PHYSICAL REVIEW D 96, 083529 (2017)

Extranatural inflation redux

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(Received 24 August 2017; published 30 October 2017)

The success of a given inflationary model crucially depends upon two features: its predictions for observables such as those of the cosmic microwave background (CMB) and its insensitivity to the unknown UV physics such as quantum gravitational effects. Extranatural inflation is a well-motivated scenario that is insensitive to UV physics by construction. In this five-dimensional model, the fifth dimension is compactified on a circle, and the zero mode of the fifth component of a bulk $U(1)$ gauge field acts as the inflaton. In this work, we study simple variations of the minimal extranatural inflation model in order to improve its CMB predictions while retaining its numerous merits. We find that it is possible to obtain CMB predictions identical to those of, e.g., the $\mathcal{R}+\mathcal{R}^2$ Starobinsky model of inflation and show that this can be done in the most minimal way by having two additional light fermionic species in the bulk, with the same $U(1)$ charges. We then find the constraints that CMB observations impose on the parameters of the model.

DOI: 10.1103/PhysRevD.96.083529

I. INTRODUCTION

In the Standard Model of cosmology, one posits, among other things, that the early Universe was incredibly homogeneous, isotropic, and spatially flat but still had small, adiabatic, and Gaussian density perturbations that had a nearly scale-invariant power spectrum. In the recent past, a successful combination of theory and experiments has validated this picture. Cosmic inflation [1–9] is a mechanism by which the microscopic theories (very similar to the ones we routinely use to explain the behavior of elementary particles in, e.g., colliders) can give rise to a Universe that would look very much like the one posited in the Standard Model of cosmology. One of the questions worth addressing in near future is whether one can learn more about the details of inflation.

In the next decade, the upcoming cosmological experiments [10–14] will help us better understand this era in the early history of the Universe. Although better observational data will surely help in this quest, a careful look at the literature suggests that realistic inflation model building is a formidable task irrespective of the available data [15]. For example, small field inflationary potentials suffer from the overshooting problem, while large field inflation can be extremely sensitive to the unknown UV physics. Thus, even though there are many models of inflation that seem to give cosmic microwave background (CMB) predictions that agree with recent observations (such as those of Planck experiment [16]), a model of inflation that does not suffer from other theoretical problems and still agrees with experiments is a rarity. In this paper, we wish to look for a model of inflation that achieves this.

A well-known example of a model in which there are supposedly no issues of UV sensitivity, at least at the level of field theory, is extranatural inflation [17].[1] As we explain in Sec. II, there are many reasons why this scenario is preferred over natural inflation [21,22]. In extranatural inflation, the zero mode of the fifth component of a bulk Abelian gauge field acts as the inflaton. This field is clearly a scalar under four-dimensional coordinate transformations. The potential of the inflaton is generated by loop corrections of light fermions (charged under the bulk gauge group) present in the bulk. However, it turns out that this minimal scenario gives CMB predictions identical to those of natural inflation, which is increasingly getting disfavored with newer CMB data. The most recent CMB data, however, are completely consistent with the famous $\mathcal{R}+\mathcal{R}^2$ model of inflation of Starobinsky [1]. It may thus be worth it to look for variations of the minimal scenario of extranatural inflation that give predictions identical to those of the Starobinsky model. We propose achieving this by adding extra fermionic species in the bulk. The observational data then constrain the charge of these fermions under the bulk gauge group.

It turns out that the effect of adding such extra fermions in the bulk is the addition of more sinusoidal functions to the potential of natural inflation. Although scenarios with such potentials have been studied for a long time in the context of both natural inflation [23–30] and extranatural

[1] It is well known that extranatural inflation requires a rather small ($\mathcal{O}(10^{-3})$) value of a four-dimensional gauge coupling, and this is challenging to achieve in known UV completions such as string theory [18–20]. We do not address UV completion of extranatural inflation in this work.

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inflation [31–35], the possibility of obtaining CMB predictions similar to those of the Starobinsky model has not been explored (see, however, Refs. [36,37]).

This paper is organized as follows. In Sec. II, we begin by explaining the issue of UV sensitivity of inflation and then review the relevant details of extranatural inflation. Then, in Sec. III, we explain our scenario and its possible CMB predictions, show that adding a single extra fermion in the bulk does not sufficiently improve the CMB predictions of extranatural inflation, and then present the minimal variation of extranatural inflation of which the CMB predictions are identical to those of the Starobinsky model. We then find the values of all the parameters and scales in the model using the constraints obtained from CMB observations. Finally, we conclude in Sec. IV.

We work with $\hbar = c = 1$ units; moreover, M_p is the four-dimensional (4D) Planck mass, $M_{\rm Pl}$ is the 4D reduced Planck mass, $M_p^{(5)}$ is five-dimensional Planck mass, ℓ_P is the 4D Planck length, R is the radius of the extra dimension, $L = 2\pi R$, and $\mathcal{R}$ is the Ricci scalar.

II. EXTRANATURAL INFLATION

Before getting into the details of extranatural inflation, we revisit the issue of UV sensitivity of inflation, particularly large field inflation [15].

A. UV sensitivity of inflation

Even though during inflation the energy density of the inflaton field dominates the Universe, the inflaton would be but one field in a Lagrangian that would, in any realistic picture, contain many fields. In fact, there would at least be Standard Model fields in the same Lagrangian. Moreover, we might need more fields to explain, e.g., neutrino masses, dark matter, or baryon asymmetry of the Universe or to solve the cosmological constant problem. Moreover, there must be new degrees of freedom that would show up near the Planck scale, which would help unitarize the graviton-graviton scattering at Planck scale.

Every theory is to be interpreted as a Wilsonian eective field theory (EFT) with a physical cutoff. The Wilsonian effective action can be obtained from the UV theory by integrating out the physics above a UV cutoff Λ_0. For example, if, in the path integral of the theory, one integrates out all the fields except the inflaton and also integrates out all the high-frequency modes (above some scale Λ_0) of the inflaton, one would obtain the Wilsonian effective Lagrangian of the inflaton, which would be of the form [38]

$$\mathcal{L}_{\rm eff}[\phi] = \mathcal{L}_{\ell}[\phi] + \sum_{i=1}^{\infty} c_i \frac{\phi^{4+2i}}{\Lambda_0^{2i}} + d_i \frac{(\partial\phi)^2\phi^{2i}}{\Lambda_0^{2i}} + e_i \frac{(\partial\phi)^{2(i+1)}}{\Lambda_0^{4i}} + \cdots, \quad (1)$$

where it is assumed that a Z_2 symmetry holds well in the UV theory.

For most observables, when one performs experiments at energies well below Λ_0, the higher-dimension operators have negligible effect. But this is not always true; e.g., the mass of an elementary scalar is highly sensitive to *all* the higher-dimension operators (see Ref. [38] for details). For example, the $m^2\phi^2$ inflation happens to be such that all the Wilson coefficients in the above Lagrangian except the m^2 term vanish. The question is what ensures that this will happen. In the context of inflation, this problem is closely related to the so-called eta problem: given that the higher-dimension operators could renormalize m^2 (with typical contributions of the order of Λ_0^2), why is $m^2 \ll H^2$?

1. Rolling beyond the cutoff

The Wilson effective action is valid only when one performs experiments at energies below the cutoff scale. If the Lagrangian contains a higher-dimension operator, at high enough energies, unitarity is violated. On the other hand, during large field inflation, the field rolls by a super-Planckian amount. Notice that if we have a large field inflation $\phi > M_p$, which means $\phi > \Lambda_0$, this means that the contribution of higher powers in Eq. (1) is even higher (unless the coefficients somehow compensate for this). This raises the question of whether, when the field rolls by an amount greater than the cutoff, the Wilson action is even valid. It is often argued that, since energy density during inflation is conveniently sub-Planckian, there is no problem if the field vacuum expectation value (vev) changes by super-Planckian values; but in the Wilson action, since the field excursion is super-Planckian, unless all the infinite Wilson coefficients are guaranteed to be small, we would surely have a problem.

Finally, one may be concerned with how the inflaton potential may get affected by unknown UV physics, e.g., loop corrections due to heavy particles, the effects of virtual black holes, or gravitational and other instantons in, e.g., string theory [39,40].

B. Symmetries of UV theory: Global and gauge

The low energy EFT of inflaton thus faces a number of problems. However, these problems can be cured if the UV completion of this EFT has certain symmetries.

For example, if one assumes that there is a global shift symmetry in the UV theory, this will set all the c_i and d_i to zero. But this will also set even m and λ to be zero. One could then generate m by breaking the global shift symmetry softly by an independent sector.

As far as the coefficients e_i are concerned, they need not be small since for a homogeneous inflating background $(\partial\phi)^2 = \dot{\phi}^2 + (\nabla\phi)^2 \approx \dot{\phi}^2 = 2\epsilon H^2$. which is suppressed by ϵ, the Hubble slow-roll parameter. Moreover, the quantum correction to the mass of the scalar due to these derivative operators is $\delta m \propto m$, while their contributions to a scattering amplitude at energies lower than Λ_0 is negligible anyway.

The most familiar example implementing these ideas is natural inflation [21,22]. To begin with, the inflaton is assumed to be the Goldstone mode of a spontaneously broken global $U(1)$ symmetry. This causes its potential to vanish at all orders in perturbation theory. If one now also assumes that this symmetry is anomalous, i.e., though it exists in the classical theory, it is broken by quantum effects, and then gauge instantons generate a potential that is a cosine at the leading order. However, the requirement of having large field slow-roll inflation causes the scale of spontaneous breaking of $U(1)$ to be super-Planckian. Since there are reasons to suspect that there can be no continuous global symmetries in quantum gravity (see Ref. [41] and the discussion in Sec. 4 of Ref. [42]), one must find out alternatives to the most basic natural inflation (e.g., by having multiple $U(1)$s [30] or by taking into account spinodal instabilities [43]). In stark contrast, in extranatural inflation [17], instead of global symmetries, a gauge symmetry forbids the coefficients c_i and d_i so that the unknown UV physics has negligible effects on the potential of the inflaton. A lot of recent work [18–20] has been devoted to trying to understand the issue of possible UV insensitivity of extranatural inflation and similar models. In this work, however, we assume that the inflaton potential for extranatural inflation can be protected from unknown UV effects and focus on improving its CMB predictions.

C. Extranatural inflation in a nutshell

In the rest of the paper, we will restrict our attention to quantum field theory in a five-dimensional (5D) spacetime in which the fifth dimension is compactified on a circle; i.e., the spacetime in the absence of gravity is $M_4 \times S^1$. The coordinates on this 5D spacetime are denoted as (x^μ, y).

1. Basic set-up

Since the extradimensional coordinate y is identified to $y + 2\pi R$, for all fields, $\Phi(x^\mu, y) \sim \Phi(x^\mu, y + 2\pi R)$.

By mode expansion, one can verify that a single 5D (i.e., bulk) Abelian gauge field A_M is equivalent to the following fields (it is easiest to see the field content in the so-called "almost-axial" gauge; see, e.g., Ref. [44] for details):

(i) $A_5^{(0)}$, which is a gauge-invariant, massless 4D scalar with no tree-level potential (this will act as the inflaton);

(ii) $A_\mu^{(0)}$, which has a residual gauge invariance, a massless 4D vector;

(iii) $A_\mu^{(n)}$, an infinite tower of massive 4D vectors [the Kaluza-Klein (KK) modes of the vector].

Before proceeding, it is worth noting that in five dimensions a gauge field has mass dimension $3/2$, while the corresponding 5D gauge coupling g_5 has mass dimension $-1/2$. Of course, the 4D gauge coupling is dimensionless; in fact,

$$g_4 = \frac{g_5}{\sqrt{2\pi R}}. \tag{2}$$

One can define the dimensionless and gauge-invariant field

$$\theta(x) = g_5 \oint dy A_5(x^\mu, y), \tag{3}$$

which is also the gauge-invariant Wilson loop of A_5 along the extra dimension (and, as we shall see, is going to be very simply related to the inflaton field). It is easy to verify that this integral will pick only the contributions from the zero mode of the fifth component of the bulk Abelian gauge field, i.e., $A_5^{(0)}$.

If bulk matter is present, a potential for $A_5^{(0)}$ and hence the inflaton is readily generated. If there is one bulk matter field that is charged under the gauge symmetry of the 5D Abelian gauge field (e.g., a bulk complex scalar field or a bulk spinor field), and hence has a charge Q, then the gauge-covariant derivative in its Lagrangian will be given by $D_M = \partial_M - iQg_5A_M$. From the 4D point of view, this bulk matter field will give rise to an infinite tower of KK modes. Thus, from the 4D point of view, the 4D-scalar-field $A_5^{(0)}$ has coupling to all these infinite matter fields. So, every KK mode of the matter field will generate a Coleman-Weinberg potential for $A_5^{(0)}$. For a bulk matter field with mass m_a and $U(1)$ charge Q_a, the Coleman-Weinberg potential of θ due to the KK modes of this bulk matter field is given by (see Refs. [45–47] for some early references, Ref. [48] for a particularly accessible derivation, and Refs. [20,31] for some relatively recent papers)

$$V(\theta) = \pm \frac{3}{64\pi^6 R^4} \left[\sum_{n=1}^{\infty} c_n e^{-2\pi R n m_a} \mathrm{Re}(e^{inQ_a\theta}) \right], \tag{4}$$

where

$$c_n = \frac{1}{n^5} + \frac{2\pi R m_a}{n^4} + \frac{(2\pi R m_a)^2}{3n^3} \tag{5}$$

and the $+$ sign is for fermionic matter, while the $-$ sign is for bosonic matter.

If the bulk matter is massless (or has a mass very small as compared to R^{-1}), taking the $m_a \to 0$ limit in the above expression gives

$$V(\theta) = \pm \frac{3}{64\pi^6 R^4} \sum_{n=1}^{\infty} \frac{\cos(nQ_a\theta)}{n^5}. \tag{6}$$

Finally, for the sake of completion, when one turns on gravity, the spacetime will have a curved geometry but will still retain the topology of $M_4 \times S^1$. The radius of the circle will then be different at different points of 4D spacetime and will be determined from the vev of a scalar field called the radius modulus (or radion). The 5D Einstein gravity gives rise to the following:

(i) $h^{(0)}_{55}$, which is gauge invariant, is a massless 4D scalar with no tree-level potential (this is the radius modulus or the radion, and its vev will have to be stabilized to a value large enough so that the inflaton potential can be kept protected from unknown UV effects).
(ii) $h^{(0)}_{5\mu}$, which has a residual gauge invariance, is a massless 4D vector (the graviphoton).
(iii) $h^{(0)}_{\mu\nu}$, which has a residual gauge invariance, is a massless spin-2 particle (the familiar 4D graviton).
(iv) $h^{(n)}_{\mu\nu}$ is an infinite tower of massive KK gravitons.

2. *The various fields in extranatural inflation*

We thus have a 5D Abelian gauge field, 5D Einstein gravity, and bulk matter (and, if required, bulk cosmological constant and brane tension). We would be interested in solutions in which the radion is stabilized (i.e., it sits at the bottom of its potential) so that the physical size of the extra dimension is fixed. On the other hand, the inflaton is rolling down, and hence the effective 4D cosmological constant is positive and dominates the dynamics of the Universe; thus, the 4D universe is undergoing inflation. In this work, the vaccum energy at the minimum of the inflaton potential is assumed to be zero.

D. Connection to natural inflation

If we specialize to the case of the potential generated due to just light fermions in the bulk and notice that the subsequent terms in Eq (6) are suppressed so that the term with $n = 1$ dominates, the potential due to only one fermion in the bulk will be of the form

$$V(\theta) \approx \frac{3}{64\pi^6 R^4}\cos(Q_a\theta). \tag{7}$$

The dimensionless field $\theta(x)$ is canonically normalized to [17]

$$\phi = \frac{\theta}{g_4(2\pi R)}, \tag{8}$$

which is the inflaton; this gives

$$V(\phi) \approx \frac{3}{64\pi^6 R^4}\cos(Q_a g_4 2\pi R\phi). \tag{9}$$

If we now define

$$f = \frac{1}{2\pi R g_4}, \tag{10}$$

then

$$V(\phi) \approx \frac{3}{64\pi^6 R^4}\cos\left(\frac{Q_a\phi}{f}\right). \tag{11}$$

If one adds an appropriate constant to this potential in order to keep the minimum of the potential at zero vacuum energy,[2] one obtains the potential of natural inflation

$$V = \Lambda^4\left[1 + \cos\left(\frac{\phi}{f_{\rm eff}}\right)\right], \tag{12}$$

where the overall factor Λ in Eq. (12) is given by

$$\Lambda^4 = \frac{3}{64\pi^6 R^4}. \tag{13}$$

It is clear from Eq. (12) that CMB data are sensitive to the "decay constant" $f_{\rm eff} = f/Q_a$. Thus, it is only the ratio of f [defined by Eq. (10)] and Q_a, which can be determined from the data.

In summary, if we have just one light fermion in the bulk with $U(1)$ charge unity, then the predictions of extranatural inflation are identical to those of natural inflation to a very good accuracy. It is, however, noteworthy that the most recent CMB data disfavor natural inflation at 2σ statistical significance [16].

E. Merits of extranatural inflation

Before proceeding, it is worth noting that in extranatural inflation, even though the effective scale f appears to be super-Planckian in the 4D description, there is no super-Planckian mass scale involved in the 5D description. This is because a super-Planckian "axion decay constant" can be obtained by having a small 4D gauge coupling,

$$\frac{f}{M_{\rm Pl}} = \frac{1}{2\pi g_4(RM_{\rm Pl})}. \tag{14}$$

Moreover, heavy particles that are uncharged under the bulk $U(1)$ gauge symmetry cannot affect the inflaton potential, while, although the potential gets affected by the loops of heavy particles that are charged under the bulk gauge symmetry, this effect is exponentially suppressed [see Eq. (4)].

The only remaining concern is the super-Planckian excursion of the inflaton since there is the possibility that quantum gravitational effects could still affect the potential. In Ref. [17], it is mentioned that, since the super-Planckian decay constant originates from sub-Planckian mass scales, the authors expect that quantum gravitational effects on the potential go at most as $\sim e^{-2\pi RM_5}$ (the exponential is

[2]This is equivalent to assuming a solution to the cosmological constant problem [49].

suppressed by the Euclidean action of a relativistic particle going around the extra dimension). A look at the recent literature lends credence to the notion that the jury is still out on the validity of such estimates [39,40]. We also note that successful large field inflation in extranatural inflation is achieved by assuming the 4D gauge coupling to be too small. This may be harmless in field theory but, as is well known [18], strongly resists any embedding in string theory. In this work, we shall not be exploring these fascinating issues any further but instead turn to observational constraints.

III. CMB OBSERVATIONS AND PARAMETERS

We saw in the last section that one light fermion in the bulk leads to a potential that, to leading order, is of the form of a cosine. To aid the discussion, we will use the phrase "first fermion" to refer to the bulk fermion of which the loop corrections generate the potential of natural inflation. We now turn our attention to variations of extra natural inflation in which additional fermions shall be present in the bulk [31,50]. It is worth noting that we only consider the additional fermions to be light as compared to the KK scale.

Let us suppose we have one light fermion with charge Q_a and then $\mathcal{N}$ copies of another light fermion with charge Q; then, the potential of the inflaton would be

$$V(\phi) = \frac{3}{64\pi^6 R^4}\left\{\sum_{n=1}^{\infty}\frac{1}{n^5}\cos\left(\frac{nQ_a\phi}{f}\right) + \mathcal{N}\sum_{n=1}^{\infty}\frac{1}{n^5}\cos\left(\frac{nQ\phi}{f}\right) + C\right\}. \quad (15)$$

The constant C in the above potential is chosen such that the vacuum energy of the minimum of the potential is zero. Since it is only the ratios f/Q_a and f/Q that determine the arguments of the cosines, we could set Q_a to 1 and hence rescale Q and f accordingly. Thus, if we have one light fermion with charge $+1$ and then $\mathcal{N}$ copies of another light fermion with charge Q, then the potential of the inflaton would be

$$V(\phi) = \frac{3}{64\pi^6 R^4}\left\{\sum_{n=1}^{\infty}\frac{1}{n^5}\cos\left(\frac{n\phi}{f}\right) + \mathcal{N}\sum_{n=1}^{\infty}\frac{1}{n^5}\cos\left(\frac{nQ\phi}{f}\right) + C\right\}. \quad (16)$$

Notice that this is quite different from the potentials dealt with in, e.g., multinatural inflation, in which the amplitudes, frequencies, and phases of the two cosines could all be arbitrarily different from each other. Thus, one comes across a more constrained scenario simply due to the extradimensional embedding of our model. In the rest of this section, we show that with this simple choice of particle content there exist parameter choices that will lead to CMB predictions identical to those of the $\mathcal{R}+\mathcal{R}^2$ model of inflation of Starobinsky [1].

A. Slow-roll inflation

Irrespective of how complicated the inflaton potential is, if the potential slow-roll parameters, defined by

$$\epsilon_V \equiv \frac{M_{\rm pl}^2}{2}\left(\frac{V'}{V}\right)^2, \quad (17)$$

$$\eta_V \equiv \frac{M_{\rm pl}^2 V''}{V}. \quad (18)$$

are small as compared to unity at the time when the pivot scale k_* crossed the Hubble radius during inflation, the primordial scalar and tensor power spectra are given by power functions of the wave number. Thus, in slow-roll inflation, the primordial scalar and tensor power spectra are given by

$$P_s(k) = A_s\left(\frac{k}{k_*}\right)^{n_s-1}, \quad (19)$$

$$P_t(k) = A_t\left(\frac{k}{k_*}\right)^{n_t}, \quad (20)$$

where the amplitude of the scalar power spectrum A_s, the tensor-to-scalar ratio, and the scalar spectral index n_s are given, respectively, by

$$A_s \approx \frac{V}{24\pi^2 M_{\rm pl}^4 \epsilon_V}, \quad (21)$$

$$n_s \approx 1 + 2\eta_V - 6\epsilon_V, \quad (22)$$

$$r \approx 16\epsilon_V. \quad (23)$$

Given these, other quantities such as the amplitude of the tensor power spectrum A_t and the tensor spectral index n_t,

$$A_t \approx \frac{2V}{3\pi^2 M_{\rm pl}^4}, \quad (24)$$

$$n_t \approx -2\epsilon_V. \quad (25)$$

can easily be found from the relations $A_t = rA_s$ and $n_t = -(\frac{r}{8})$.

Recall that if the pivot scale k_* goes out of the Hubble radius during inflation at an epoch that was N_* e-foldings from the end of inflation, then we expect, for grand unified theory scale inflation, N_* to be between 50 and 60, and in the following, we shall set N_* to 60. For the potential given in Eq (6), for any choice of the parameters $\mathcal{N}$, R, f, and Q, one can numerically find $\phi_{\rm end}$, the value of the inflaton field when inflation ends, and then use[3]

$$N(\phi) = \int_{\phi_{\rm end}}^{\phi} \frac{d\phi}{M_{\rm Pl}\sqrt{2\epsilon_V(\phi)}} \tag{26}$$

and find ϕ_*, the value of the inflaton field when the pivot scale exited the Hubble radius. Finally, one can find the corresponding value of slow-roll parameters corresponding to ϕ_* and hence the scalar and tensor power spectra.

B. Numerical results

Lets us now see how the CMB predictions (i.e., A_s, n_s, and r) in this scenario change as we explore the parameter space of $\mathcal{N}$, R, f, and Q. We would like to restrict our attention to the region of parameter space that offers slow-roll inflation and that yields n_s and r that are most compatible with the Planck measurements, i.e., 2015 Planck TT,TE,EE+lowP data [16]. These data imply that $n_s = 0.9652 \pm 0.0047$ at 1σ C.L., while $r < 0.099$ (for $k_* = 0.002\ \text{Mpc}^{-1}$) and $A_s = 2.2065^{+0.0763}_{-0.0738} \times 10^{-9}$.

Notice that in slow-roll inflation n_s and r are completely determined by the slow-roll parameters and hence do not depend on any overall multiplicative factor in the potential. The value of any overall factor in the potential, e.g., R in Eq. (16), can be adjusted to ensure that A_s matches the observed value and is thus determined by A_s and not n_s and r.

When $\mathcal{N} = 0$, one recovers the minimal version of extranatural inflation. Its predictions for the spectral index and the tensor-to-scalar ratio are identical to those of natural inflation, and this model is mildly ($> 2\sigma$) disfavored by the Planck data [16]. If we restrict our attention to the case $Q = 1$, then, no matter what value of $\mathcal{N}$ one works with, it is only R that will be redefined. This will not change the slow-roll parameters and hence will not change the spectral index and the tensor-to-scalar ratio. Thus, the CMB predictions in this case will not be any better than those of natural inflation. Similarly, the case $Q = 0$ will only redefine C. Moreover, since cosine is an even function, for any given $\mathcal{N}$, the sign of Q is unimportant. By numerically solving the underlying equations, one can also find the following: i. for $f \leq 1.5$, the assumption $\eta_V \ll 1$ no longer remains valid, and therefore one of the slow-roll conditions gets violated. Since we wish to restrict our attention to slow-roll inflation, we choose to investigate the cases with $f > 2$ in this work. Large values of f yield values of n_s and r which are inconsistent with the most recent data, and hence $2 \leq f \leq 3$. ii. For a fixed value of f,

[3]It is worth noting that we have replaced ϵ_H with $\epsilon_V(\phi)$ in order to get this relation.

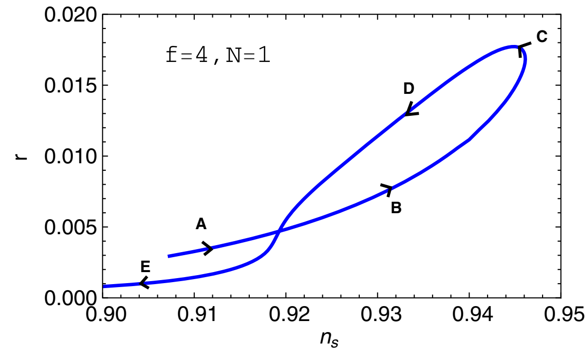

FIG. 1. For a fixed $f = 4M_{\rm Pl}$, changing the charge Q on the additional fermion leads to a trajectory in the $n_s - r$ plane parametrized by Q as shown here (for the case $\mathcal{N} = 1$). As we increase the charge Q from 0.5 to 0.71, we go from point **A** to **C** via point **B**. At **C**, there is a turning, the corresponding $Q = 0.71$. Further increasing Q from 0.71 to 1.55 takes us from **C** to the points **D** and **E** along the path shown.

the slow-roll conditions get violated if $Q \geq 2$ or $Q < 0.5$. So, we restrict ourselves to the range $0.5 \leq Q \leq 1.5$.

For a given combination of $\mathcal{N}$ and f, as one changes Q, the charge of the additional fermion in the bulk, the predictions for n_s and r change, and we get trajectories in the $n_s - r$ plane that are parametrized by Q. Although the detailed shape of the curve depends on the choice of $\mathcal{N}$ and f, for any such choice, there is typically a range of Q that will yield slow-roll inflation, and the corresponding trajectories in the $n_s - r$ plane can then be found. For example, for the case $\mathcal{N} = 1$, $f = 4$, as one increases the charge from $Q_i = 0.5$ to $Q_f = 1.55$, one gets curves of the form shown in Fig. 1. For any given f, one can find the trajectories in the $n_s - r$ plane for the various values of Q, and one can then change f and repeat this. Thus, for a given $\mathcal{N}$, one obtains a family of trajectories in the $n_s - r$ plane. Let us now look at what happens as we choose various values of $\mathcal{N}$.

1. $\mathcal{N} = 1$

The case $\mathcal{N} = 1$ corresponds to two fermions in the bulk. The charge of the first fermion has been set to $+1$, while that of the second one is Q. For this case, as shown in Fig. 2, as one plots the family of trajectories corresponding to different f and different Q, one finds that no choice of parameters brings us inside the 1σ contours.

2. $\mathcal{N} = 2$

We now argue that for the case with $\mathcal{N} = 2$, i.e., three fermions in the bulk, there exists a range of values of f and Q for which the CMB predictions improve significantly. For example, if one chooses $R = 28.7$ (in units of reduced Planck length), $f = 2.5$, and $Q = 0.582$, then one finds (for $N_* = 60$) that $n_s = 0.973$, $A_s = 2.19 \times 10^{-9}$, and $r = 0.0036$. This must be compared with the predictions

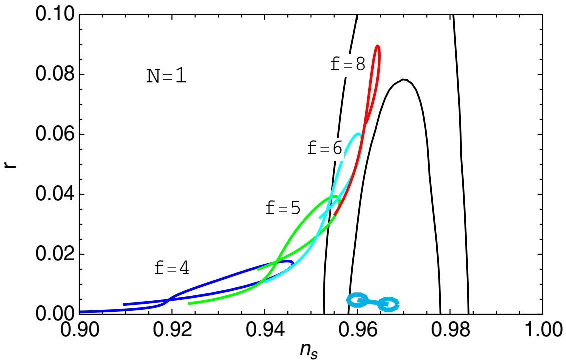

FIG. 2. When $\mathcal{N} = 1$, no matter what values of f and Q are chosen, one never obtains the values of n_s and r that are inside the 1σ contours of Planck TT,TE,EE+lowP data. Thus, there is no hope of obtaining CMB predictions similar to those of the Starobinsky model.

for the Starobinsky model (for $N_* = 54$): $r = 0.004$ while $n_s = 0.963$.

If we restrict our attention to the range

$$2.0 < f < 3, \qquad 0.5 < Q < 0.75, \tag{27}$$

first, as is shown in Fig. 3, for each value of f in the above range, there exist values of Q for which n_s lies within the 1σ allowed region. When f is in this range, increasing Q from 0.5 first increases n_s such that for a small range of values of Q n_s does fall within the observationally preferred range. It then reaches its local maximum value and then decreases. While decreasing, the value of n_s again falls within the observationally preferred range. Second, as Figs. 4 and 5 indicate, for this range of parameters, the tensor-to-scalar ratio stays below 0.02, and this is true irrespective of the values of Q and f.

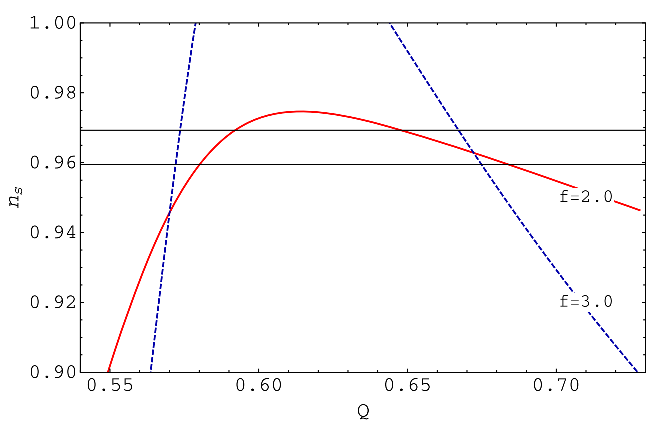

FIG. 3. The behavior of n_s as one changes Q for a fixed value of f (in units of the reduced Planck mass). The horizontal lines correspond to the 1σ limits on n_s for the 2015 Planck TT,TE,EE +lowP data.

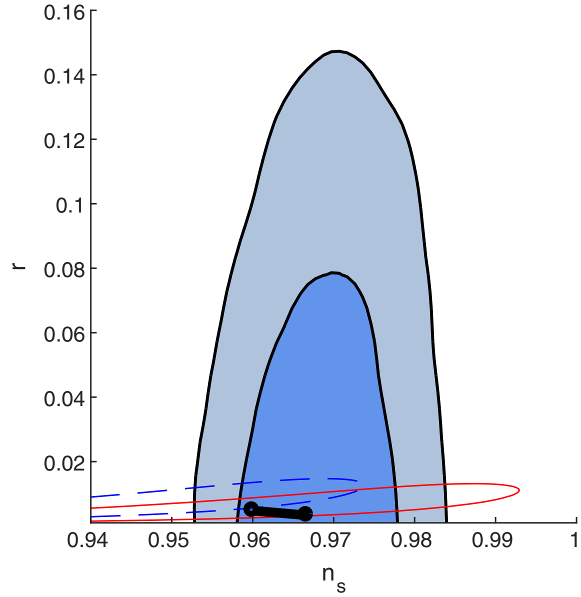

FIG. 4. The trajectories in the $n_s - r$ plane as one increases Q for a fixed f. The two curves correspond to $f = 3.0M_{\mathrm{Pl}}$ (dashed) and $f = 2.5M_{\mathrm{Pl}}$ (solid), respectively. The predictions of the Starobinsky model for n_s and r are also shown for reference. The shaded regions show the 1σ and 2σ contours for the 2015 Planck TT,TE,EE+lowP data.

In summary, when one is in the range specified by Eq. (27), r stays below 0.02; changing Q essentially changes n_s, f has a small effect on n_s and r, and (as was mentioned earlier) R essentially determines A_s. One finds that there must be some combination of f and Q that leads to n_s and r that are identical to those obtained in the Starobinsky model (see, e.g., Fig. 5).

We thus learn that adding fermions in the bulk with appropriately chosen charges can improve the CMB predictions of extranatural inflation. Moreover, CMB data suggest that $\mathcal{N} = 2$, $R \approx 29\ M_{\mathrm{Pl}}^{-1}$, $f \approx 2.5\ M_{\mathrm{Pl}}$, and $Q \approx 0.58$.

C. Constraints on derived parameters

Since the basic parameters of this scenario are constrained by CMB data, it may be a good idea to find the constraints on the other, derived parameters. But before we do so, since there are many scales in the problem, it will be a good idea to get some feel for their relative hierarchies before proceeding. Let ℓ_P be the 4D Planck length, L $(= 2\pi R)$ be the size of the extra dimension, M_p be the 4D Planck mass, and $M_p^{(5)}$ be the 5D Planck mass; then [51],

$$L = \left(\frac{M_p}{M_p^{(5)}}\right)^3 \ell_P, \tag{28}$$

and this implies that

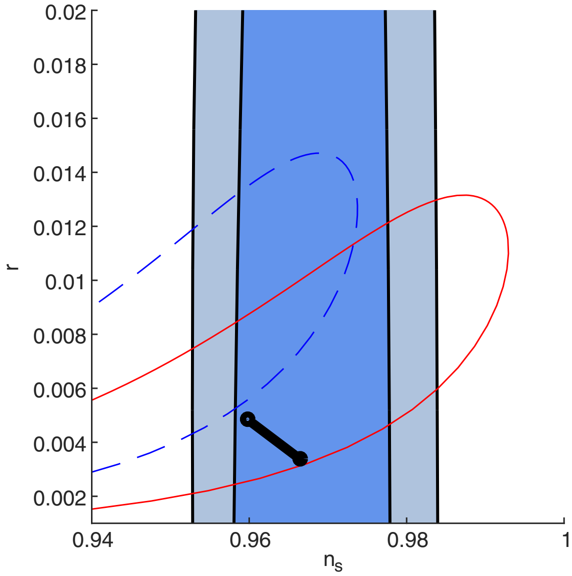

FIG. 5. This is a "zoomed-in" form of the previous figure; this shows that for f between $2.5M_{\rm Pl}$ (solid) and $3.0M_{\rm Pl}$ (dashed) and Q between 0.55 to 0.75 there exists a combination of f and Q that will give predictions identical to those of the Starobinsky model.

$$\frac{L^{-1}}{M_p^{(5)}} = \frac{L^{-1}}{M_p}\frac{M_p}{M_p^{(5)}} = \left(\frac{M_p^{(5)}}{M_p}\right)^2. \tag{29}$$

Now, in five dimensions, the energy scale of quantum gravity is $M_P^{(5)}$, while L^{-1} is the cutoff scale of 4D EFT. So, L^{-1} must be smaller than $M_p^{(5)}$, so from Eqs. (28) and (29),

$$M_p > M_p^{(5)}. \tag{30}$$

Using Eq. (10) and the fact that we typically need $f > M_p$ to be consistent with data, one concludes that

$$g_4\left(\frac{M_p}{M_p^{(5)}}\right)^3 < 1. \tag{31}$$

Since the 5D Planck mass (i.e., $M_p^{(5)}$) is smaller than the 4D Planck mass (i.e., M_p), the above inequality implies that we must have

$$g_4 \ll 1. \tag{32}$$

Moreover, since the 5D gauge coupling has a negative mass dimension, Eq. (2) implies that the unitarity bound of the theory is $E_{\rm strong} = 1/(2\pi R g_4^2)$ so that when the energy scale of any process is of this order perturbative unitarity gets violated. Thus,

$$\frac{E_{\rm strong}}{f} = \frac{1}{g_4} \gg 1. \tag{33}$$

Similarly, Eq. (6) implies that $V \sim L^{-4}$, while the Friedman equation implies that the Hubble parameter during inflation is given by $H^2 = \frac{8\pi V}{3M_p^2}$ so that

$$H = \sqrt{\frac{8\pi}{3}}\frac{1}{L^2 M_p}, \tag{34}$$

and hence

$$\frac{H}{L^{-1}} = \sqrt{\frac{8\pi}{3}}\frac{L^{-1}}{M_p} \ll 1. \tag{35}$$

The above analysis implies that

$$E_{\rm strong} \gg f \gg M_p > M_p^{(5)} > L^{-1} \gg H. \tag{36}$$

Let us now find the numerical values of all of these ratios, given the best-fit values of R and f. Obviously, $M_p = \sqrt{8\pi}M_{\rm Pl} = 5.013M_{\rm Pl}$, and the 4D gauge coupling $g_4 = (2\pi R f)^{-1}$ is given by

$$g_4 = 0.0022. \tag{37}$$

The energy scale $E_{\rm strong} = 1.13 \times 10^3 M_{\rm Pl}$, the five-dimensional Planck scale is given by $M_p^{(5)} = 0.518\, M_{\rm Pl}$, the scale $R^{-1} = 0.035\ M_{\rm Pl}$, and, finally, $H = 1.77 \times 10^{-5}\ M_{\rm Pl}$.

Any scenario in which the UV completes extranatural inflation then needs a mechanism to keep g_4 at 2.2×10^{-3} and $R \approx 30\ M_{\rm Pl}^{-1}$. Adding bulk fermions will surely destabilize the potential of the radion and cause the extra dimension to contract to Planckian values [52]. To keep the size of the extra dimension fixed to the desired value, the vev of the radius modulus must stay put at the desired value. There are many ways of dealing with this problem, and the most harmless one is to use stabilizer fields *à la* Goldberger-Wise [53] to ensure that the extra dimension stays sufficiently large. Since these fields need not be charged under the gauge group of the bulk gauge field, the inflaton potential and hence inflationary predictions will not be affected by employing this mechanism. This mechanism will work only if the radion is not excited during inflation. The condition for this is that the mass of the radion must be large as compared to H.

IV. SUMMARY AND DISCUSSION

There seems little doubt that cosmological experiments of the next generation will be able to reduce the uncertainties in various inflationary observables to a very high degree. For example, it is expected that within a decade the uncertainties in the tensor-to-scalar ratio, i.e., $\sigma(r)$, will be of the order of 0.0005 [11], while those in the running of the scalar spectral index (α_s) will be of the order of 0.0025 [14]. It thus seems that in the near future we shall uncover the shape of the inflaton potential with unprecedented accuracy. Given this optimistic state of affairs, one must ensure that the mechanisms that lead to the inflaton potential are completely trustworthy from a theoretical perspective.

The CMB observations of even the current generation have imposed tight observational constraints on many scenarios of cosmic inflation. A model of which the CMB predictions are compatible with the observations of Planck experiment is the $\mathcal{R} + \mathcal{R}^2$ model of inflation of Starobinsky [1] (see also Ref. [54]). If one considers higher-dimensional corrections to Einstein-Hilbert action, i.e., $\frac{\mathcal{R}}{2\kappa^2} + \alpha\mathcal{R}^2 + \beta\mathcal{R}_{\mu\nu}\mathcal{R}^{\mu\nu} + \gamma\mathcal{R}_{\mu\nu\rho\sigma}\mathcal{R}^{\mu\nu\rho\sigma}$, one can use the Chern-Gauss-Bonnet theorem (i.e., the fact that $\mathcal{R}^2 - 4\mathcal{R}_{\mu\nu}\mathcal{R}^{\mu\nu} + \mathcal{R}_{\mu\nu\rho\sigma}\mathcal{R}^{\mu\nu\rho\sigma}$ is a topological invariant) to eliminate the term with a Riemann tensor. Now, the Starobinsky model is based on the assumption that among κ, α, and β in

$$S = \int d^4x\sqrt{-g}\left(\frac{\mathcal{R}}{2\kappa^2} + \alpha\mathcal{R}^2 + \beta\mathcal{R}_{\mu\nu}\mathcal{R}^{\mu\nu}\right) + \cdots \tag{38}$$

one can pick only the $\frac{\mathcal{R}}{2\kappa^2} + \alpha\mathcal{R}^2$ terms and adjust the coefficients to ensure that one obtains CMB predictions. This is sensible because, among one-loop quantum corrections to the Einstein-Hilbert action, the terms other than $\frac{\mathcal{R}}{2\kappa^2} + \alpha\mathcal{R}^2$ vanish when the metric is conformally flat (e.g., a spatially flat FRW metric) as happens after a sufficient duration of inflation. However, if κ is 4D reduced Planck length, CMB observations imply that the dimensionless coefficient α is $\mathcal{O}(10^9)$. Given this, one might wonder whether assuming α to be too large and β to be negligible may appear to involve some fine-tuning, but notice that, since β is small due to a symmetry reason, its smallness may be protected from loop corrections. Similarly, if one resorts to coupling the inflaton field nonminimally to gravity in order to obtain successful CMB predictions, super-Planckian vevs may make it hard to justify including just the one higher-dimensional operator $\mathcal{R}\phi^2$ while ignoring all others (see, however, the discussion in Sec. 5.1 of Ref. [55]). But the super-Planckian field excursion in the Einstein frame field in both the cases discussed (Starobinsky as well as nonminimal coupling) suggests that one may need to include even higher-order terms. Given this, a scenario of large field inflation that does not suffer from such uncertainties and nonetheless makes successful predictions that can be tested in next-generation experiments must be seriously sought. In this work, an attempt has been made to find one such model.

Ultraviolet sensitivity of large field inflation persuades one to resort to symmetries to protect the inflaton potential. The minimal natural inflation uses global symmetries to deal with this, but since there can be no global symmetries in quantum gravity [41,42], a better alternative is to employ gauge symmetries for this task. Extranatural inflation does exactly this; the inflaton is the zero mode of the fifth component of a bulk Abelian gauge field of which the potential is generated by charged fermions present in the bulk. The required value of gauge coupling to achieve the same, however, turns out to be quite small. This small value of this 4D gauge coupling has been a topic of discussion in the literature since the failure to obtain such small values of the gauge coupling in known UV completions such as string theory [18] has inspired the famous weak gravity conjecture [19]. The exact value of gauge coupling is thus of paramount importance.

However, the most minimal version of extranatural inflation gives predictions identical to those of natural inflation, which is mildly disfavored by current CMB data. One may wonder whether one could have a variation of the minimal version of extranatural inflation which fits the data with a reasonable value of 4D gauge coupling g_4. In this work, we studied the effect of additional charged, light, fermions in the bulk on the CMB predictions of extranatural inflation. We have found that the one needs to add at least two more fermion species in the bulk in order to improve the fit to CMB data and if the radius of the extra dimensions is $R \approx 29\ M_{\rm Pl}^{-1}$, $f = 2.5\ M_{\rm Pl}$ and the charges of the additional fermions are $Q = 0.58$ (in units of charge of the fermion that generates the potential of extranatural inflation), one obtains CMB predictions very close to those of the Starobinsky model. One can readily determine the corresponding value of g_4, and it turns out to be $g_4 = 0.0022$. Since this value still turns out to be too small, we have essentially shown that, although adding more fermions can help improve the fit to CMB, this does not resolve the problem of the smallness of gauge coupling. In any case, it is this value of the 4D gauge coupling that needs to be targeted in the UV completions of extranatural inflation.

In summary, we have presented a model of cosmic inflation that has the merit that its predictions can be identical to those of the Starobinsky model and that can potentially be free from all issues of UV sensitivity, provided one can find a UV completion in which the 4D gauge coupling turns out to be as small as required. Apart from looking for appropriate UV completions, in the future, one could make use of the light bulk fermions in the model to tackle with other problems of the physics of the early Universe, e.g., matter-antimatter asymmetry or dark matter. Since the masses and charges of such fermions are already constrained by CMB data, this can be a very interesting exercise.

ACKNOWLEDGMENTS

G. G. would like to thank Koushik Dutta (Saha Institute of Nuclear Physics, Kolkata) and Tirtha Sankar Ray (Indian Institute of Technology, Kharagpur) for discussions. G. G. would also like to thank Gary Shiu (University of Wisconsin, Madison) for clarifying several issues related to extranatural inflation. J. P. would like to thank the Navajbai Ratan Tata Trust for financial support.

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