

No dark matter axion during minimal Higgs inflation

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We study minimal versions of Higgs inflation in the presence of a massless QCD axion. While the inflationary energy scale of the metric variant is too high to accommodate isocurvature bounds, it was argued that Palatini Higgs inflation could evade these constraints. We show, however, that an energy-dependent decay constant enhances isocurvature perturbations, implying that axions can at most constitute a tiny fraction $<10^{-5}$ of dark matter. This conclusion can be avoided in Einstein-Cartan gravity by an additional coupling of the axion to torsion, albeit for a very specific choice of parameters. Analogous constraints as well as the possibility to alleviate them are relevant for all inflationary models with a nonminimal coupling to gravity, including Starobinsky inflation and certain classes of attractor models.

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Inflation and dark matter. In our Universe, about 85% of matter has only been observed through gravitational effects. Understanding the microscopic nature of this dark matter (DM) remains one of the greatest challenges in cosmology. Another fundamental mystery is the origin of the initial conditions for the hot Big Bang. The leading explanation is cosmic inflation [1–4]—an early phase of accelerated expansion—which successfully accounts for the near-homogeneity and isotropy of the Universe and is strongly supported by precision measurements of the cosmic microwave background (CMB) [5,6]. A wide range of DM models have been developed (see Ref. [7]). Similarly, a plethora of inflationary scenarios have been put forward (see Ref. [8]). The key difficulty lies in distinguishing viable proposals from the multitude of possibilities, with the ultimate goal of identifying the mechanisms that govern our Universe. In this work, we reveal an incompatibility between two leading candidates: QCD axions [9–11] and Higgs inflation (HI) [12].

Originally, axions were introduced to address the strong CP problem of QCD, originating from the experimental fact that strong interactions are CP -conserving to a very good accuracy [13]. This observation can be explained by promoting the CP -violating angle of QCD to a dynamical

pseudoscalar field—the axion [9–11], which does not only resolve the strong CP problem but moreover contributes to DM (see, e.g., [14]). The QCD axion can arise, among others, from the original proposal of a spontaneously broken Peccei-Quinn (PQ) symmetry [9–11], from extra dimensions [15] (see the review in [16]), or from local gauge invariance [17–20].

HI agrees excellently with CMB observations and is unique among inflationary proposals since it does not require more particles than have already been observed in experiments. HI is highly sensitive to the different equivalent formulations of general relativity (GR) [21–25] (see the overview in [26]). While the most commonly employed metric version suffers from a low perturbative cutoff scale [27,28], no such problems exist in Palatini HI [29], providing strong motivation for considering the Palatini variant of HI.¹

¹Whether the low cutoff scale of metric HI, above which perturbation theory breaks down, invalidates inflationary dynamics remains debated [27,28,30–34]. It does, however, lead to large and partly uncontrollable corrections to the renormalization group (RG) running, so the high-energy values of the coupling constants cannot be computed uniquely and an uncertainty arises in inflationary predictions [30–32,35–38]. In addition, the reheating temperature exceeds the cutoff scale, so the end of inflation cannot be determined unambiguously [39] and becomes sensitive to the UV completion [40,41]. No such issues arise in Palatini HI because the cutoff scale is significantly higher [29]. As a result, RG running remains under perturbative control and inflationary predictions can be linked to low-energy collider measurements [38,42], while reheating can be calculated unambiguously since the temperature remains below the cutoff [43,44].

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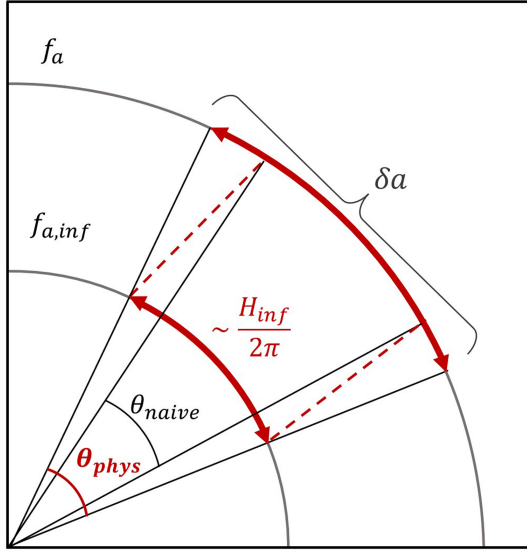


FIG. 1. Representation of varying decay constant. Since the effective decay constant is reduced during inflation, isocurvature perturbations are enhanced. The dashed line corresponds to the result of [47], where the change of the decay constant was not taken into account. The *physical* fluctuation of the axion field is shown with δa . Figure inspired by [48].

A massless axion present during inflation generically sources isocurvature perturbations [45], which are tightly constrained by their nondetection in the CMB. These bounds can be avoided if axions constitute only a negligible fraction of DM, or if the axion field is absent during inflation (or reheating). The latter option can be realized in certain UV completions, most notably those featuring a PQ symmetry, provided its spontaneous breaking occurs after inflation. In this case, however, many models are excluded due to the overproduction of topological defects [46]. In addition, owing to the comparatively low inflationary energy scale in Palatini HI, the viable parameter space for postinflationary symmetry breaking is significantly smaller than in metric HI.

In this paper, we shall investigate the other alternative and check the compatibility of preinflationary axions with HI. While metric HI leads to excessive isocurvature perturbations, it was suggested that Palatini HI could tolerate a massless QCD axion during inflation [47]. However, the energy dependence of the decay constant [48,49] (see also [50–53]) was not taken into account in [47]. Since it is smaller during inflation than in the late Universe, isocurvature perturbations are enhanced (see the illustration in Fig. 1). We point out that this is a generic feature in all inflationary models with a large nonminimal coupling to gravity. Applied to Palatini HI, we show that the presence of a massless axion field that can later act as DM is excluded. Consequently, preinflationary axions constituting DM are incompatible with minimal versions of HI.

Review of isocurvature bound. CMB observations bound the magnitude of uncorrelated isocurvature perturbations Δ_a relative to the adiabatic curvature perturbations $\Delta_{\mathcal{R}}$ as $\Delta_a^2/(\Delta_{\mathcal{R}}^2 + \Delta_a^2) \lesssim 0.038$ [5,14]. Allowing for correlations between isocurvature and adiabatic mode modifies this number slightly, as discussed in Supplemental Material A [54] and Refs. [55,56] therein, but the conclusion of the paper remains unchanged. Plugging in the measured value $\Delta_{\mathcal{R}}^2 \approx 2.1 \times 10^{-9}$, we get [5]

$$\Delta_a \lesssim 9.1 \times 10^{-6}. \quad (1)$$

Now an axion with initial misalignment angle θ_i induces isocurvature perturbations [55,56]

$$\Delta_a = \mathcal{F}_{\text{DM}}^a \frac{\sigma_\theta \sqrt{2(\sigma_\theta^2 + 2\theta_i^2)}}{\theta_i^2 + \sigma_\theta^2}, \quad (2)$$

where $\mathcal{F}_{\text{DM}}^a$ is the relative contribution of axions to DM and σ_θ stands for the typical quantum fluctuation of the angular axion field. Moreover, θ_i is the initial misalignment angle, which can be expressed through $\mathcal{F}_{\text{DM}}^a$ and the axionic decay constant f_a (see Supplemental Material B [54] and Refs. [14,47,57–60] therein),

$$\theta_i = \mathcal{F}_{\text{DM}}^{a1/2} \left(\frac{1.02 \times 10^{12} \text{ GeV}}{f_a} \right)^{7/12}. \quad (3)$$

For $\sigma_\theta \gg \theta_i$, Eq. (2) reduces to $\Delta_a \approx \sqrt{2} \mathcal{F}_{\text{DM}}^a$ and so the isocurvature bound can only be satisfied for a tiny $\mathcal{F}_{\text{DM}}^a$. Having a sizable contribution of axions to DM, $\mathcal{F}_{\text{DM}}^a \approx 1$, is only possible for $\sigma_\theta \ll \theta_i$, in which case Eq. (2) gives $\Delta_a \approx \mathcal{F}_{\text{DM}}^a 2\sigma_\theta/\theta_i$. In summary, the isocurvature constraint (2) implies that one of the following two conditions must be satisfied (see Supplemental Material A [54]):

$$\mathcal{F}_{\text{DM}}^a \lesssim 6.4 \times 10^{-6} \quad \text{OR} \quad \sigma_\theta \lesssim 4.6 \times 10^{-6} \frac{\theta_i}{\mathcal{F}_{\text{DM}}^a}. \quad (4)$$

Since we are interested to have a sizable fraction of DM in axions, we shall focus on fulfilling the second condition.

Argument of [47]. In Palatini HI, the inflationary Hubble scale is (see Refs. [21,38,42] and derivation below)

$$H_I \sim 2.6 \times 10^{-6} \frac{M_P}{\sqrt{\xi}}. \quad (5)$$

Here, ξ sets the strength of nonminimal coupling of the Higgs field to the Ricci scalar. The value of ξ is not known exactly, as will be discussed later. We shall use the largest possible $\xi \sim 10^9$ since this leads to the weakest isocurvature constraint. Then, the Hubble scale (5) yields $H_I \sim 2.2 \times 10^8 \text{ GeV}$, in accordance with [47].

For incorporating the axionic isocurvature bound (4) into Palatini HI, one is tempted to estimate $\sigma_\theta = H_I/(2\pi f_a)$, as in [47], where f_a is the axionic decay constant. With the Hubble scale (5), we would then get (for $\xi \sim 10^9$)

$$f_a \theta_i \gtrsim 9.0 \times 10^{-2} \mathcal{F}_{\text{DM}}^a \frac{M_P}{\sqrt{\xi}} \sim \mathcal{F}_{\text{DM}}^a 6.9 \times 10^{12} \text{ GeV}. \quad (6)$$

If Eq. (6) were to hold, plugging in Eq. (3) would show that having all DM in axions, $\mathcal{F}_{\text{DM}}^a = 1$, can be achieved for $f_a \sim 10^{14} \text{ GeV}$ corresponding to $\theta_i/(2\pi) \sim 10^{-2}$ [47]. We shall demonstrate, however, that this conclusion is premature because it overlooks the need to canonically normalize the axion field in the early Universe.

QCD axion coupled to palatini HI. For a fundamental derivation, the relevant action of the axion a and the Higgs field h (in unitary gauge) coupled to gravity is

$$S = \int d^4x \sqrt{-g} \left[\frac{M_P^2}{2} \Omega^2 R - \frac{1}{2} \partial_\alpha h \partial^\alpha h - \frac{\lambda}{4} h^4 - \frac{1}{2} \partial_\alpha a \partial^\alpha a - \frac{1}{2} \text{Tr} G^{\mu\nu} G_{\mu\nu} + \frac{a}{f_a} c_G \text{Tr} G^{\mu\nu} \tilde{G}_{\mu\nu} \right], \quad (7)$$

where M_P denotes the reduced Planck mass, R the Ricci scalar, λ the Higgs self-coupling, and we defined

$$\Omega^2 = 1 + \frac{\xi h^2}{M_P^2}, \quad (8)$$

with nonminimal coupling constant ξ . Moreover, $G_{\mu\nu}$ and $\tilde{G}_{\mu\nu}$ correspond to the field strength tensor of QCD and its dual, respectively, and we have the dimensionless parameter $c_G \sim \alpha$, where α is the gauge coupling. We view Eq. (7) as the low-energy effective field theory of a pseudoscalar a coupled to QCD via an operator of mass dimension 5 that is suppressed by a fixed mass scale f_a . Crucially, the shift symmetry of the axion prevents the appearance of a nonminimal interaction of a with R .

As usual, we remove the nonminimal coupling to curvature with the conformal transformation $g_{\mu\nu} \rightarrow \Omega^{-2} g_{\mu\nu}$. In contrast to the situation in metric gravity, the Ricci tensor $R_{\mu\nu}$ of Palatini GR is independent of the metric $g_{\mu\nu}$ and, therefore, curvature simply transforms as $R \rightarrow \Omega^2 R$. Thus, action (7) becomes

$$S = \int d^4x \sqrt{-g} \left[\frac{M_P^2}{2} \dot{R} - \frac{1}{2\Omega^2} \partial_\alpha h \partial^\alpha h - \frac{\lambda}{4\Omega^4} h^4 - \frac{1}{2\Omega^2} \partial_\alpha a \partial^\alpha a - \frac{1}{2} \text{Tr} G^{\mu\nu} G_{\mu\nu} + \frac{a}{f_a} c_G \text{Tr} G^{\mu\nu} \tilde{G}_{\mu\nu} \right], \quad (9)$$

where it is important to note that the gauge kinetic term is invariant under the conformal transformation so that the gauge field remains canonical (see Ref. [61]). Since in the

form (9) the coupling of gravity to matter is minimal, the Palatini and metric formulations of GR are equivalent. This has allowed us to replace curvature by its Riemannian counterpart $\dot{R}$, which is a function of the metric only.

Still following the standard analysis of Palatini HI, we next perform a field transformation for the Higgs field, introducing a new field χ defined by [21,38]

$$h = \frac{M_P}{\sqrt{\xi}} \sinh\left(\frac{\sqrt{\xi}\chi}{M_P}\right), \quad (10)$$

so that the action (9) becomes

$$S = \int d^4x \sqrt{-g} \left[\frac{M_P^2}{2} \dot{R} - \frac{1}{2} \partial_\alpha \chi \partial^\alpha \chi - U - \frac{\partial_\alpha a \partial^\alpha a}{2 \cosh^2\left(\frac{\sqrt{\xi}\chi}{M_P}\right)} - \frac{1}{2} \text{Tr} G^{\mu\nu} G_{\mu\nu} + \frac{a}{f_a} c_G \text{Tr} G^{\mu\nu} \tilde{G}_{\mu\nu} \right], \quad (11)$$

with inflationary potential

$$U = \frac{\lambda M_P^4}{4\xi^2} \tanh^4\left(\frac{\sqrt{\xi}\chi}{M_P}\right). \quad (12)$$

Evaluating the first slow-roll parameter, we can express χ as a function of the number $N_\star$ of e-foldings (see Ref. [42]),

$$\chi(N) \approx \frac{M_P \text{arccosh}(16\xi N_\star)}{2\sqrt{\xi}}, \quad (13)$$

and then match the amplitude of CMB perturbations to obtain the constraint $\xi = 1.2 \times 10^{10} \lambda$, where we used $N_\star \approx 51$, as in [38]. Then, the potential (12) yields the Hubble scale (5). Due to uncertainties in the measurement of the top Yukawa coupling, the inflationary value of λ is unknown [38,62]. Although RG analysis indicates a small $\lambda \sim 10^{-3}$ at high energies [38], we shall follow [47] and use $\lambda = 0.1$ corresponding to $\xi \sim 10^9$ since a large ξ weakens isocurvature perturbations.

Finally, we define an approximately canonical axion field

$$A = \frac{a}{\Omega} = \frac{a}{\cosh\left(\frac{\sqrt{\xi}\chi}{M_P}\right)}, \quad (14)$$

where we plugged the field (10) into Eq. (8). We arrive at

$$S = \int d^4x \sqrt{-g} \left[\frac{M_P^2}{2} \dot{R} - \frac{1}{2} \partial_\alpha \chi \partial^\alpha \chi - U - \frac{1}{2} \partial_\alpha A \partial^\alpha A - \frac{1}{2} \text{Tr} G^{\mu\nu} G_{\mu\nu} + \frac{\Omega A}{f_a} c_G \text{Tr} G^{\mu\nu} \tilde{G}_{\mu\nu} + \Delta\mathcal{L} \right], \quad (15)$$

where $\Delta\mathcal{L}$ contains correction terms,

$$\Delta\mathcal{L} = -\frac{\sqrt{\xi}A}{M_P} \tanh\left(\frac{\sqrt{\xi}\chi}{M_P}\right) \left[\partial_\alpha \chi \partial^\alpha A + \frac{\sqrt{\xi}A}{2M_P} \tanh\left(\frac{\sqrt{\xi}\chi}{M_P}\right) \partial_\alpha \chi \partial^\alpha \chi \right]. \quad (16)$$

Since during inflation $A \sim H_I \sim \sqrt{\lambda} M_P / \xi$, the correction terms in $\Delta\mathcal{L}$ are at least suppressed as $\sqrt{\xi}A/M_P \lesssim \sqrt{\lambda/\xi}$, and we shall neglect them in the following (see also [48,49,63,64], as well as Supplemental Material C [54] and Ref. [25] therein for a consistency check).

Field-dependent decay constant. Crucially, Eqs. (14) and (15) make evident that the effective decay constant $f_{a,\text{inf}}$ during inflation differs from its low-energy value f_a ,

$$f_{a,\text{inf}} = \frac{f_a}{\Omega} \approx \frac{f_a}{\sqrt{8\xi N_\star}}, \quad (17)$$

where we used the solution (13) in Eq. (8). Therefore, plugging $\sigma_\theta = H_I/(2\pi f_{a,\text{inf}})$ together with Eq. (5) into the second condition of the isocurvature bound (4) implies

$$f_a \theta_i \gtrsim 0.25 \sqrt{N_\star} \mathcal{F}_{\text{DM}}^a M_P. \quad (18)$$

This constraint is much stronger than the previously proposed Eq. (6). Interestingly, it is not sensitive to the uncertainty in λ (and equivalently ξ). Upon inserting θ_i , as in Eq. (3), Eq. (18) would bound the abundance of axions as

$$\mathcal{F}_{\text{DM}}^a \lesssim 1.1 \times 10^{-8} \left(\frac{f_a}{M_P} \right)^{5/6}. \quad (19)$$

Thus, the first condition of the isocurvature bound (4) is relevant for any sub-Planckian decay constant,

$$\mathcal{F}_{\text{DM}}^a \lesssim 6.4 \times 10^{-6}. \quad (20)$$

Therefore, QCD axions that are present as a massless field during Palatini HI can at most contribute a tiny fraction $\sim 6 \times 10^{-6}$ to DM. We show in Supplemental Material D [54] and Refs. [9–11,15–20,48,49,63,64] therein that the same conclusion can be obtained in a UV completion by a PQ symmetry, where a arises as phase of a complex PQ field.

Axion and generic inflationary models. Clearly, the existence of these bounds does not depend on the particular structure of Palatini HI. Only important is a nonminimal coupling $\Omega^2 \gg 1$ so that the axion kinetic term is multiplied by $1/\Omega^2 \ll 1$, as in the action (9). This directly translates to a decay constant that is smaller during inflation than now, according to the first equality in Eq. (17). Importantly, the same conclusion also holds in the metric formulation of GR. In this case, the conformal

transformation $g_{\mu\nu} \rightarrow \Omega^{-2} g_{\mu\nu}$ yields an additional contribution proportional to $(\partial_\alpha \Omega)^2$ (see Ref. [26]), but this only contributes to the kinetic term of the inflaton. Therefore, the enhancement of isocurvature perturbations is not limited to HI but generic in all inflationary models with a large nonminimal coupling to gravity.²

Consequently, the strengthening of isocurvature bounds that we have identified applies to leading inflationary plateau models, in particular Starobinsky inflation [1] (for details see Supplemental Material E [54] and Refs. [26,65,66] therein), metric HI [12], and certain classes of attractor models [67,68]. Even without the effect of the conformal transformation, these scenarios are typically incompatible with isocurvature bounds because of their high inflationary Hubble scale. Crucially, any attempt to alleviate this must additionally account for the further suppression of the inflationary decay constant induced by the nonminimal coupling.³ We note that α -attractor models represent an exception. While they can be constructed without invoking a nonminimal coupling [70,71], some can also be obtained from a negative nonminimal coupling, $\xi < 0$ [70,72]. In this case, $\Omega < 1$, which could open up the possibility of alleviating isocurvature bounds.

Conversely, the nonminimal interaction can lead to new channels for detecting axions in those models that are not ruled out: we expect that the rapid change of Ω during reheating leads to the production of axions (cf. [39]), which may result in an observable signal in the effective number of massless degrees of freedom [73] (see also [14]).

Ways out. Several options exist to alleviate isocurvature bounds. Since Palatini gravity is part of the Einstein-Cartan formulation (see Ref. [26] for terminology), we can add many more terms composed of torsion to the action (cf. [26,61,74]). Some of them can alleviate isocurvature bounds, as discussed in Supplemental Material F [54] and references therein [48,49,63,64] (see also [75]). Moreover, it is possible to consider nonminimal models of HI, such as [24,25], or to include a direct coupling between the norm of the PQ field and torsion, $|\Phi_{\text{PQ}}|^2 T_\alpha T^\alpha$.

²In a UV completion by a PQ field, inflationary models driven by the radial mode $|\Phi_{\text{PQ}}|$ represent an exception [48,49,63,64]. Since in this case $|\Phi_{\text{PQ}}|$ is displaced from its minimum, the effective inflationary decay constant can be larger than its low-energy counterpart in spite of the presence of a nonminimal coupling to gravity.

³In the recent paper [69], which appeared after our work, a mechanism was constructed for alleviating isocurvature constraints based on a nonminimal coupling ξ_{PQ} of the radial PQ field to curvature. That work discusses isocurvature constraints in Starobinsky inflation, but without accounting for the modification of the axion kinetic term induced by the conformal transformation. In [66], we study the combined effect of both nonminimal couplings ξ and ξ_{PQ} , showing that the inflaton coupling ξ significantly reduces the viable interval of ξ_{PQ} in which isocurvature bounds can be alleviated.

Independently of torsion, one can consider an early phase of confinement [76], which can arise from an effective coupling of the Higgs field to the gauge kinetic term $\text{Tr} G^{\mu\nu} G_{\mu\nu}$ and would drive the QCD axion to the CP -conserving value already during inflation, thereby reducing isocurvature perturbations. We do not expect this mechanism to be generic in our approach since there is no simple coupling of torsion to gauge fields and so new interactions involving $G^{\mu\nu}$ arise [61]. In a UV completion by a PQ field, one can also directly include a coupling $|\Phi_{\text{PQ}}|^2 h^2$ to change the inflationary vacuum expectation value of the PQ field (see also [77]),⁴ or even allow for an explicit breaking of PQ symmetry [51, 78–81]. However, all of these options require additional nonminimal interactions with coupling constants in specific intervals—cf. the highly tuned parameter choice in Supplemental Material F [54]—and so arguably should be considered as nonminimal. Finally, one may turn to more generic axionlike particles (see Ref. [82]) as a solution to the strong CP problem, as, e.g., in [83].

Conclusion. The microscopic origins of inflation and dark matter remain among the most important open questions in cosmology. Given that many models are still consistent with all observations, additional constraints are needed to distinguish among the plethora of proposals. In this work, we have corrected a previous claim in the literature and demonstrated that a massless QCD axion—if abundant enough to account for dark matter in the late Universe—must not exist during simple Higgs inflation models. This finding resonates with the approach of Higgs inflation,

which is motivated by the lack of detected particles beyond the standard model.

Our result stems from the energy dependence of the axionic decay constant. An inflationary value that is smaller than its late-time counterpart enhances isocurvature perturbations. As a result, a significant axion abundance is incompatible with observational isocurvature bounds. Our findings have important implications for all inflationary models with nonminimal couplings to gravity. First, the enhancement of isocurvature perturbations is a general feature of such scenarios. Second, coupling the axion to torsion can relax isocurvature constraints, in particular, Einstein-Cartan gravity can rescue proposals otherwise deemed excluded due to excessive isocurvature perturbations.

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