Pure Natural Inflation

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We point out that a simple inflationary model in which the axionic inflaton couples to a pure Yang-Mills theory may give the values of the quadratic potential

\[ V(\phi) = \Lambda^4 \left[ 1 - \cos \left( \frac{\phi}{f} \right) \right], \]

(2)

where \( \Lambda \) is the dynamical scale of the Yang-Mills theory. The resulting inflation model is called natural inflation [4, 5], which has been extensively studied in the literature. The potential of Eq. (2), however, is not favored by the current data, and it would soon be excluded at a higher confidence level if the bound on \( r \) improves with \( n_s \) staying at the current value; see Fig. 1.

It is known since long ago, however, that the cosine potential in Eq. (2) is not correct in general, as argued by Witten [6, 7] in the large \( N \) limit [8] with the ’t Hooft coupling \( \lambda \equiv g^2N \) held fixed.1 In particular, while the physics is periodic in \( \phi \) with the period of \( 2\pi f \) (because \( \theta \equiv \phi/f \) is the \( \theta \) angle of the Yang-Mills theory), the multi-valued nature of the potential allows for the potential of \( \phi \) in a single branch

\[ V(\phi) = N^2 \Lambda^4 \mathcal{V} \left( \frac{\lambda \phi}{8\pi^2 N f} \right), \]

(3)

not to respect the periodicity under \( \phi \to \phi + 2\pi f \). Here,
the combination
\[ x \equiv \frac{\lambda \phi}{8\pi^2 N f}, \] (4)

appearing in the argument of \( V(x) \) is determined by analyzing the large \( N \) limit. This allows for building axionic models of inflation in which the range of the field excursion exceeds the decay constant \( f \) [14–16].

The potential of Eq. (3) has an expansion of the form
\[ V(\phi) = \sum_{n=1}^{\infty} b_{2n} \left( \frac{\phi}{F} \right)^{2n}, \] (5)

where \( F \propto f \). The values of the coefficients \( b_{2n} \)—more precisely their signs and double ratios—are important for how the predictions for \( n_* \) and \( r \) change as \( F \) is varied. If the cosine potential in Eq. (2) were valid, then we would obtain
\[ \text{sgn}(b_{2n}) = (-1)^{n-1}, \] (6)

and
\[ \frac{b_2}{b_4} = \frac{2}{5}, \quad \frac{b_6}{b_4} = \frac{15}{28}, \quad \ldots, \] (7)

which lead to the curves labeled as “cosine” in Fig. 1. The correct values of the double ratios, however, are expected to be different from these values. In fact, \( b_{2n} \)’s obtained by lattice gauge theory disfavor the cosine form of Eq. (2) and are rather consistent with those expected from large \( N \) expansion [17].

While \( b_{2n} \)’s may in principle be determined by lattice calculations, their errors are still large. Instead, we may infer the form of the potential by the following arguments. First, invariance under the \( CP \) transformation \( \phi \to -\phi \) implies that \( V(x) \) is a function of \( x^2 \), where we have absorbed a possible bare \( \theta \) parameter in the definition of \( \phi \). Second, \( V(\phi) \) is expected to flatten as the potential energy approaches the point of the deconfining phase transition with increasing \( |\phi| \) (since the dynamics generating the potential will become weaker). Assuming that the potential is given by a simple power law, we thus expect \( V(x) \sim 1/(x^2)^p \) \((p > 0) \). This potential is singular at \( x \to 0 \), and a simple way to regulate it is to replace \( x^2 \) with \( x^2 + \text{const} \). After setting the minimum of the potential to be zero, these considerations give
\[ V(x) = M^4 \left[ 1 - \frac{1}{(1 + x^2)^p} \right] \quad (p > 0), \] (8)

where \( M \sim \sqrt{N\lambda} \), and \( c > 0 \) is a parameter of order unity. Here, we have used the well-established fact that the coefficient of \( x^2 \) is positive when \( V(x) \) is expanded around \( x = 0 \). We call the model of inflation in which the axionic inflaton potential is generated by a pure Yang-Mills theory (whose potential we expect to take the form of Eq. (8)) pure natural inflation.

As in the cosine potential, the potential of Eq. (8) gives \( \text{sgn}(b_{2n}) = (-1)^{n-1} \). It, however, gives different values of the double ratios
\[ \frac{b_2}{b_4} = \frac{2(p + 2)}{3(p + 1)}, \quad \frac{b_{2n}}{b_{2n+2}} = \frac{(n+1)(p+n+1)}{(n+2)(p+n)}, \quad \ldots. \] (9)

Therefore, predictions of this model are different from those of conventional natural inflation. (For example, by equating \( (b_6/b_4) / (b_4/b_2) \) we obtain \( p = -7/2 < 0 \).) Here, we have assumed that the effect of a transition between different branches can be neglected, which we will argue to be the case.

The potential of Eq. (8) can be obtained by a holographic calculation [10, 18], which is applicable in the limit of large \( N \) and ‘t Hooft coupling. In this calculation, \( N \) D4-branes in type IIA string theory are considered, with the D4-branes wrapping a circle. Below the Kaluza-Klein scale \( M_{KK} \) for the circle, the theory reduces to a 4d (non-supersymmetric) pure \( SU(N) \) Yang-Mills theory, with the dynamical scale
\[ \lambda = M_{KK} e^{-\frac{2+2}{4\pi^2}}, \] (10)

where \( \lambda \) is the ‘t Hooft coupling at \( M_{KK} \). Considering the backreaction to the geometry of the constant Wilson line of the Ramond-Ramond one-form, which represents the \( \theta \) angle of the gauge theory, the potential of the form of Eq. (8) is obtained with \( c = 1 \) and \( p = 3 \). Specifically, the potential of \( \phi \) for a single branch is given by
\[ V(\phi) = M^4 \left[ 1 - \frac{1}{\left(1 + \left(\frac{\phi}{F}\right)^2\right)^3} \right], \] (11)

where\(^2\)
\[ M^4 = \frac{\lambda N^2}{3^7 \pi^2} M_{KK}^4, \quad F = \frac{8\pi^2 N}{\lambda} f. \] (12)

The potentials for the other branches are obtained by replacing \( \phi \) with \( \phi + 2\pi k f \) \((k \in \mathbb{Z})\); see Fig. 2.

To illustrate the parameter region we consider, let us choose
\[ \lambda \simeq 8\pi^2. \] (13)

Strictly speaking, the holographic calculation is not quite valid with this value of the ‘t Hooft coupling—it requires a larger value of \( \lambda \). However, we may expect, e.g. based on the success of the AdS/QCD program [19, 20], that

\(^2\) The ‘t Hooft coupling (and the gauge coupling squared) defined in Ref. [10] is a factor of 2 smaller than \( \lambda \) (\( g^2 \)) here.
this reasonably approximates the true dynamics of the 4d Yang-Mills theory. With this choice of \( \lambda \), we find

\[
M \approx \sqrt{N} A, \quad F \approx N f,
\]

(14)
as one naively expects from dimensional and \( N \)-scaling considerations. From the analysis of Ref. [10], we expect that for sufficiently large \( N \) the effect of a transition between different branches is not important, unless \( \phi \) becomes much larger than \( F \) (which does not occur in our analysis below). To reproduce the observed amplitude of the scalar perturbation, we will need to take

\[
M \sim 10^{16} \text{ GeV}.
\]

(15)
The precise value depends on other parameters, e.g. \( F \).

The second expression in Eq. (14) implies that for \( N \gg 1 \) the characteristic scale for the field excursion, \( F \), can be much larger than the axion decay constant \( f \). This, however, does not mean that \( F \) can be much larger than the Planck scale \( M_{\text{Pl}} \sim 1.22 \times 10^{19} \text{ GeV} \). In general, the decay constant \( f \) is expected to be smaller than the field theoretic cutoff (string) scale \( M_s \)

\[
f \lesssim M_s.
\]

(16)
On the other hand, the Planck scale is related with \( M_s \), as \( M_{\text{Pl}}^2 \sim N^2 M_s^2 \) (see, e.g., Ref. [21]), so that \( N \) drops from the relation between \( F \) and \( M_{\text{Pl}} \):

\[
F \lesssim O(M_{\text{Pl}}).
\]

(17)
We argue that this is a desired feature. If \( F \gg M_{\text{Pl}} \), the inflaton potential would be well approximated by the first, quadratic term in Eq. (5), which is excluded by the data. Because of Eq. (17), however, we expect that higher terms in Eq. (5) are important. This makes the predictions of the model deviate from those of the quadratic potential \( V(\phi) = m^2 \phi^2 / 2 \).

In Fig. 1, we plot the values of \( n_s \) and \( r \) predicted by the potential of Eq. (11). For \( F \gg M_{\text{Pl}} \), the predictions approach those of the quadratic potential, denoted by the black dots for \( e \)-folding \( N_e = 50 \) (left) and 60 (right).

\[
V(\phi) = M^4 \left[ 1 - \frac{1}{\left( 1 + (\phi / f)^2 \right)^p} \right],
\]

(18)
and, as in Fig. 1, varied \( F/M_{\text{Pl}} \) in the range between 100 and 0.1, with the dots representing \( F/M_{\text{Pl}} = 10, 5, 1 \) indicated by the red dots (from top to bottom). We find that the model gives the values of \( n_s \) and \( r \) consistent with the data at the 95\% (68\%) CL for \( F/M_{\text{Pl}} \lesssim 3.3 \) (0.7) for \( N_e = 50 \) and \( F/M_{\text{Pl}} \lesssim 6.8 \) (4.4) for \( N_e = 60 \). Given Eq. (17), this is quite satisfactory.

In Fig. 3, we plot the predictions arising from the potential in which \( p \) in Eq. (8) takes more general values \( p = 1, 2, 3, 4, 10 \). We have parameterized the potentials as

As \( F \) decreases, however, they deviate from these values. In particular, the tensor-to-scalar ratio \( r \) decreases while \( n_s \) being (almost) kept, as represented by the lines indicated as “pure natural.” This is because higher terms in Eq. (5) start contributing.\(^3\) In the figure, we have varied \( F \) from 100\(M_{\text{Pl}} \) to 0.1\(M_{\text{Pl}} \), with the predictions for \( F/M_{\text{Pl}} = 10, 5, 1 \) indicated by the red dots (from top to bottom). We find that the model gives the values of \( n_s \) and \( r \) consistent with the data at the 95\% (68\%) CL for \( F/M_{\text{Pl}} \lesssim 3.3 \) (0.7) for \( N_e = 50 \) and \( F/M_{\text{Pl}} \lesssim 6.8 \) (4.4) for \( N_e = 60 \). Given Eq. (17), this is quite satisfactory.

\(^{3}\) A special case of \( V(\phi) \approx A - B/\phi^6 \) leading to \( n_s \approx 0.965 \) and \( r \approx 8 \times 10^{-4} \) was discussed in Ref. [16].
Finally, we mention that a large value of $N$—despite the fact that it does not help to make $F$ larger than $M_{Pl}$—can make the decay constant $f$ smaller than $F$; see Eq. (14). This allows for enhancing couplings of the inflaton $\phi$ to the standard model gauge fields for fixed $F$, i.e. for a fixed inflation potential. This in turn allows for raising reheating temperature $T_R$. The reheating temperature is given by $T_R \simeq 0.2 \sqrt{\Gamma_\phi M_{Pl}}$, where

$$\Gamma_\phi \approx \frac{nm_\phi^3}{4096\pi^2 f^2},$$  \hspace{1cm} (19)$$

is the inflaton decay width. Here,

$$m_\phi = \sqrt{2\rho} \frac{M^2}{F} \simeq \frac{10^{13}}{F/M_{Pl}} \text{ GeV},$$  \hspace{1cm} (20)$$
is the inflaton mass, and $n$ is the final state multiplicity. ($n = 12$ if $\phi$ decays only to the standard model gauge fields.) This yields

$$T_R \sim 10^9 \text{ GeV} \left( \frac{N}{10} \right) \left( \frac{0.5}{F/M_{Pl}} \right)^{5/2},$$  \hspace{1cm} (21)$$
The model can, therefore, be made consistent with thermal leptogenesis, which requires $T_R \gtrsim 2 \times 10^9$ GeV [22], even for $F/M_{Pl} \approx 0(1)$.

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