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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

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

(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 $S = s + i\,c$ 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

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

and the flux-induced superpotential

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

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 coupling and large radius, so we can ignore higher-order corrections to the scalar potential. The above fluxes also induce D3- and a D7-brane tadpole charges, $N_{D3} = -\lambda^2 \tilde{\mathfrak{f}} h$ and $N_{D7} = -\lambda^2 \tilde{\mathfrak{f}} q$, that need to be canceled by D-branes.

To compute the mass eigenvalues and eigenstates in the canonically-normalized basis, we consider the matrix $M^i{}_j = K^{ik} V_{kj}$, with $V_{ij} = \frac{1}{2}\partial_i \partial_j V$, evaluated at the minimum. We then find mass eigenvalues

$$M^2_{\text{mod},i} = \mu_i \, \frac{\lambda^2 h q^3}{16 \tilde{\mathfrak{f}}^2} \frac{M_{\text{Pl}}^2}{4\pi}, \tag{17}$$

with the numerical factors

$$\mu_i = \left(0, \tfrac{185}{54\lambda^2}, \tfrac{25(17-\sqrt{97})}{108}, \tfrac{25(17+\sqrt{97})}{108}\right). \tag{18}$$

The first two eigenstates are axionic while the last two are saxionic. In particular, the massless state corresponds to an axionic linear combination. Note that for sufficiently large λ, the axion can be parametrically lighter than the two saxions.

Since the volume and the dilaton are fixed by fluxes, we can explicitly evaluate the string scale as

$$M_s = \frac{\sqrt{\pi}\, M_{\text{Pl}}}{s^{\frac{1}{4}} (2\tau)^{\frac{3}{4}}}. \tag{19}$$

Recalling then (17) and (16), we derive the ratio

$$\frac{M_s}{M_{\text{mod},i}} = \frac{13.03}{\sqrt{\mu_i}} \frac{1}{\lambda h^{\frac{1}{4}} q^{\frac{3}{4}}}. \tag{20}$$

4. Axion monodromy inflation

We now consider the backreaction effect of a slowly rolling and sufficiently light axion θ, i.e. we take into account that during the rolling the moduli τ and s adjust adiabatically. Solving the extremum conditions for a non-vanishing value of θ, we find

$$\tau_0(\theta) = \frac{3}{20q}\left(4\tilde{\mathfrak{f}} + \sqrt{10\left(\tfrac{\theta}{\lambda}\right)^2 + 16\tilde{\mathfrak{f}}^2}\right),$$
$$s_0(\theta) = \frac{1}{4h}\sqrt{10\left(\tfrac{\theta}{\lambda}\right)^2 + 16\tilde{\mathfrak{f}}^2}. \tag{21}$$

For large λ, the motion in the full four-dimensional field space is well-described by (21); for λ of order one the trajectories differ, but qualitatively our results are still valid. Note also that for large excursions of θ, the values of τ_0 and s_0 are in the perturbative regime, so that higher-order α'- and g_s-corrections to the scalar potential are under control. Using (21) in the potential (15) and performing a constant uplift to vanishing cosmological constant in the minimum, gives the following backreacted effective inflaton potential (in units of $M_{\text{Pl}}^2/4\pi$)

$$V_{\text{back}}(\theta) = \frac{25\lambda^2 h q^3}{108\tilde{\mathfrak{f}}^2} \frac{5\left(\frac{\theta}{\lambda}\right)^2 - 4\tilde{\mathfrak{f}}\left(4\tilde{\mathfrak{f}} - \sqrt{10\left(\frac{\theta}{\lambda}\right)^2 + 16\tilde{\mathfrak{f}}^2}\right)}{\left(4\tilde{\mathfrak{f}} + \sqrt{10\left(\frac{\theta}{\lambda}\right)^2 + 16\tilde{\mathfrak{f}}^2}\right)^2}. \tag{22}$$

Note that the initial simple quadratic potential is changed; the expected flattening of the potential becomes evident in Fig. 1. For small values of θ the potential still takes a quadratic form, whereas for large values of θ it becomes hyperbolic. In the intermediate regime there is a turning point, around which the potential is linear. We remark that a non-constant uplift term of the form $V_{\text{up}} = \varepsilon/\tau^\beta$ with β small is also possible.

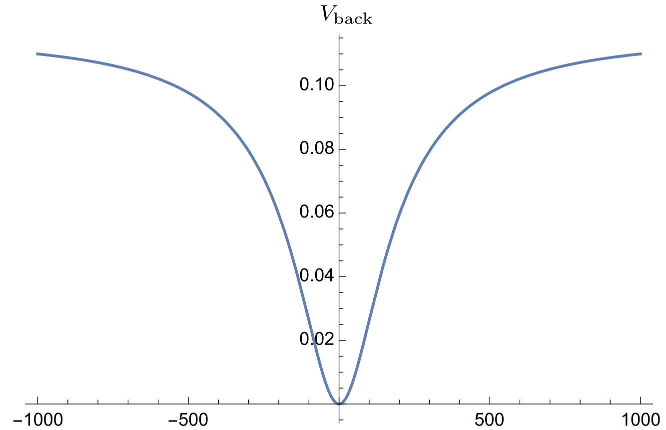

Fig. 1. The potential $V_{\text{back}}(\theta)$ shown in (22) in units of $M_{\text{Pl}}^4/(4\pi)$ for fluxes $h = 1$, $q = 1$, $\tilde{\mathfrak{f}} = 10$ and $\lambda = 10$. For this large value of λ, the trajectory (21) correctly describes the full motion in field space.

Computing the mass eigenvalues for each value of θ, we find that for large θ the eigenvalue along the trajectory becomes tachyonic while the two transversal ones are positive. Note also that for $\lambda \gg 1$, the mass hierarchy (17) remains intact for $\theta \neq 0$.

Let us now analyze the potential for the canonically-normalized inflaton in more detail. For $\theta/\lambda \ll \tilde{\mathfrak{f}}$ the shift in the value of the minimum (21) is small and the potential takes a quadratic form

$$V_{\text{back}}(\theta) \approx \frac{125 h q^3}{3456 \tilde{\mathfrak{f}}^4} \theta^2. \tag{23}$$

Employing (21), the total kinetic energy

$$\mathcal{L}_{\text{kin}} = 3\left(\frac{\partial \tau}{2\tau}\right)^2 + \left(\frac{\partial s}{2s}\right)^2 + 3\left(\frac{\partial \rho}{2\tau}\right)^2 + \left(\frac{\partial c}{2s}\right)^2, \tag{24}$$

determines the kinetic term for θ. To find the latter, we need to determine for every value of θ the orthogonal combination σ of axions. This can be fixed as

$$\partial\theta = h\,\partial c + q\,\partial\rho, \qquad \partial\sigma = -\frac{q}{s^2}\partial c + \frac{3h}{\tau^2}\partial\rho, \tag{25}$$

so that the axionic terms in (24) become

$$\mathcal{L}^{\text{ax}}_{\text{kin}} = \frac{3(\partial\theta)^2 + \tau^2 s^2 (\partial\sigma)^2}{4(3h^2 s^2 + q^2\tau^2)}. \tag{26}$$

For small θ this leads to a θ-dependence of the form $\mathcal{L}_{\text{kin}} \approx \frac{25}{148\tilde{\mathfrak{f}}^2}(\partial\theta)^2$, and thus the canonically-normalized inflaton takes the form $\Theta \approx \sqrt{25/74}\,\theta/\tilde{\mathfrak{f}}$. Note also that $\theta/\lambda \ll \tilde{\mathfrak{f}}$ implies $\Theta \ll \lambda$.

Next, we consider the large-field regime $\theta/\lambda \gg \tilde{\mathfrak{f}}$. Here we expand the backreacted potential (22) as

$$V_{\text{back}}(\theta) \approx \frac{25}{216}\frac{h q^3 \lambda^2}{\tilde{\mathfrak{f}}^2} - \frac{20}{27}\frac{h q^3 \lambda^4}{\theta^2}. \tag{27}$$

We also approximate (21) by $\tau_0(\theta) \approx \frac{3}{2\sqrt{10}\,q\,\lambda}\theta$ and $s_0(\theta) \approx \frac{\sqrt{10}}{4h\lambda}\theta$. Then, taking into account all fields from (24), we derive

$$\mathcal{L}_{\text{kin}} \approx \frac{2}{\gamma^2}\left(\frac{\partial\theta}{\theta}\right)^2 + \frac{15}{896 h^2 q^2 \lambda^2}\theta^2(\partial\sigma)^2, \tag{28}$$

with $\gamma^2 = 28/(14 + 5\lambda^2)$. Note that γ is independent of the fluxes, but it depends on λ. It can also be shown that, for appropri-

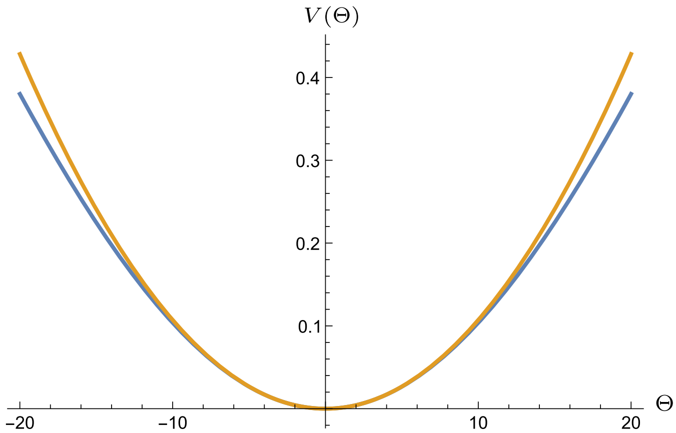

Fig. 2. The potentials $V_{\text{back}}(\Theta)$ and (23) in units of $M_{\text{Pl}}^4/(4\pi)$ for fluxes $h=1$, $q=1$, $\tilde{\mathfrak{f}}=10$ and $\lambda=60$. The lower (blue) curve is the exact backreacted potential. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

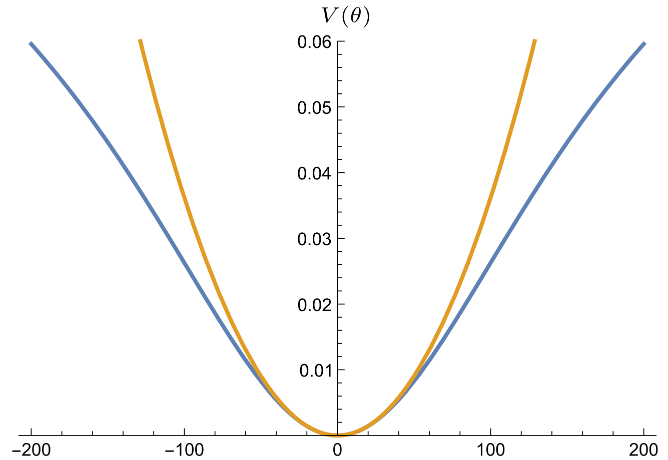

Fig. 3. The potentials $V_{\text{back}}(\theta)$ and (23) in units of $M_{\text{Pl}}^4/(4\pi)$ for fluxes $h=1$, $q=1$, $\tilde{\mathfrak{f}}=10$ and $\lambda=10$.

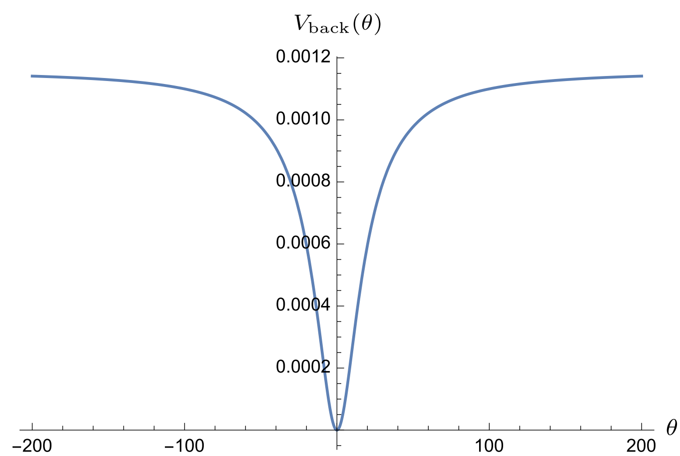

Fig. 4. The potential $V_{\text{back}}(\theta)$ in units of $M_{\text{Pl}}^4/(4\pi)$ for fluxes $h=1$, $q=1$, $\tilde{\mathfrak{f}}=10$ and $\lambda=1$.

ate initial conditions, $\partial\sigma$ vanishes along the trajectory. Canonically normalizing the inflaton via

$$\theta = 4\sqrt{\frac{2}{5}}\,\tilde{\mathfrak{f}}\,\lambda\,\exp\left(\frac{\gamma}{2}\,\Theta\right), \tag{29}$$

the potential in the large-field regime becomes Starobinsky-like

$$V_{\text{back}}(\Theta) = \frac{25}{216}\,\frac{h\,q^3\,\lambda^2}{\tilde{\mathfrak{f}}^2}\left(1-e^{-\gamma\,\Theta}\right). \tag{30}$$

We emphasize that in contrast to the potential (23) in the small-field regime, the potential (30) is exponential due to the backreaction.

In the intermediate regime $\theta/\lambda \approx 1$, it is not possible to take the canonical normalization into account analytically. However, we show below that the full backreacted potential interpolates between a quadratic and a Starobinsky-like form.

5. Qualitative picture of inflation

Let us discuss how inflation can take place in this set-up, and how its features depend on the value of λ. Here our intuition has to be based on the potential $V(\theta)$ for the non-canonically normalized potential, whereas a more accurate computation is presented in Section 6.

Quadratic inflation For sufficiently large λ the backreacted potential (22) is well approximated by the quadratic term for the region $0<\Theta<15$, i.e. slowly rolling down the potential one collects 60 e-foldings. As is illustrated in Fig. 2, this is expected to happen for $\lambda \gtrsim 60$. In this case, the inflaton is the lightest state and the heavy moduli having masses of the order $M_{\text{mod}} \sim \lambda\,M_\Theta$, which is larger than the Hubble scale $H \sim \sqrt{2N_e/3}\,M_\Theta \sim 6.32\,M_\Theta$. Therefore, we have a model of single-field inflation and all predictions agree with the ones of chaotic inflation, in particular $r \sim 0.133$. However, for such a large value of λ, the relation (20) implies that the string scale becomes smaller than the heavy moduli masses. Therefore, from the UV-complete point of view, using the effective supergravity approximation becomes questionable.

Linear inflation Lowering the value of λ, the non-trivial backreaction becomes more and more relevant, that is the potential becomes flatter in the large-field region. For $\lambda=10$ the full potential and the quadratic approximation are shown in Fig. 3. Thus it is expected that most of the 50–60 e-foldings occur along the approximately linear potential. A more precise computation would require the determination of the canonically normalized inflaton. However, the expectation is that the tensor-to-scalar ratio becomes smaller, namely $r \sim 0.08$ for linear inflation. The tension between the string scale and the heavy moduli masses becomes weaker, while the heavy masses come closer to the Hubble scale.

Starobinsky-like inflation As Fig. 4 shows, for $\lambda = O(1)$ the number of e-foldings mainly would occur on the Starobinsky-like plateau. In this case the tensor-to-scalar ratio becomes even smaller and approaches the value $r \sim 0.0015$. However, even though the heavy-moduli masses are lighter than the string scale, they are now even lighter than the Hubble scale. Therefore, this is not a model of single-field inflation and the discussion of the inflationary trajectory becomes more involved.

6. Numerical analysis

In this section, we numerically evaluate the tensor-to-scalar ratio and the number of e-foldings, taking also the kinetic term into account. We see that qualitatively, the intuition from the previous section is confirmed.

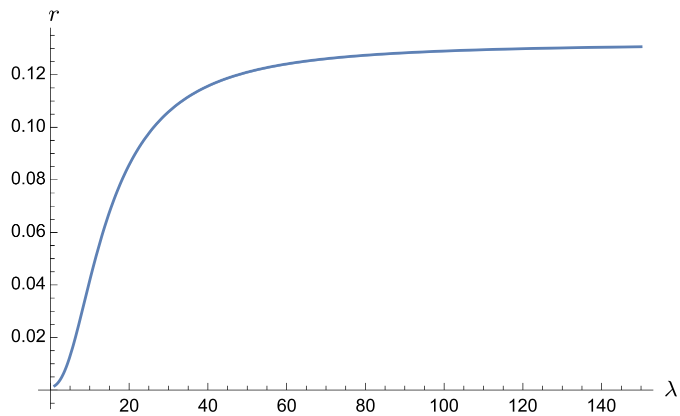

Fig. 5. The tensor-to-scalar ratio as a function of λ for fixed $n_s = 0.967$.

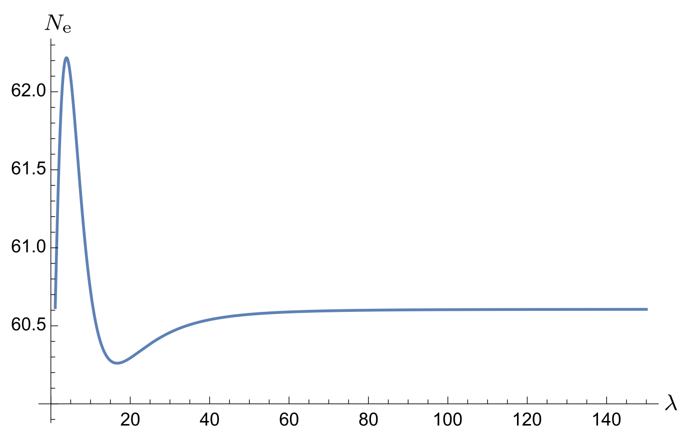

Fig. 6. The number of e-foldings as a function of λ for fixed $n_s = 0.967$.

Starting from the kinetic terms (24), we can determine an effective Lagrangian for the field θ of the form

$$\mathcal{L} = \frac{1}{2} f(\theta)\,(\partial\theta)^2 + V(\theta)\,. \tag{31}$$

Expressing the Lagrangian in terms of the canonically-normalized field is not always possible analytically, but one can determine the slow-roll parameters also in terms of θ via

$$\epsilon = \frac{1}{2f}\left(\frac{V'}{V}\right)^2, \qquad \eta = \frac{V''}{fV} - \frac{f'V'}{2f^2V}\,, \tag{32}$$

where the prime denotes the derivative with respect to θ. The number of e-foldings is given by

$$N_e = \int_{\theta_{\rm end}}^{\theta^*} d\theta\, \frac{fV}{V'}\,. \tag{33}$$

To evaluate $f(\theta)$ we substitute (21) and (26) in (24). We can then numerically determine the tensor-to-scalar ratio in terms of λ (for fixed fluxes) by fixing $n_s = 0.967$. The resulting behavior is displayed in Fig. 5, whereas Fig. 6 shows the corresponding number of e-foldings.

The curves show the expected behavior, namely that with decreasing λ the model changes from chaotic to Starobinsky-like inflation.

7. Conclusions

In a simple string compactification with two complex moduli, after introducing by hand a scaling parameter λ, we were able to stabilize all moduli except a single axion using NS–NS and R–R three-form flux together with non-geometric Q-flux. The hierarchically-light but massive axion served as an inflaton candidate. Taking into account the backreaction and assuming an uplift to Minkowski, we evaluated the resulting potential, which turned out to interpolate between a quadratic and a Starobinsky-like potential. We analyzed the cosmological consequences for three different regimes of λ. Depending on λ, the tensor-to-scalar ratio interpolates between the one for chaotic and the one for Starobinsky-like inflation.

From a controllable UV-complete theory point of view, large-field inflation models require a hierarchy of the form

$$M_{\rm Pl} > M_{\rm s} > M_{\rm KK} > M_{\rm mod} > H_{\rm inf} > M_\Theta\,, \tag{34}$$

where neighboring scales can differ by (only) a factor of $O(10)$. Our main observation was: the larger λ, the more difficult it becomes to separate the high scales on the left of (34). Contrarily, for small λ, the smaller (Hubble-related) scales on the right of (34) become difficult to separate.

Corrections to the scalar potential are expected to be under control, due to the shift symmetry of the axion/inflaton and due to the adiabatic adjustment of the saxionic moduli into the perturbative regime. Our string-motivated analysis shows how a Starobinsky-like inflation model could arise from string-theory axions. For a more realistic scenario, however, more string model building is needed, including the introduction of an MSSM-like D7-brane set-up and the computation of soft-supersymmetry breaking terms [22]. Note also that since the inflaton in our model is a linear combination of the universal axion and a Kähler axion, we can realize the stringy reheating mechanism proposed in [17].

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