Axion isocurvature fluctuations with extremely blue spectrum

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We construct an axion model for generating isocurvature fluctuations with blue spectrum, $n_{\rm iso} = 2$–4, which is suggested by recent analyses of admixture of adiabatic and isocurvature perturbations with independent spectral indices, $n_{\rm ad} \neq n_{\rm iso}$. The distinctive feature of the model is that the spectrum is blue at large scales while scale invariant at small scales. This is naturally realized by the dynamics of the Peccei-Quinn scalar field.

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I. INTRODUCTION

Large-scale structures of the Universe, such as galaxies and clusters of galaxies, have formed through gravitational instabilities, initiated by the primordial seed density fluctuations, which were created during inflation. The simplest initial condition seeded these inhomogeneities is the (almost) scale-invariant adiabatic curvature perturbations. They can fit to very precise measurements of the cosmic microwave background temperature and polarization anisotropies, large-scale structures, and supernovae [1–3]. It is usually realized by the single-field inflation where the inflaton fluctuations are responsible for the adiabatic perturbations.

Generally, there will exist other light fields whose fluctuations during inflation become isocurvature perturbations [4]. Therefore, the admixture of isocurvature and adiabatic fluctuations could be what really happened in the early Universe. Observational analyses with the assumption that the spectral indices of the adiabatic and isocurvature fluctuations are the same, $n_{\rm ad} = n_{\rm iso} \simeq 1$, have revealed that the contributions from the isocurvature perturbation should be small [1,5].

However, there is *a priori* no reason for the isocurvature fluctuations to have (almost) scale-invariant spectrum. In fact, more general analyses with independent spectral indices of adiabatic and isocurvature modes based on recent observations result in the favor of much more contribution of the isocurvature component with an extremely blue tilt ($n_{\rm iso} \simeq 2$–4) [6–8].

In this article, we provide a concrete model for generating isocurvature fluctuations with an extremely blue spectrum for the first time.[1] It is the axion model in supersymmetry (SUSY) [10]. Since the axion is a good candidate of the cold dark matter of the Universe, and has nothing to do with radiation, it gives rise to uncorrelated isocurvature fluctuations. The axion isocurvature fluctuation is usually expressed as $\delta a/a \simeq H/(2\pi F_a \theta)$, where H is the Hubble parameter during inflation, F_a the axion decay constant, and θ a misalignment angle. The key to produce the blue spectrum is that we promote F_a as a dynamical field $\varphi \equiv$ "F_a", which initially has a large value $\simeq M_P$, evolves toward smaller values, and stops at $\simeq F_a$ during inflation. It is realized very simply and naturally in the SUSY axion model, and we can obtain extremely blue spectrum such as $n_{\rm iso} \simeq 4$ at large scales, which is connected to the scale-invariant spectrum at small scales.

The structure of the article is as follows: In the next section, we explain the essence to generate the extremely blue spectrum in a simple model, reduced from the concrete model that we provide based on SUSY in Sec. III. We then show the dynamics of the fields, which leads to the favorable spectrum in Sec. IV. Our conclusions are given in Sec. V.

II. HOW TO GET THE BLUE SPECTRUM

Let us consider a toy model of a complex scalar field Φ, whose energy density is negligible during inflation. Fluctuations in the phase direction give rise to an isocurvature perturbation, while fluctuations in the radial direction are negligibly small due to large effective mass in that direction as shown shortly. Thus, the isocurvarture fluctuation is given by

$$\frac{\delta\theta}{\theta} \simeq \frac{H}{2\pi\varphi\theta}, \tag{1}$$

where we denote $\Phi = \varphi e^{i\theta}/\sqrt{2}$. Since the Hubble parameter during inflation is (almost) constant, it is the decreasing amplitude of φ that makes the isocurvature perturbation blue tilted. When the field φ has mass of $O(H)$, it can roll down in the potential during inflation, and, in addition, its fluctuation $\delta\varphi$ is suppressed. The reduced potential is given by

$$V \simeq \tfrac{1}{2} c H^2 \varphi^2, \tag{2}$$

when φ has a large field value, and $c \sim O(1)$ is a constant.

[1]The possibility to obtain isocurvature fluctuations with some deviation from scale-invariant was investigated in Ref. [9].

Then the φ field obeys the equation

$$\ddot{\varphi} + 3H\dot{\varphi} + cH^2\varphi = 0, \tag{3}$$

which has a solution of the form $\varphi \propto e^{-\lambda Ht}$ with

$$\lambda = \frac{3}{2} - \frac{3}{2}\sqrt{1 - \frac{4}{9}c} \tag{4}$$

for $0 \le c \le 9/4$.[2] Since the isocurvature fluctuation is estimated as

$$\Delta^2_{\rm iso} \propto \left(\frac{\delta a}{\varphi}\right)^2 \sim \left(\frac{H}{\varphi}\right)^2 \propto e^{2\lambda Ht}, \tag{5}$$

its spectral index is given by

$$n_{\rm iso} - 1 \equiv \frac{d\ln\Delta^2_{\rm iso}}{d\ln k} = 2\lambda = 3 - 3\sqrt{1 - \frac{4}{9}c}. \tag{6}$$

Therefore, we obtain the blue spectrum, even extremely blue such as $n_{\rm iso} = 4$ for $c = 9/4$. As shown explicitly in the following sections, the field φ eventually settles down in the minimum of the potential placed at $\varphi \simeq F_a$. Thereafter the isocurvature flucutation becomes scale invariant.

III. AXION MODEL IN SUSY

The axion [11,12] is a Nambu-Goldstone boson associated with the Peccei-Quinn (PQ) symmetry breaking, and is the most natural solution to the strong CP problem in QCD [13]. The PQ symmetry breaking scale F_a is astrophysically and cosmologically constrained within the range of 10^{10}–10^{12} GeV [14]. The axion can be cold dark matter for the higher values.

Let us consider a concrete model of the axion in SUSY. We take the following superpotential [10]:

$$W = h(\Phi_+\Phi_- - F_a^2)\Phi_0. \tag{7}$$

Here, Φ_+, Φ_-, and Φ_0 are chiral superfields with PQ charges $+1$, -1, and 0, respectively, and h is a constant of $O(1)$. The scalar potential is obtained as

$$V_{\rm SUSY} = h^2|\Phi_+\Phi_- - F_a^2|^2 + h^2(|\Phi_+|^2 + |\Phi_-|^2)|\Phi_0|^2, \tag{8}$$

where we denote the scalar components with the same symbols as the superfields. One can easily see the existence of the flat direction, which satisfies

$$\Phi_+\Phi_- = F_a^2, \qquad \Phi_0 = 0. \tag{9}$$

In addition, SUSY breaking effects lift the flat direction by soft mass terms

$$V_{\rm m} = m_+^2|\Phi_+|^2 + m_-^2|\Phi_-|^2 + m_0^2|\Phi_0|^2, \tag{10}$$

at low energy scales, where m_+, m_-, and m_0 are of O(TeV), as well as the so-called Hubble-induced mass terms during inflation,

$$V_{\rm H} = c_+H^2|\Phi_+|^2 + c_-H^2|\Phi_-|^2 + c_0H^2|\Phi_0|^2, \tag{11}$$

where c_+, c_-, and c_0 are positive constants of $O(1)$, which stem from the supergravity effects [15,16].[3] Notice that no Hubble-induced A terms will appear due to PQ symmetry. We assume $H \ll F_a$ in order not to destroy the flat direction (9). Taking into account the fact that $\Phi_0 = 0$ and $m_i \ll H(i = \pm, 0)$ during inflation, we only consider the potential of the form

$$V = h^2|\Phi_+\Phi_- - F_a^2|^2 + c_+H^2|\Phi_+|^2 + c_-H^2|\Phi_-|^2. \tag{12}$$

Owing to the Hubble-induced mass terms, the minimum of the potential is given by

$$|\Phi_+^{\rm min}| \simeq \left(\frac{c_-}{c_+}\right)^{1/4} F_a, \qquad |\Phi_-^{\rm min}| \simeq \left(\frac{c_+}{c_-}\right)^{1/4} F_a. \tag{13}$$

Since it is symmetric between Φ_+ and Φ_-, we consider $|\Phi_+| > |\Phi_-|$ without loss of generality.[4]

Now we must identify the axion field a. Rewriting $\Phi_\pm$ as

$$\Phi_\pm \equiv \frac{1}{\sqrt{2}}\varphi_\pm \exp(i\theta_\pm), \qquad \theta_\pm \equiv \frac{a_\pm}{\sqrt{2}\varphi_\pm}, \tag{14}$$

we can define the fields a and b as

$$a = \frac{\varphi_+}{(\varphi_+^2 + \varphi_-^2)^{1/2}} a_+ - \frac{\varphi_-}{(\varphi_+^2 + \varphi_-^2)^{1/2}} a_-, \tag{15}$$

$$b = \frac{\varphi_-}{(\varphi_+^2 + \varphi_-^2)^{1/2}} a_+ + \frac{\varphi_+}{(\varphi_+^2 + \varphi_-^2)^{1/2}} a_-. \tag{16}$$

From Eqs. (8) or (12), the potential $V(b)$ for the field b is obtained as

$$V(b) = -h^2F_a^2\varphi_+\varphi_- \cos\left(\frac{(\varphi_+^2 + \varphi_-^2)^{1/2}}{\varphi_+\varphi_-} b\right), \tag{17}$$

which implies that the mass of b is given by $\sim h(\varphi_+^2 + \varphi_-^2)^{1/2}$. Since the field value is $\varphi_+ \simeq M_P$ initially, and decreases until it reaches to F_a, as shown in the next section, $m_b \gg H$ during inflation and hence the b field quickly settles down into the minimum of the potential. On the other hand, the potential for the a field is flat, and we can regard it as the axion. During inflation, the quantum

[2] One obtains the damping oscillating solution for $c > 9/4$. Since it does not suit for our purpose, we only consider for $c \le 9/4$.

[3] The coefficients of the Hubble-induced mass terms are determined as $c_i \simeq 3(1 - y_i)(i = 0, \pm)$ for the nonrenormalizable coupling in Kähler potential $\delta K = y_i|\Phi_i|^2|I|^2/M_P^2$, where I is the inflaton and y_i's are the coupling constants.

[4] Notice that, as shown shortly, the amplitude of the isocurvature fluctuation is determined by the larger between $|\Phi_+|$ and $|\Phi_-|$, so the spectrum cannot be red tilted.

fluctuations of a develop as

$$\delta a \simeq \delta a_+ \simeq \frac{H}{2\pi}, \tag{18}$$

where $\varphi_+ \gg \varphi_-$, while $\delta b \simeq 0$ because it is very massive, $m_b \gg H$. Thus, $\delta a_- \simeq -(\varphi_-/\varphi_+)\delta a_+$. Therefore, we have

$$\delta\theta_\pm = \frac{\delta a_\pm}{\sqrt{2}\varphi_\pm} \simeq \pm\frac{H}{2\sqrt{2}\pi\varphi_+}. \tag{19}$$

The crucial point is that the amplitude of the fluctuation is determined solely by the larger field value φ_+. Also notice that the fluctuations in the radial directions $\delta\varphi_+$ and $\delta\varphi_-$ are both suppressed due to large curvatures in their potentials, which stem from the first term in V_{SUSY} [Eq. (8)] and the Hubble-induced mass terms [Eq. (11)].

The axion isocurvature perturbation is given by[5]

$$\mathcal{S}_a \equiv \frac{\delta n_a}{n_a} - \frac{\delta n_\gamma}{n_\gamma} = 2\frac{\delta a}{a} \simeq \frac{H}{\sqrt{2}\pi\varphi_+\theta_+}, \tag{20}$$

where n_a and n_γ denote the number densities of the axion and photon, respectively, and we use Eqs. (14) and (18) in the last equality. Therefore, the isocurvature fluctuation is written as

$$\Delta^2_{\mathcal{S}}(k) = A_{\mathrm{iso}}\left(\frac{k}{k_0}\right)^{n_{\mathrm{iso}}-1}, \qquad A_{\mathrm{iso}} \simeq \frac{H^2}{2\pi^2\varphi_+^2\theta_+^2}\Bigg|_{k=k_0}. \tag{21}$$

Recent analyses of the admixture of adiabatic and isocurvature perturbations with independent spectral indices, $n_{\mathrm{ad}} \neq n_{\mathrm{iso}}$, reveal that the isocurvature contribution can be as large as the adiabatic mode at the pivot scale k_0, and the blue spectral index of the isocurvature fluctuation is favored such as $n_{\mathrm{iso}} \sim 4$ [8].

IV. DYNAMICS OF THE FIELDS AND ISOCURVATURE FLUCTUATIONS

As shown in the previous section, the amplitude of the isocurvature fluctuation is solely determined by the larger field value φ_+ with the constant Hubble parameter during inflation, $H \simeq$ const. We therefore need to investigate the dynamics of φ_+ only. It is reasonable to consider that the fields slide only along the flat direction, so that $\varphi_- = 2F_a^2/\varphi_+$, thus the potential can be approximated as

$$V \simeq \frac{1}{2}c_+H^2\varphi_+^2 + 2c_-H^2F_a^4\frac{1}{\varphi_+^2} \simeq \frac{1}{2}c_+H^2\varphi_+^2, \tag{22}$$

where the last equality holds when φ_+ has a large field value.[6] Now we must just follow the same argument discussed in Sec. II. Since the φ_+ field obeys the equation

$$\ddot{\varphi}_+ + 3H\dot{\varphi}_+ + c_+H^2\varphi_+ = 0, \tag{23}$$

whose solution is given by the form $\varphi_+ \propto e^{-\lambda Ht}$ with

$$\lambda = \frac{3}{2} - \frac{3}{2}\sqrt{1-\frac{4}{9}c_+}, \tag{24}$$

the isocurvature fluctuation is obtained as

$$\Delta^2_{\mathrm{iso}} \propto \left(\frac{\delta a}{\varphi_+}\right)^2 \sim \left(\frac{H}{\varphi_+}\right)^2 \propto e^{2\lambda Ht}, \tag{25}$$

so that its spectral index becomes

$$n_{\mathrm{iso}} - 1 \equiv \frac{d\ln\Delta^2_{\mathrm{iso}}}{d\ln k} = 2\lambda = 3 - 3\sqrt{1-\frac{4}{9}c_+}. \tag{26}$$

Therefore, we obtain the blue spectrum with $1 < n_{\mathrm{iso}} \leq 4$ for $0 < c_+ \leq 9/4$. The prominent feature of this model is that φ_+ eventually settles down to the minimum of the potential,

$$\varphi_+^{\min} \simeq \sqrt{2}\left(\frac{c_-}{c_+}\right)^{1/4}F_a, \qquad \varphi_-^{\min} \simeq \sqrt{2}\left(\frac{c_+}{c_-}\right)^{1/4}F_a, \tag{27}$$

and hence we have scale-invariant spectrum afterwards, smoothly connected from the blue spectrum at large scales. The e-folds during the field evolving from $\simeq M_P$ to $\simeq F_a$ are estimated as

$$N_{\mathrm{blue}} \simeq \frac{1}{\lambda}\ln\left(\frac{M_P}{F_a}\right), \tag{28}$$

which gives $N_{\mathrm{blue}} \simeq 10$ for $F_a = 10^{12}$ GeV and $\lambda = 3/2$ ($c_+ = 9/4$), for example.

In order to confirm what we have obtained above, we solve numerically the equations for Φ_+ and Φ_- with the potential (12). For the sake of numerical calculations, we decompose the field into real and imaginary parts as $\Phi_\pm = (\phi^R_\pm + i\phi^I_\pm)/\sqrt{2}$, which leads to the following equations:

$$\ddot{\phi}^R_+ + 3H\dot{\phi}^R_+ + c_+H^2\phi^R_+ + \frac{h^2}{2}[\{\phi^R_+\phi^R_- - \phi^I_+\phi^I_- - 2F_a^2\}\phi^R_- + (\phi^R_+\phi^I_- + \phi^I_+\phi^R_-)\phi^I_-] = 0, \tag{29}$$

$$\ddot{\phi}^I_+ + 3H\dot{\phi}^I_+ + c_+H^2\phi^I_+ + \frac{h^2}{2}[-\{\phi^R_+\phi^R_- - \phi^I_+\phi^I_- - 2F_a^2\}\phi^I_- + (\phi^R_+\phi^I_- + \phi^I_+\phi^R_-)\phi^R_-] = 0, \tag{30}$$

$$\ddot{\phi}^R_- + 3H\dot{\phi}^R_- + c_-H^2\phi^R_- + \frac{h^2}{2}[\{\phi^R_+\phi^R_- - \phi^I_+\phi^I_- - 2F_a^2\}\phi^R_+ + (\phi^R_+\phi^I_- + \phi^I_+\phi^R_-)\phi^I_+] = 0, \tag{31}$$

[5] The actual observable is $\mathcal{S}_{\mathrm{CDM}} = (\Omega_a/\Omega_{\mathrm{CDM}})\mathcal{S}_a \propto \theta_+$, where the axion density parameter is $\Omega_a \propto \theta_+^2$.

[6] For large φ_+, kinetic terms are reduced to the normal one as $\frac{1}{2}\sum_{i=\pm}\partial_\mu\varphi_i\partial^\mu\varphi_i = \frac{1}{2}(1+\frac{4F_a^4}{\varphi_+^4})\partial_\mu\varphi_+\partial^\mu\varphi_+ \simeq \frac{1}{2}\partial_\mu\varphi_+\partial^\mu\varphi_+$.

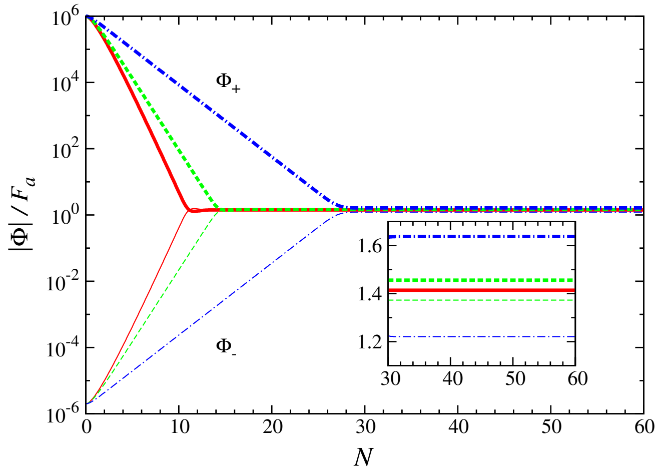

FIG. 1 (color online). Evolution of the fields Φ_+ (upper thick lines) and Φ_- (lower thin lines) for $c_- = 9/4$ and $c_+ = 9/4$ ($n_{\rm iso} = 4$, red solid), 2 ($n_{\rm iso} = 3$, green dashed), and $5/4$ ($n_{\rm iso} = 2$, blue dotted-dashed). The inset shows the minima where the fields settle down.

$$\ddot{\phi}^I_- + 3H\dot{\phi}^I_- + c_- H^2 \phi^I_- + \frac{h^2}{2}[-\{\phi^R_+\phi^R_- - \phi^I_+\phi^I_- - 2F_a^2\}\phi^I_+ + (\phi^R_+\phi^I_- + \phi^I_+\phi^R_-)\phi^R_+] = 0. \quad (32)$$

Some of the examples are shown in Fig. 1. Here, we set the initial condition as $|\Phi_+(0)| = M_P$. The initial values for Φ_- and the phases $\theta_\pm$ are taken so as to stay along the flat direction (9). We also take $h = 1$ and $c_- = 9/4$. In the figure, we show the results for $c_+ = 9/4$, 2, and $5/4$, which correspond to $n_{\rm iso} = 4$, 3, and 2, respectively. One can see that the amplitude of the field φ_+ decreases exponentially, and eventually stays at the constant value, which coincides to Eq. (27). This is very attractive, since there is no blowup of the spectrum at smaller scales, while having extremely blue tilt even as large as $n_{\rm iso} = 4$ at large scales over a few orders of magnitude. Notice that the results remain unchanged even if we vary the Hubble parameter provided that $H \ll F_a$; here, we take a particular value as $H/F_a = 10^{-2}$.

Finally, we comment on the initial amplitude of the fields. The φ_+ should be at large field values in the beginning. One of the simple mechanism to realize this situation is to consider preinflation, where the preinflaton and the Φ_+ have nonrenomalizable coupling in the Kähler potential so as to give rise to a negative Hubble-induced mass term during preinflation. In this case, the initial condition will be $\varphi_+ \simeq M_P$.

V. CONCLUSIONS

We have proposed the concrete model for generating isocurvature fluctuations with extremely blue spectrum for some range of the scale. It is based on the axion model in supersymmetry. The supergravity effects raise the Hubble-induced mass terms in the potential of the $\varphi_\pm$ fields. These Hubble-induced mass terms play two roles. One is that they suppress the fluctuations in the radial directions, $\delta\varphi_\pm$. The other is to make the fields evolve during inflation. In particular, the field value of φ_+ determines the amplitude of the axion isocurvature perturbation: the blue tilt is due to the dynamics of φ_+ moving from the large initial value ($\sim M_P$) down to the PQ symmetry breaking scale F_a during inflation. Depending on the coefficient of the Hubble-induced mass term, c_+, we can obtain $1 < n_{\rm iso} \leq 4$. The prominent feature of this model is that the blue spectrum is realized only while φ_+ is evolving and after it settles down into the potential minimum the spectrum becomes scale invarinat.

The actual scale L where the spectrum is blue is determined by e-folds N' after φ_+ settles down to F_a. For example, $n_{\rm iso} > 1$ at $L \gtrsim 1$ Mpc for $N' \simeq 47$ assuming that the present Hubble radius corresponds to $N \simeq 55$. Observations of large-scale structures, or even PLANCK, could see the existence of the isocurvature fluctuations with a huge blue tilt, which may approve our model in the near future.

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