, H. Quinn, *Phys. Rev. Lett.* **38** (1977) 1440. *Phys. Rev. D* **16** (1977) 1791.
- [2] S. Weinberg, *Phys. Rev. Lett.* **40** (1978) 223. F. Wilczek, *ibid* **40** (1978) 279.
- [3] J.E. Kim, *Phys. Rev. Lett.* **43** (1979) 103. M. A. Shifman, A. I. Vainshtein, V. I. Zakharov, *Nucl. Phys. B* **166** (1980) 493.
- [4] M. Dine, W. Fischler, M. Srednicki, *Phys. Lett. B* **104** (1981) 199. A. R. Zhitnitsky, *Sov. J. Nucl. Phys.* **31** (1980) 260.
- [5] A. Caputo, G. Raffelt, *Astrophysical axion bounds: the 2024 edition*, 2024, [arXiv:2401.13728 \[hep-ph\]](#)
- [6] J. Preskill, M.B. Wise, F. Wilczek, *Phys. Lett.* **120B** (1983) 127. L. F. Abbott, P. Sikivie, *Phys. Lett.* **120B** (1983) 133. M. Dine, W. Fischler, *Phys. Lett.* **120B** (1983) 137.
- [7] R.T. Co, L.J. Hall, K. Harigaya, *Phys. Rev. Lett.* **124** (25) (2020) 251802. [arXiv:1910.14152 \[hep-ph\]](#)
- [8] A.G. Cohen, D.B. Kaplan, *Phys. Lett. B* **199** (1987) 251–258. *Nucl. Phys. B* **308** (1988) 913–928.
- [9] R.T. Co, K. Harigaya, *Phys. Rev. Lett.* **124** (11) (2020) 111602. [arXiv:1910.02080 \[hep-ph\]](#)
- [10] P. Minkowski, *Phys. Lett. B* **67** (1977) 421–428. T. Yanagida, *Conf. Proc. C* **7902131** (1979) 95–99. M. Gell-Mann, P. Ramond, R. Slansky, *Conf. Proc. C* **790927** (1979) 315–321. S. L. Glashow, *NATO Sci. Ser. B* **61** (1980) 687. R. N. Mohapatra, G. Senjanovic, *Phys. Rev. Lett.* **44** (1980) 912.
- [11] E.J. Chun, T.H. Jung, *Phys. Rev. D* **109** (9) (2024) 095004. [arXiv:2311.09005 \[hep-ph\]](#)
- [12] E.J. Chun, H.M. Lee, J.H. Song, 2026, [arXiv:2512.06413 \[hep-ph\]](#)
- [13] I. Affleck, M. Dine, *Title Missing*, *Nucl. Phys. B* **249** (1985) 361–380.
- [14] H.M. Lee, A.G. Menkara, M.J. Seong, J.H. Song, *Eur. Phys. J. C* **84** (12) (2024) 1260. [arXiv:2408.17013 \[hep-ph\]](#) H.M. Lee, A.G. Menkara, M.J. Seong, J.H. Song, *JHEP* **05** (2024) 295. [arXiv:2310.17710 \[hep-ph\]](#)
- [15] R.T. Co, N. Fernandez, A. Ghalsasi, L.J. Hall, K. Harigaya, *JHEP* **03** (2021) 017. [arXiv:2006.05687 \[hep-ph\]](#), [https://doi.org/10.1007/JHEP03\(2021\)017](#) P. Barnes, R. T. Co, A. Ghalsasi, K. Harigaya, A. Pierce, *JHEP* **08** (2025) 004. [arXiv:2402.10263 \[hep-ph\]](#), [https://doi.org/10.1007/JHEP08\(2025\)004](#)
- [16] J. Kawamura, S. Raby, *JHEP* **04** (2022) 116. [arXiv:2109.08605 \[hep-ph\]](#), [https://doi.org/10.1007/JHEP04\(2022\)116](#) P. Barnes, R. T. Co, K. Harigaya, A. Pierce, *JHEP* **05** (2023) 114. [arXiv:2208.07878 \[hep-ph\]](#), [https://doi.org/10.1007/JHEP05\(2023\)114](#)
- [17] M. Berbig, 2025, [arXiv:2506.23290 \[hep-ph\]](#)
- [18] S. Davidson, E. Nardi, Y. Nir, *Leptogenesis*, *Phys. Rept.* **466** (2008) 105–177. [arXiv:0802.2962 \[hep-ph\]](#)
- [19] S. Borsanyi, Z. Fodor, J. Guenther, K.H. Kampert, S.D. Katz, T. Kawanai, T.G. Kovacs, S.W. Mages, A. Pasztor, F. Pittler, et al., *Calculation of the axion mass based on high-temperature lattice quantum chromodynamics*, *Nature* **539** (7627) (2016) 69–71. [arXiv:1606.07494 \[hep-lat\]](#)
- [20] V. Domcke, Y. Ema, K. Mukaida, M. Yamada, *Spontaneous baryogenesis from axions with generic couplings*, *JHEP* **08** (2020) 096. [arXiv:2006.03148 \[hep-ph\]](#)
- [21] S. Navas et al., [Particle Data Group], *Phys. Rev. D* **110** (3) (2024) 030001. [https://doi.org/10.1103/PhysRevD.110.030001](#)
- [22] R. Kallosh, A. Linde, *JCAP* **12** (2013) 006. [arXiv:1309.2015 \[hep-th\]](#); R. Kallosh, A. Linde, D. Roest, *JHEP* **11** (2013) 198. [arXiv:1311.0472 \[hep-th\]](#); M. Galante, R. Kallosh, A. Linde, D. Roest, *Phys. Rev. Lett.* **114** (14) (2015) 141302. [arXiv:1412.3797 \[hep-th\]](#); B. J. Broy, M. Galante, D. Roest, A. Westphal, *JHEP* **12** (2015) 149. [arXiv:1507.02277 \[hep-th\]](#); T. Terada, *Phys. Lett. B* **760** (2016) 674–680. [arXiv:1602.07867 \[hep-th\]](#)
- [23] S. Clery, H.M. Lee, A.G. Menkara, *JHEP* **10** (2023) 144. [arXiv:2306.07767 \[hep-ph\]](#)
- [24] M. Kamionkowski, J. March-Russell, *Phys. Lett. B* **282** (1992) 137–141. [arXiv:hep-th/9202003 \[hep-th\]](#) S. M. Barr, J. March-Russell, *Phys. Rev. D* **46** (1992) 539–549.
- [25] M.A.G. Garcia, K. Kaneta, Y. Mambrini, K.A. Olive, *JCAP* **04** (2021) 012. [arXiv:2012.10756 \[hep-ph\]](#)
- [26] M.A.G. Garcia, M. Gross, Y. Mambrini, K.A. Olive, M. Pierre, J.H. Yoon, *JCAP* **12** (2023) 028. [arXiv:2308.16231 \[hep-ph\]](#), [https://doi.org/10.1088/1475-7516/2023/12/028](#)