



QCD axion dark matter from long-lived domain walls during matter domination

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ABSTRACT

The domain wall problem of the Peccei–Quinn mechanism can be solved if the Peccei–Quinn symmetry is explicitly broken by a small amount. Domain walls decay into axions, which may account for dark matter of the universe. This scheme is however strongly constrained by overproduction of axions unless the phase of the explicit breaking term is tuned. We investigate the case where the universe is matter-dominated around the temperature of the MeV scale and domain walls decay during this matter dominated epoch. We show how the viable parameter space is expanded.

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1. Introduction

The standard model has a source of CP violation from the QCD dynamics [1], which has however not been observed [2]. This is so-called the strong CP problem. One of the most attractive solutions to it is provided by the Peccei–Quinn (PQ) mechanism [3,4], where anomalous $U(1)_{\text{PQ}}$ symmetry is introduced. The mechanism predicts the existence of a light pseudo-scalar field called axion [5,6], which is also a candidate of dark matter in the universe [7–9].

The axion model is however not always cosmologically safe. The QCD dynamics explicitly break the $U(1)_{\text{PQ}}$ symmetry down to a discrete subgroup $Z_{N_{\text{DW}}}$. If the subgroup is non-trivial (i.e. $N_{\text{DW}} > 1$), there exist stable domain walls which eventually dominate the energy density of the universe [10,11]. The domain problem can be avoided if the PQ symmetry is already broken during inflation, but this scenario requires a small inflation scale to suppress the isocurvature perturbation [12–18] or a flat potential of the PQ symmetry breaking field so that the effective PQ symmetry breaking scale is large during inflation [18].

The domain wall problem can be solved if the $Z_{N_{\text{DW}}}$ symmetry is also explicitly broken and domain walls are unstable [11]. Actually the PQ “symmetry” is not at all symmetry since it is explicitly broken by the quantum anomaly. There would be no wonder that the PQ symmetry is also explicitly broken by a small amount at the

classical level, and there is no residual $Z_{N_{\text{DW}}}$ symmetry. For example, if the PQ symmetry is accidental symmetry as a result of other exact symmetry [19–30], we expect explicit PQ symmetry breaking by higher dimensional operators, which may also break the $Z_{N_{\text{DW}}}$ symmetry.¹

The unstable domain wall mainly decays into axions, which can explain the observed amount of dark matter [40–42] for a range of the decay constant f_a much smaller than the one required for the misalignment mechanism [7–9] ($f_a > 10^{11}$ GeV). Such a small decay constant is relevant for several future searches for solar axions [31–33] and halo axions [34–37]. For a recent review of axion searches see e.g. [38]. This scheme is however strongly constrained by non-observation of the effect of the strong CP phase. The explicit PQ symmetry breaking gives a small axion mass in addition to the one given by the QCD strong dynamics, and hence the cancellation of the strong CP phase becomes incomplete. The large enough explicit PQ symmetry breaking to obtain the dark matter abundance tends to yield too much strong CP phase unless f_a is close to the astrophysical lower bound, or the phase of the explicit symmetry breaking term is accidentally small.

In the previous works [40–42] it is assumed that the universe is radiation-dominated after the QCD phase transition. In this letter

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¹ This seems to be difficult to achieve in models with one-step symmetry breaking, as the exact symmetry results in stable topological defects. Models with several symmetry breakings like the one in Ref. [22] can have a desired property. Ref. [39] considers one-field models, but the symmetry imposed there is anomalous and needs further completion of the model.

we investigate the case where the universe is matter-dominated around the temperature of the MeV scale and domain walls decay during this matter dominated epoch. Such a cosmological scenario is actually expected if there is a very weakly coupled field such as a moduli field. The matter dominance results in the dilution of the axions emitted from domain walls, and hence the required magnitude of the explicit PQ symmetry breaking becomes smaller. We show how the viable parameter space is extended.

2. Axions from long-lived domain walls

We assume that the PQ symmetry is restored at some time after inflation and is spontaneously broken later. This includes the cases where the Hubble scale during inflation is larger than the PQ symmetry breaking scale and thermal/non-thermal effects restore the PQ symmetry. Once the $U(1)_{\text{PQ}}$ symmetry is broken, cosmic strings are produced [43]. The reconnection and the decay of strings maintain the number of strings per horizon size roughly one, which is so-called the scaling law [44,45].

As the temperature of the universe becomes as low as the QCD phase transition temperature, the explicit breaking of the PQ symmetry by the QCD dynamics becomes effective, and domain walls are formed between the strings, with the wall-tension given by [41,49]

$$\sigma_{\text{dw}} \simeq 9m_a f_a^2, \quad (1)$$

where m_a is the axion mass. In the parameter space of interest the domain walls decay after the QCD phase transition. Thus we consider the axion mass at the zero-temperature given by [5]

$$m_a \simeq 6 \text{ meV} \frac{10^9 \text{ GeV}}{f_a}. \quad (2)$$

If there remains a discrete subgroup $Z_{N_{\text{DW}}}$ ($N_{\text{DW}} > 1$) of the PQ symmetry, the domain wall-string network is stable. This is for example the case with the DFSZ model [46,47]. The network roughly follows the scaling-law [40,48], with the energy density given by

$$\rho_{\text{dw}}(t) \simeq \mathcal{A} \frac{\sigma_{\text{dw}}}{t}. \quad (3)$$

Here $\mathcal{A}$ is an $O(1)$ factor which can be determined by numerical simulations. The energy density of the network decreases to the second power of the scale factor of the universe, and eventually dominates the universe. This is inconsistent with the success of the standard cosmology, and is called the domain wall problem [11].

The problem may be solved if the $Z_{N_{\text{DW}}}$ symmetry is explicitly broken so that the domain walls are unstable. We parametrize the effect of the explicit breaking by the following bias term

$$V = -2\Xi v^4 \cos\left(\frac{a}{v} + \delta\right), \quad (4)$$

where $v = N_{\text{DW}} f_a$ is the PQ symmetry breaking scale, δ and Ξ are constants. We use the phase convention where the QCD non-perturbative effect creates the minima at $a = 2\pi k \times f_a$ ($k = 0, 1 - N_{\text{DW}} - 1$). This bias term resolves the degeneracy between the N_{DW} minima, and puts pressure on domain walls. The network collapses around the time t_{dec} when the pressure beats the tension of domain walls [42],

$$t_{\text{dec}} \simeq C_d \frac{\mathcal{A} \sigma_{\text{dw}}}{\Xi v^4 (1 - \cos(2\pi/N_{\text{DW}}))}, \quad (5)$$

where C_d is a numerical constant.

The energy of the domain walls is transferred into axions. The resultant number density of the axion is given by

$$n_a(t_{\text{dec}}) \simeq \frac{\rho_{\text{dw}}(t_{\text{dec}})}{\tilde{\epsilon}_a m_a}, \quad (6)$$

where $\tilde{\epsilon}_a$ is the average energy of the radiated axions normalized by the axion mass. We find that Eq. (6) reproduces the number density estimated by numerically solving the equation of motion of the domain wall energy density and the axion number density with an accuracy of few ten percents, if we determine t_{dec} (i.e. C_d) by the “10% criteria” in Ref. [42].

In order to suppress the axion abundance, the bias term must be sufficiently large so that the domain walls decay early enough. The bias term, which explicitly breaks the PQ symmetry, results in a non-zero strong CP phase. In the limit where the bias term is much smaller than the potential given by the QCD non-perturbative effect, the phase is given by

$$\theta \simeq -2N_{\text{DW}}^3 \Xi \sin \delta \frac{f_a^2}{m_a^2}. \quad (7)$$

The relation between the strong CP phase and the neutron electric dipole moment is estimated by various methods and approximation [50–56], and the estimation are different from each others by a factor of $O(10)$. This results in the uncertainty in the upper bound on θ . We conservatively accept the bound of $|\theta| < 10^{-10}$.

2.1. Axion abundance without entropy production

Now we evaluate the axion abundance and the required magnitude of Ξ , assuming that the universe is radiation-dominated below the QCD phase transition temperature. The axion abundance is estimated as [42]

$$\begin{aligned} \frac{\rho_a}{s} &\simeq 7 \times 10^{-10} \text{ GeV} \left(\frac{10^9 \text{ GeV}}{f_a} \right)^{1/2} \left(\frac{10^{-50}}{\Xi} \right)^{1/2} \\ &\times \frac{\mathcal{A}^{3/2} C_d^{1/2}}{\tilde{\epsilon}_a N_{\text{DW}}^2 \sin(\pi/N_{\text{DW}})}, \end{aligned} \quad (8)$$

where s is the entropy density. The observed dark matter abundance, $\rho_{\text{DM}}/s \simeq 4 \times 10^{-10} \text{ GeV}$, is reproduced with the magnitude of the explicit PQ symmetry breaking given by

$$\Xi \simeq 2 \times 10^{-50} \frac{10^9 \text{ GeV}}{f_a} \frac{\mathcal{A}^3 C_d}{\tilde{\epsilon}_a^2 N_{\text{DW}}^4 \sin^2(\pi/N_{\text{DW}})}, \quad (9)$$

which yields a non-zero strong CP phase,

$$\theta \simeq 9 \times 10^{-10} \sin \delta \left(\frac{f_a}{10^9 \text{ GeV}} \right)^3 \frac{\mathcal{A}^3 C_d}{\tilde{\epsilon}_a^2 N_{\text{DW}} \sin^2(\pi/N_{\text{DW}})}. \quad (10)$$

Using Eq. (9), the temperature at which the domain walls decay is given by

$$T_{\text{dec,DW}} \simeq 30 \text{ MeV} \frac{f_a}{10^9 \text{ GeV}} \frac{\mathcal{A}}{\tilde{\epsilon}_a}. \quad (11)$$

In Fig. 1, we show the prediction on the strong CP phase by a red solid line. We list the numerical values of $\mathcal{A}$, C_d and $\tilde{\epsilon}_a$ taken from [42] in Table 1. For $N_{\text{DW}} = 2$, to satisfy the upper bound $|\theta| < 10^{-10}$, $f_a < 8 \times 10^8 \text{ GeV}$ is required for $\sin \delta \sim 1$. Such parameter region is disfavored by the cooling of white dwarfs [57], if there exists a significant axion–electron coupling as is the case with the DFSZ model. Accidentally small δ allows for larger f_a . For example,

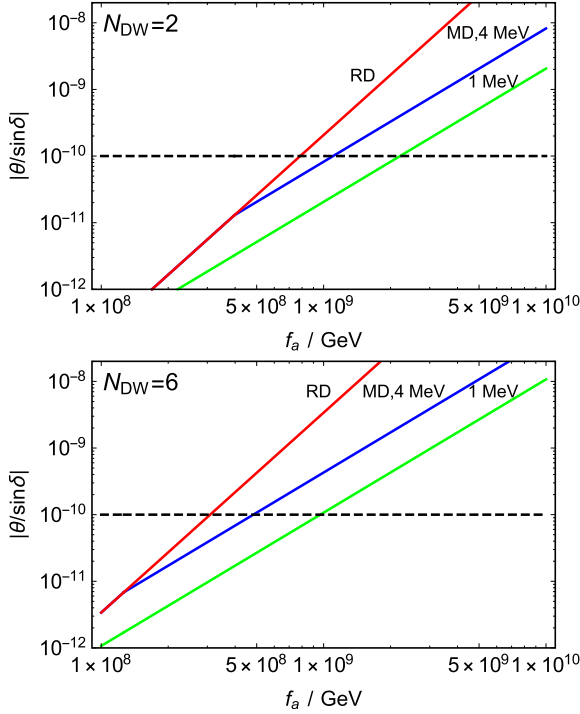


Fig. 1. The prediction on the strong CP phase as a function of f_a . The red lines assume that the universe is radiation-dominated after the QCD phase transition. The blue (green) lines assume that the universe is matter-dominated when the domain wall decays, and the matter decays at the temperature of 4 (1) MeV. (For interpretation of the colors in the figure(s), the reader is referred to the web version of this article.)

Table 1
Numerical values taken from [42].

	$\mathcal{A}$	C_d	$\tilde{\epsilon}_a$
$N_{\text{DW}} = 2$	0.7	5.3	2.0
$N_{\text{DW}} = 6$	2.2	2.1	2.0

if $\delta < 0.1$, $f_a < 2 \times 10^9$ GeV is allowed.² The constraint is stronger for $N_{\text{DW}} = 6$. This is because of a larger $\mathcal{A}$ for larger N_{DW} and hence larger domain wall energy, which can be understood by the fact that more domain walls are attached to a string. To satisfy the bound on the strong CP phase, $f_a < 3 \times 10^8$ GeV is required for $\sin\delta \sim 1$.

2.2. Axion abundance with entropy production

As is shown in Eq. (11), for the parameters which reproduce the observed dark matter abundance, the domain walls decay above the temperature of the MeV scale. The era is before the onset of the Big-Bang Nucleosynthesis, and the universe may be matter-dominated. Then the axion abundance from the domain walls is smaller than the simple case of the radiation dominance and the predicted strong CP phase is suppressed.

Actually if there exists a light, very weakly interacting particle, after the inflaton decays and the universe becomes radiation-dominated, the particle eventually dominates the energy density of the universe. We call such a particle as a modulus, although it is not necessarily a scalar particle. For example, a particle with a mass m which couples to the standard model particle with a

dimension-5 operator suppressed by a scale M_* decays around the temperature

$$T_{\text{dec},m} \sim 10 \text{ MeV} \left(\frac{m}{10^5 \text{ GeV}} \right)^{3/2} \frac{10^{18} \text{ GeV}}{M_*}. \quad (12)$$

A gravitationally coupled particle ($M_* \sim 10^{18}$ GeV) with a mass around 100 TeV gives $T_{\text{dec},m} \sim \text{MeV}$. In the supersymmetric standard model with the supersymmetry breaking mediated by gravitational interaction, the masses of moduli fields are as large as the masses of the scalar partners of the standard model fermions. The scalar masses of $O(100)$ TeV can explain the observed Higgs mass of 125 GeV [58–62]. If $M_* \sim f_a$, which is the case for the radial direction of the PQ symmetry breaking field, $m = O(0.1 - 1)$ GeV yields a MeV scale $T_{\text{dec},m}$.

When the domain walls decay, the ratio of the number density of the axion to the energy density of the moduli, ρ_m , is given by

$$\frac{n_a}{\rho_m} = \frac{3n_a(t_{\text{dec}})}{4M_{\text{Pl}}/t_{\text{dec}}^2}, \quad (13)$$

which does not change under the expansion of the universe until the moduli decays. Converting ρ_m to the entropy density when the moduli decay (see the appendix), we obtain

$$\frac{\rho_a}{s} \simeq 1.2 T_{\text{dec},m} m_a \frac{n_a}{\rho_m} = 4 \times 10^{-11} \text{ GeV} \left(\frac{10^9 \text{ GeV}}{f_a} \right)^2 \times \frac{10^{-50}}{\Xi} \frac{T_{\text{dec},m}}{1 \text{ MeV}} \frac{\mathcal{A}^2 C_d}{\tilde{\epsilon}_a N_{\text{DW}}^4 \sin^2(\pi/N_{\text{DW}})}, \quad (14)$$

where $T_{\text{dec},m}$ is defined by the decay rate of the moduli Γ as

$$\Gamma = 3H(T_{\text{dec},m}), \quad (15)$$

with the Hubble scale evaluated by that during radiation dominated era. The dark matter abundance is obtained for

$$\Xi \simeq 1 \times 10^{-51} \left(\frac{10^9 \text{ GeV}}{f_a} \right)^2 \frac{T_{\text{dec},m}}{1 \text{ MeV}} \frac{\mathcal{A}^2 C_d}{\tilde{\epsilon}_a N_{\text{DW}}^4 \sin^2(\pi/N_{\text{DW}})}, \quad (16)$$

which yields a non-zero strong CP phase,

$$\theta \simeq 3 \times 10^{-11} \sin\delta \left(\frac{f_a}{10^9 \text{ GeV}} \right)^2 \frac{T_{\text{dec},m}}{1 \text{ MeV}} \frac{\mathcal{A}^2 C_d}{\tilde{\epsilon}_a N_{\text{DW}} \sin^2(\pi/N_{\text{DW}})}. \quad (17)$$

The temperature $T_{\text{dec},m}$ is constrained to be $T_{\text{dec},m} > 0.7$ MeV, as the low temperature leads to incomplete thermalization of neutrinos, affecting the Big-Bang nucleosynthesis [63].³ This also reduces the energy density of the relativistic component of the universe, which changes the expansion history of the universe and affects the spectrum of the cosmic microwave background. The Planck satellite obtains the constraint $N_{\text{eff}} > 2.8$ (95% C.L.) [64]. Comparing this with the results in [63], we obtain $T_{\text{dec},m} > 4$ MeV. This bound can be evaded if the moduli has a significant branching fraction to radiations such as neutrinos or axions. We consider two reference values $T_{\text{dec},m} = 1$ and 4 MeV.

The prediction on the strong CP phase is shown in Fig. 1 by the solid blue and green lines, with $T_{\text{dec},m} = 4$ and 1 MeV, respectively. For $N_{\text{DW}} = 2$ with $T_{\text{dec},m} = 1$ MeV, the fine-tuning of δ is not necessary as long as $f_a < 2 \times 10^9$ GeV. For that large f_a , fine-tuning of about 10% is required without the matter domination. Allowing 10% tuning, f_a may be as large as 7×10^9 GeV.

² Ref. [42] assumes an aggressive bound $|\theta| < 7 \times 10^{-12}$ and finds much more restricted results.

³ If the moduli decays into hadrons with significant hadronic branch the BBN gives more stringent constraint as $T_{\text{dec},m} > 3\text{--}4$ MeV [63].

3. Summary and discussion

We comment on the production mechanisms of axion dark matter for a small decay constant discussed in the literature. If the misalignment angle is very close to the maximal one, the unharmonic effect delays the commencement of the oscillation of the axion. This enhances the axion abundance and reproduces the observed dark matter abundance even if $f_a \ll 10^{11}$ GeV [16,65,66]. This scenario requires the fine-tuning of the misalignment angle, as well as a very small inflation scale to suppress the isocurvature perturbation. If the radial direction of the PQ symmetry breaking field takes a large value in the early universe, the production of the axions by the parametric resonance [67,68] can produce axion dark matter [69]. The scenario requires that the potential of the PQ symmetry breaking field is flat, which would call for supersymmetry. Kinetic energy domination around the QCD phase transition enhances the axion abundance produced by the misalignment mechanism and allows for a small decay constant [70].

In this paper we have investigated the production of axions from unstable domain walls because of the explicit PQ symmetry breaking, assuming that the universe is matter-dominated around the temperature of the MeV scale and domain walls decay during this matter dominated epoch. This scenario can explain the observed amount of dark matter by the axion without tuning the phase of the explicit breaking as long as $f_a < O(10^9)$ GeV. Allowing the tuning of $O(10)\%$, $f_a < O(10^{10})$ GeV is allowed.

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Appendix A. Derivation of Eq. (14)

The equations governing the evolution of the energy densities of the moduli, ρ_m , and radiation, ρ_r , and the scalar factor, a , is given by

$$\dot{\rho}_m + 3H\rho_m = -\Gamma\rho_m, \quad (18)$$

$$\dot{\rho}_r + 4H\rho_r = +\Gamma\rho_m, \quad (19)$$

$$\frac{\dot{a}}{a} = H = \frac{1}{\sqrt{3}M_{\text{Pl}}} \sqrt{\rho_m + \rho_r}, \quad (20)$$

with the initial conditions given at $t_i \ll 1/\Gamma$,

$$\rho_m(t = t_i) \equiv \rho_{m,i}, \quad (21)$$

$$\rho_r(t = t_i) = 0, \quad (22)$$

$$a(t = t_i) \equiv a_i \quad (23)$$

After the change of variables

$$t = x/\Gamma, \quad \rho_m = a^{-3}m\Gamma^2 M_{\text{Pl}}^2, \quad \rho_r = a^{-4}r\Gamma^2 M_{\text{Pl}}^2, \quad (24)$$

the equations are given by

$$\frac{dm}{dx} = -m, \quad (25)$$

$$\frac{dr}{dx} = am, \quad (26)$$

$$\frac{da}{dx} = \frac{a}{\sqrt{3}} \sqrt{a^{-3}m + a^{-4}r}. \quad (27)$$

The solution for m is

$$m = e^{-(x-x_i)} m_i, \quad (28)$$

where $m_i \equiv m(x_i)$. The evolution equations of r and a are given by

$$\frac{dr}{dx} = ae^{-(x-x_i)} m_i, \quad (29)$$

$$\frac{da}{dx} = \frac{a}{\sqrt{3}} \sqrt{a^{-3}e^{-(x-x_i)} m_i + a^{-4}r}. \quad (30)$$

For $x \rightarrow \infty$, r becomes a constant $\equiv r_f$.

The quantity we are interested in is

$$\frac{n_a}{s}(t \rightarrow \infty) = \frac{n_a}{\rho_r^{3/4}}(t \rightarrow \infty) \frac{\rho_r^{3/4}}{s} = \frac{n_a}{\rho_m}(t = t_i) \frac{m_i}{r_f^{3/4}} \sqrt{\Gamma M_{\text{Pl}}} \frac{\rho_r^{3/4}}{s}. \quad (31)$$

Using the definition of $T_{\text{dec},m}$ in Eq. (15), we obtain

$$\frac{n_a}{s}(t \rightarrow \infty) = \frac{n_a}{\rho_m}(t = t_i) T_{\text{dec},m} \frac{m_i}{r_f^{3/4}} \frac{3^{5/4}}{4}. \quad (32)$$

A numerical evaluation of Eqs. (29) and (30) gives $m_i/r_f^{3/4} \simeq 1.2$, and we obtain

$$\frac{n_a}{s}(t \rightarrow \infty) \simeq 1.2 \times T_{\text{dec},m} \frac{n_a}{\rho_m}(t = t_i). \quad (33)$$

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