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Relaxing isocurvature bounds on string axion dark matter

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ABSTRACT

If inflation scale is high, light scalars acquire large quantum fluctuations during inflation. If sufficiently long-lived, they will give rise to CDM isocurvature perturbations, which are highly constrained by the Planck data. Focusing on string axions as such light scalars, we show that thermal inflation can provide a sufficiently large entropy production to dilute the CDM isocurvature perturbations. Importantly, efficient dilution is possible for the string axions, because effectively no secondary coherent oscillations are induced at the end of thermal inflation, in contrast to the moduli fields. We also study the viability of the axion dark matter with mass of about 7 keV as the origin of the 3.5 keV X-ray line excess, in the presence of large entropy production.

1. Introduction

Inflation [1] elegantly solves theoretical problems of the standard cosmology such as the horizon problem,[1] and the slow-roll inflation paradigm is consistent with the observations including the comic microwave background (CMB) and large-scale structure data. Measuring the primordial B-mode polarization of CMB is therefore of crucial importance, as it would provide a definitive proof of inflation [6].

Recently the BICEP2 collaboration reported detection of the B-mode polarization, which could be due to the primordial gravitational waves [7]. If this is the case, the BICEP2 results can be explained by the tensor mode perturbations with a tensor-to-scalar ratio, $r = 0.20^{+0.07}_{-0.05}$, whereas the central value (and therefore significance for the signal) depends on models for foreground dust polarization [8].

Taken at face value, the BICEP2 result strongly suggests high-scale inflation with the Hubble parameter, $H_{\rm inf} \sim 10^{14}$ GeV [7]. Importantly, any light scalar particles acquire large quantum fluctuations of order $H_{\rm inf}/2\pi$ during inflation. Those scalars are copiously produced by coherent oscillations when their mass becomes comparable to the Hubble parameter after inflation. If some of them are long-lived, the coherent oscillations will contribute to dark matter, giving rise to isocurvature perturbations, which are tightly constrained by the CMB observations [9].

The QCD axion, a pseudo Nambu–Goldstone boson associated with the spontaneous break down of the Peccei–Quinn (PQ) symmetry [10,11], is an ideal candidate for such light scalars. As is well known, the QCD axion can explain the observed dark matter density for the decay constant $f_a = \mathcal{O}(10^{11-12})$ GeV, barring fine-tuning of the initial misalignment angle. The isocurvature bounds on the QCD axion dark matter were extensively studied in the literature (see Refs. [12–14] for early works, Refs. [15–17] for non-Gaussianity of isocurvature perturbations, and Refs. [19–23] for the recent works after BICEP2). The upper bound on the inflation scale is roughly given by $H_{\rm inf} \lesssim 10^7\ {\rm GeV}(f_a/10^{11}\ {\rm GeV})^{0.408}$ [9], which shows clear tension with high-scale inflation. There have been proposed several solutions to the tension between the high-scale inflation and the QCD axion dark matter. The simplest solution is to assume that the PQ symmetry is restored [24,25]. Another one is to make the QCD axion sufficiently heavy during inflation [18,19]. Alternatively, it is also known that the axion quantum fluctuations can be suppressed or enhanced if the kinetic term coefficient of the phase of the PQ scalar evolves during and after inflation. The axion isocurvature perturbations are suppressed if the effective PQ breaking scale decreases after inflation by the saxion dynamics [24,26]. Recently the saxion dynamics was studied in Ref. [27] taking account of the parametric resonance effects. A similar suppression was found in the presence of an enhanced non-minimal coupling to the Einstein tensor [28]. Another possibility is to modify the axion potential by introducing other shift

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[1] The exponentially expanding universe was also studied in Refs. [2–5].

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symmetry breaking terms. In particular, if the axion potential can be approximated by the quartic coupling in a certain range of the potential, the axion density decreases in proportional to a^{-4} as radiation, where a is the scale factor. Then the axion abundance as well as the associated isocurvature perturbations can be suppressed at the expense of the fine-tuning of the axion potential.[2]

If there are other light scalars, we can similarly apply the isocurvature constraints. A plausible candidate for such light scalars is string axions, the imaginary components of the moduli fields [30]. There appear many moduli fields through compactifications in the string theory. In order to have a sensible low-energy theory, those moduli fields must be stabilized properly. Many of them can be stabilized with a heavy mass because of the fluxes [31,32], whereas the remaining ones can be stabilized by non-perturbative superpotentials induced by gaugino condensation and/or stringy instantons [33]. In particular, the axions respect the following axionic shift symmetry,

$$a \to a + C, \tag{1}$$

where C is a real transformation parameter, and therefore, some of them may remain relatively light and play an important cosmological role.

Indeed, depending on the nature of the shift-symmetry breaking, the axion mass can be extremely light [34,35]. Then such light axions become so long-lived that they contribute to dark matter.

In this Letter we study a possibility to solve the tension between the string axion dark matter and high-scale inflation suggested by BICEP2 by a late-time entropy production due to thermal inflation [36–41]. One of the main differences of the string axions from general moduli fields in this context is that effectively no secondary coherent oscillations are induced at the end of thermal inflation, because the string axions do not receive the so-called Hubble-induced mass. This greatly helps to dilute light axions efficiently, as the effect of the secondary coherent oscillations becomes prominent for light moduli fields. We will show that a large portion of the parameter space can be indeed consistent with the isocurvature bound in the presence of large entropy production by thermal inflation. We will also study the viability of axion dark matter with mass about 7 keV [42–44][3] as the origin of the 3.5 keV X-ray line excess [47,48] in the presence of such late-time entropy production.

2. Isocurvature perturbations of axions

We briefly discuss isocurvature perturbations induced by axion coherent oscillations. Throughout this Letter we assume a large axion decay constant, f_a, of order 10^{15} GeV, for which thermal production is negligible. Also neglected is non-thermal production by the saxion decays [49–51]. In general both contributions do not induce isocurvature perturbations.

The potential of the axion is given by

$$V(a) = m_a^2 f_a^2 \left[1 - \cos\left(\frac{a}{f_a}\right)\right], \tag{2}$$

where the potential minimum is located at the origin. The axion potential can be well approximated as the quadratic potential with mass m_a in the vicinity of the minimum. The axion starts to oscillate about the potential minimum when the mass becomes comparable to the Hubble parameter, unless it is initially very close to the top of the potential. Assuming that it starts to oscillate in the radiation dominated era after reheating, we can write down the energy-to-entropy ratio, ρ_a/s, as

$$\frac{\rho_a}{s} = \frac{1}{8} T_{\rm osc} \left(\frac{a_{\rm osc}}{M_P}\right)^2, \tag{3}$$

where $a_{\rm osc}$ is the initial oscillation amplitude. The $a_{\rm osc}$ can be expressed in terms of the initial misalignment angle θ_i as

$$a_{\rm osc} = \max(f_a \theta_{\rm osc},\ H_{\rm inf}/2\pi) \tag{4}$$

$$\theta_{\rm osc}^2 \equiv \theta_i^2 F(\theta_i). \tag{5}$$

Here $F(\theta_i)$ represents the anharmonic effect for the axion coherent oscillation [52–54] given by

$$F(\theta_i) = \left[\ln\left(\frac{e}{1-\theta_i^2/\pi^2}\right)\right]^{3/2}, \tag{6}$$

which is unity for $\theta_i \ll \pi$, while it diverges as θ_i approaches π where the axion potential becomes maximum. Note that the power in Eq. (6) is different from the one adopted in [54] because there is no finite temperature effect on the axion mass in our case. The reason is as follows. The thermal effects on the axion mass appear only if the hidden gauge fields are in the thermal bath. We simply assume that the hidden gauge sector was not reheated by the inflaton decay. Then, the hidden gauge interactions become strong and the axion potential arises when the Hubble parameter becomes comparable to the dynamical scale of the hidden gauge symmetry. At that time, the axion mass is still less than the Hubble parameter as long as the decay constant is smaller than the Planck scale. Therefore, the axion mass is constant and does not depend on the cosmic temperature (of the visible sector) when it starts to oscillate.

Also we have defined $T_{\rm osc}$ as the temperature at the beginning of the axion oscillation:

$$T_{\rm osc} = \left(\frac{90}{\pi^2 g_*(T_{\rm osc})}\right)^{1/4} \sqrt{m_a M_P}, \tag{7}$$

where $g_*(T)$ is the relativistic degrees of freedom at cosmic temperature T, and M_P is the reduced Planck mass.

The present density parameter of the axion can then be calculated by using $\rho_{{\rm cr},0}/s_0 = 3.64 \times 10^{-9} h^2$ GeV as following;

$$\Omega_a h^2 \simeq \begin{cases} 4 \times (\frac{228.75}{g_*(T_{\rm osc})})^{1/4} (\frac{m_a}{10^{-9}\ {\rm eV}})^{1/2} (\frac{f_a}{10^{15}\ {\rm GeV}})^2 \theta_i^2 F(\theta_i) \\ \quad \text{for } \theta_i > \theta_c \\ 1 \times 10^{-3} (\frac{228.75}{g_*(T_{\rm osc})})^{1/4} (\frac{m_a}{10^{-9}\ {\rm eV}})^{1/2} (\frac{H_{\rm inf}}{10^{14}\ {\rm GeV}})^2 \\ \quad \text{for } \theta_i < \theta_c \end{cases} \tag{8}$$

where $\theta_c \approx 0.016 (H_{\rm inf}/10^{14}\ {\rm GeV})(10^{15}\ {\rm GeV}/f_a)$ represents the critical value where the quantum fluctuation δa becomes equal to the classical field deviation $\theta_i f_a$. Note that the high-scale inflation implies the axion overproduction unless the axion mass is extremely light such as, $m_a \lesssim 10^{-9}$ eV, for $f_a \sim 10^{15}$ GeV.

In linear perturbation theory, the power spectrum of the axion isocurvature perturbation can be expressed in terms of the axion fluctuation as follows [55];

$$\mathcal{P}_{\mathcal{S},a}^{1/2} = \frac{\delta\rho_a}{\rho_a} = \frac{2 f_a \theta_{\rm osc} (\frac{\partial \theta_{\rm osc}}{\partial \theta_i}) \delta a + (\delta a)^2}{(f_a \theta_{\rm osc})^2 + (\delta a)^2} \tag{9}$$

where axion field fluctuation is given by $\delta a = H_{\rm inf}/2\pi$, and we have taken account of the anharmonic effect on the isocurvature

[2] A similar fine-tuning of the potential enables the axion hilltop inflation [29].

[3] See also Refs. [40,45,46] for the early works on the X-ray constraint on a light modulus field.