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Towards axionic Starobinsky-like inflation in string theory ☆

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ABSTRACT

It is shown that Starobinsky-like potentials can be realized in non-geometric flux compactifications of string theory, where the inflaton involves an axion whose shift symmetry can protect UV-corrections to the scalar potential. For that purpose we evaluate the backreacted, uplifted F-term axion-monodromy potential, which interpolates between a quadratic and a Starobinsky-like form. Limitations due to the requirements of having a controlled approximation of the UV theory and of realizing single-field inflation are discussed.

1. Introduction

The recent release of the PLANCK 2015 data provides improved experimental results and bounds on the ΛCDM cosmology [1]. In particular, the BICEP2 observation [2] of a tensor-to-scalar ratio as big as $r = 0.2$ can now be completely explained by a foreground dust contamination of the signal, and is replaced by the upper bound $r < 0.113$. Moreover, for the spectral index PLANCK 2015 reports $n_s = 0.9667 \pm 0.004$ and for its running $\alpha_s = -0.002 \pm 0.013$.

As a consequence, large-field inflationary potentials of the type $V \sim \Theta^p$ are essentially ruled out for $p \geq 2$, and the currently best class of models fitting the data is plateau-like [3]. This class contains the Starobinsky model [4], as well as more general Starobinsky-like models

$$V(\Theta) = M_{\rm Pl}^4 \Big(A - B\, e^{-\gamma\, \Theta} \Big) , \tag{1}$$

(see also [5], as well as [6] for a historical perspective on the Starobinsky model). Starobinsky-like models have been constructed in string theory in the LARGE volume scenario (LVS), where the role of the inflaton is played by a canonically normalized Kähler modulus [7,8].

When working with a model of large-field inflation, Planck suppressed higher-order operators need to be controlled, since otherwise they lead to an η-problem. For the LVS, corrections are suppressed by an exponentially large volume, while in the case of the inflaton being an axion, the shift symmetry of the latter can protect the potential against perturbative corrections. Various scenarios for axion inflation have been constructed, such as natural inflation [9], N-flation [10], or aligned inflation [11].

Another promising string-theoretic approach, still allowing for some control over the higher-order corrections, is axion monodromy inflation [12,13], for which a field theory version has been proposed in [14] (for a review see [15]). In [16–18] this scenario has been realized via the F-term scalar potential induced by background fluxes, which has the advantage that supersymmetry is broken spontaneously by the very same effect by which usually moduli are stabilized. Such models were studied in [19,20] for the possibility to provide a quadratic potential for the axion.

In this letter, we analyze a toy model for the flux-induced scalar potential for a large excursion of a prospective axion/inflaton. For concreteness, we consider a simple type IIB superstring flux compactification, where a superpotential for all moduli is generated by turning on NS–NS and R–R three-form fluxes as well as non-geometric fluxes. Following [21], taking the backreaction of the stabilized moduli onto the evolution of the inflaton into account, we find the expected flattening of the uplifted potential, which af-

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ter canonical normalization interpolates between a quadratic and a Starobinsky-like form. Here we discuss the cosmological consequences of this model, whereas more details on the formal framework and on the phenomenology will be discussed elsewhere [22].

2. Large-field inflation

Let us recall the expressions of the cosmological parameters for the large-field polynomial and Starobinsky-like inflationary models. For polynomial inflation with $V \sim \Theta^p$, the slow-roll parameters $\epsilon = \frac{1}{2}(\frac{V'}{V})^2$ and $\eta = \frac{V''}{V}$ can be written in the following way

$$\epsilon = \frac{1}{2}\frac{p^2}{\Theta^2}, \qquad \eta = \frac{p\,(p-1)}{\Theta^2}, \tag{2}$$

and the number of e-foldings is expressed as

$$N_e = \int_{\Theta_{\rm end}}^{\Theta_*} \frac{V}{V'}\, d\Theta = \frac{1}{p}\int_{\Theta_{\rm end}}^{\Theta_*} \Theta\, d\Theta \simeq \frac{\Theta_*^2}{2\,p}. \tag{3}$$

This implies that $N_e \simeq \frac{p}{4\epsilon}$, and the spectral index, its running and the tensor-to-scalar ratio are obtained as

$$n_s = 1 + 2\,\eta - 6\,\epsilon \sim 1 - \frac{p+2}{2N_e},$$
$$\alpha_s = -\frac{(p-1)(9\,p-14)}{2\,N_e^2}, \qquad r = 16\,\epsilon \sim \frac{4p}{N_e}. \tag{4}$$

For the Starobinsky-like model (1), the slow-roll parameters become

$$\epsilon = \frac{1}{2\,\gamma^2}\,\eta^2, \qquad \eta = -\frac{1}{N_e}, \tag{5}$$

so that

$$n_s = 1 - \frac{2}{N_e}, \qquad \alpha_s = -\frac{2}{N_e^2}, \qquad r = \frac{8}{(\gamma\, N_e)^2}. \tag{6}$$

Independently of the parameters A, B and γ, for $N_e = 60$ e-foldings this gives the experimental value $n_s \sim 0.967$. Note that n_s and α_s in (6) agree with the values for a quadratic potential in (4), except that the tensor-to-scalar ratio comes out smaller.

The amplitude of the scalar power spectrum takes the experimental value $\mathcal{P} = (2.142 \pm 0.049) \cdot 10^{-9}$, and can be expressed as

$$\mathcal{P} \sim \frac{H_{\rm inf}^2}{8\,\pi^2\,\epsilon\, M_{\rm Pl}^2}. \tag{7}$$

From this one can extract the Hubble constant during inflation, and consequently the mass of the inflaton via $M_\Theta^2 = 3\,\eta\, H^2$. The relation $V_{\rm inf} = 3\, M_{\rm Pl}^2\, H_{\rm inf}^2$ then fixes the mass scale of inflation.

3. Fluxes and moduli stabilization

We now turn to the framework of type IIB orientifolds on Calabi–Yau (CY) threefolds, equipped with geometric and non-geometric fluxes. The NS–NS and R–R fluxes H_3 and F_3 generate a potential for the complex-structure and axio-dilaton moduli, where the latter is written as $S = s + ic$ with $s = \exp(-\phi)$ and c denoting the R–R zero-form. Non-geometric Q-fluxes can generate a tree-level potential for the Kähler moduli $T_\alpha = \tau_\alpha + i\rho_\alpha$, where τ_α denotes a four-cycle volume (in Einstein frame) and ρ_α is the R–R four-form reduced on that cycle. The details for such flux compactifications have been worked out in [23–26] (see also [27]). The resulting scalar potential reads

$$V = \frac{M_{\rm Pl}^4}{4\pi}\, e^K \left(K^{i\bar{j}} D_i W\, D_{\bar{j}} \overline{W} - 3\,|W|^2\right), \tag{8}$$

which is computed from the Kähler potential

$$K = -\log\!\left(-i\int \Omega \wedge \overline{\Omega}\right) - \log\!\left(S + \overline{S}\right) - 2\log \mathcal{V} \tag{9}$$

and the flux-induced superpotential

$$W = -\left(\mathfrak{f}_\lambda X^\lambda - \tilde{\mathfrak{f}}^\lambda F_\lambda\right) + i\, S\left(h_\lambda X^\lambda - \tilde{h}^\lambda F_\lambda\right)$$
$$+\, i\, T_\alpha \left(q_\lambda{}^\alpha X^\lambda - \tilde{q}^{\lambda\alpha} F_\lambda\right). \tag{10}$$

Note that here $\mathcal{V}$ denotes the volume of the Calabi–Yau manifold (in Einstein frame), and that we have assumed a large-volume and small string-coupling regime so that higher-order corrections can safely be ignored. As usual, X^λ and F_λ denote the periods of the holomorphic three-form Ω, and $\{\mathfrak{f}, \tilde{\mathfrak{f}}\}$, $\{h, \tilde{h}\}$ and $\{q, \tilde{q}\}$ denote the flux quanta of F_3, H_3 and Q.

To be more specific, let us consider a simple case of a CY manifold with no complex-structure moduli and just one Kähler modulus. One might think of it as an isotropic six-torus with fixed complex structure. In this case the Kähler potential is given by

$$K = -3\log(T + \overline{T}) - \log(S + \overline{S}). \tag{11}$$

Next, we turn on fluxes to generate the superpotential

$$W = -i\,\tilde{\mathfrak{f}} + i\, h\, S + i\, q\, T, \tag{12}$$

with $\tilde{\mathfrak{f}}, h, q \in \mathbb{Z}$. The resulting scalar potential in units of $M_{\rm Pl}/(4\pi)$ reads

$$V = \left(\frac{(hs + \tilde{\mathfrak{f}})^2}{16s\,\tau^3} - \frac{6hqs - 2q\tilde{\mathfrak{f}}}{16s\,\tau^2} - \frac{5q^2}{48s\,\tau}\right) + \frac{\theta^2}{16s\,\tau^3}, \tag{13}$$

which only depends on s, τ, and the linear combination of axions $\theta = h\,c + q\,\rho$. One linear combination of axions is not stabilized by (13), but can receive a mass from non-perturbative effects. Such ultra-light axions can become part of dark radiation [28].

In [19,22] a mechanism to realize axion inflation together with moduli stabilization in string theory has been proposed. There the idea was to first stabilize all moduli except one axion by turning on fluxes proportional to a large parameter λ, and in a second step stabilize the massless axion by introducing a deformation depending on additional small fluxes. The resulting superpotential takes the schematic form

$$W_{\rm ax} = \lambda\, W + f_{\rm ax}\, \Delta W. \tag{14}$$

Instead of analyzing the resulting potential for the rather complicated fully fledged models presented in [22], in this letter we mimic the resulting structure of the scalar potential by introducing a flux parameter λ in (13) as

$$V = \lambda^2\left(\frac{(hs + \tilde{\mathfrak{f}})^2}{16s\,\tau^3} - \frac{6hqs - 2q\tilde{\mathfrak{f}}}{16s\,\tau^2} - \frac{5q^2}{48s\,\tau}\right) + \frac{\theta^2}{16s\,\tau^3}. \tag{15}$$

We consider this as a (partly exactly solvable) toy model to analyze the possibility of realizing large-field inflation in string theory.

Solving now $\partial_s V = \partial_\tau V = \partial_\theta V = 0$, we find three solutions. Besides the supersymmetric AdS minimum with a tachyonic mode, there exists a non-supersymmetric, tachyon-free AdS minimum at

$$\tau_0 = \frac{6\,\tilde{\mathfrak{f}}}{5\,q}, \qquad s_0 = \frac{\tilde{\mathfrak{f}}}{h}, \qquad \theta_0 = 0. \tag{16}$$

To ensure $\tau_0, s_0 > 0$, for definiteness we chose all flux-quanta to be positive. Furthermore, $\tilde{\mathfrak{f}}/h \gg 1$ and $\tilde{\mathfrak{f}}/q \gg 1$ implies weak string