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Axion Monodromy Inflation with Warped KK-Modes

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Abstract

We present a particularly simple model of axion monodromy inflation: Our axion is the lowest-lying KK-mode of the RR-2-form-potential C_2 in the standard Klebanov-Strassler throat. One can think of this inflaton candidate as being defined by the integral of C_2 over the S^2 cycle of the throat. It obtains an exponentially small mass from the IR-region in which the S^2 shrinks to zero size. Crucially, the S^2 cycle has to be shared between two throats, such that the second locus where the S^2 shrinks is also in a warped region. Well-known problems like the potentially dangerous back-reaction of brane/antibrane pairs and explicit supersymmetry breaking are not present in our scenario. The inflaton back-reaction on the geometry turns out to be controlled by the string coupling g_s. We hope that our setting is simple enough for many critical consistency issues of large-field inflation in string theory to be addressed at a quantitative level.

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1 Introduction

An important question in string cosmology is whether string theory compactifications allow for large-field inflation. On the one hand, many proposals for realizing inflation in string theory exist. At the same time, no-go theorems for large-field inflation have been put forward in various corners of the string theory landscape [1–14]. By studying large-field inflation in string theory one may thus hope to learn about fundamental properties of string theory compactifications.

Furthermore, observation may force us to address these questions. For models of single-field slow-roll inflation, there is a direct link between the tensor-to-scalar ratio r and the nature of inflation. To achieve $r \gtrsim 0.01$ the inflaton has to traverse a trans-Planckian field range during inflation, thus requiring inflation to be of large-field type [15]. The most recent observational constraint by BICEP2 and the Keck Array on the tensor-to-scalar ratio is $r \leq 0.07$ at 95 % confidence [16], which is compatible with large-field inflation. Currently, considerable effort is being expended towards more precise measurements of r.

One challenge faced by models of large-field inflation is their sensitivity to an infinite tower of corrections to the inflaton potential. One way of controlling these corrections is to identify the inflaton with an axion-like field (henceforth axion), so that the shift symmetry of the axion protects the potential from dangerous corrections. A promising approach for realizing axion inflation in string theory is axion monodromy inflation inflation [17, 18][1]. By introducing a monodromy the periodic field space of the axion is effectively unfolded, while the underlying periodicity of the theory continues to protect the inflaton potential from corrections. Further, an effective trans-Planckian field range for the inflaton can be achieved in theories involving more than one axion [24–26]. See [27] for a review including advances until 2014. For more recent progress and further references see [28].

A monodromy for axions can be induced by couplings to branes [17,18] (see [29] for very recent progress), but also due to background fluxes [30] (recently established in the supergravity context under the name of F-term axion monodromy inflation [31–33]). All these approaches are not without their problems. For example, the original axion monodromy inflation constructions employ setups with both branes and anti-branes [17, 18]. As a result, in addition to the issue of back-reaction of the inflaton, the problem of brane-anti-brane back-reaction has to be addressed [34, 35] (see also [36]). Such models then require complicated warped throat geometries, which has hampered further quantitative studies of these constructions. Recent progress towards realizing such warped geometries has been made in [37], where a $Z_2 \times Z_3$-orbifold of the conifold is used.

The situation is better in axion monodromy inflation models employing background fluxes, as the tools of flux compactifications can be used to examine these proposals in more detail.

[1] See [19, 20] for an early, purely field theoretic version and [21–23] for a more recent, closely related string-theoretic proposal.

In [38] it was shown that models of axion monodromy inflation inflation in the complex structure moduli sector of Calabi-Yau 3- and 4-folds require a significant level of tuning to avoid excessive back-reaction and the destabilization of Kähler moduli. The required level of tuning can only be achieved in 4-folds which further complicates the model. These difficulties can be avoided if Kähler moduli are stabilized using non-geometric fluxes [39–41]. However, it remains a challenge to implement a consistent hierarchy of scales in the resulting models.

Given the technical difficulties encountered in most constructions of axion monodromy inflation inflation, it would be desirable to realize as minimal a model of axion monodromy inflation inflation as possible. In such a simple construction one may hope that questions regarding the consistency and detailed phenomenology can be addressed explicitly and quantitatively. This is what we set out to do in this work. Here, we present a simple model of axion monodromy inflation inflation which is based on the standard Klebanov-Strassler-throat [42] (i.e. the deformed conifold) with shrinking S^2. Our axion is the RR-2-form C_2 wrapped on the homologically trivial S^2, similarly to some of the settings in [31]. We do not need to include branes in our setup, the main point being that the axion acquires its monodromic potential from the homological triviality of the S^2 (in contrast to models where the potential is due to the tension of the NS5-brane). Thus we do not need to include anti-branes either and therefore evade the dangerous brane/antibrane back-reaction described in [34, 35]. We note that our results might also be useful in the context of recently proposed Relaxion-models [43–52].

We find that the mass of the lightest 4d-Kaluza-Klein mode is lighter than the next heavier mode by a *relative* warp-factor which makes it an interesting candidate for single field inflation. Thus the inflaton potential is suppressed by warping [53] without the need for an additional tuning. Since this is due to the S^2 ending in the infrared-region we need a second throat into which the S^2 can bend around in the UV such that its second end lies in an infrared region as well. Such a geometry has been constructed in [54, 55] which we very briefly review in Section 2.

This paper is organized as follows: In Section 3 we calculate the IR-localized 5d-mass-term, finding that $M \sim 1/R$ where R is the typical radius of the KS-region. In Section 4, starting from the 5d-effective model we perform a Kaluza-Klein-reduction along the radial coordinate of the throat, thereby obtaining the effective 4d-theory with an infinite tower of KK-modes with the above mentioned mass-suppression of the lightest mode. In Section 5 we compare the energy-densities of the inflaton with those stabilizing the throat, concluding that the inflaton does not back-react strongly on the geometry in the weak coupling regime $g_s \ll 1$.

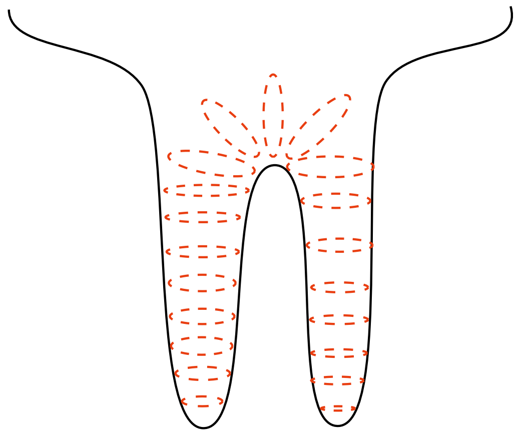

Figure 1: The Double-Throat: The dashed line indicates the family of S^2's bending around into the second throat, shrinking to zero size at the tips.

2 The Double-Throat

Let us briefly review the construction of the double-throat (see Figure 1) following the discussion in [55][2]. The conifold can be described as the subset of $\mathbb{C}^4$ solving

$$uv = y^2 + x^2 \quad , \quad (u, v, x, y) \in \mathbb{C}^4 \; . \tag{1}$$

The conifold singularity sits at $x = y = u = v = 0$. We can construct a two-conifold-setup by replacing x with a polynomial $W'(x)$ in the conifold equation (1). We take W' to have two simple roots at $x \in \{a_1, a_2\}$:

$$uv = y^2 + W'(x)^2 \quad \text{where} \quad W'(x) = g(x - a_1)(x - a_2) \; . \tag{2}$$

If $g = 0$ this gives a curve of A_1-singularities parametrized by x. Blowing up the singularity gives a curve of P^1's. Setting $g \neq 0$ there is still a family of S^2's related in homology. After a geometric transition [54] the system is deformed by means of a polynomial f_1 of degree one, to give two deformed conifolds with shrinking S^2:

$$uv = y^2 + W'(x)^2 + f_1(x) \; . \tag{3}$$

This is precisely the geometry we will use.

[2]This geometry can also be viewed as a Z_2-orbifold of the conifold [37].

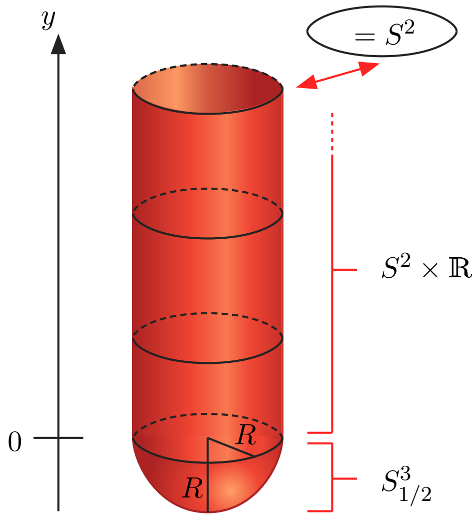

Figure 2: The geometry close to the tip of the throat.

3 A Simple Geometric Setup and Reduction to 5d

Consider the standard KS-throat with a blown up S^3_{KS} but trivial S^2-cycle. Due to the homological triviality of the S^2 there is no harmonic 2-form and thus no massless axion $c = \int_{S^2} C_2$. Our axion will hence be the (massive) lightest KK-mode of C_2.

As a first approximation, let us take the geometry of the compact space to be simply $M_6 = S^3_{KS} \times X_3$ where X_3 is a 'cylinder' ($= S^2 \times \mathbb{R}$) of constant radius R, which is closed by one half of a three-sphere ($\equiv S^3_{1/2}$) in the IR. In the UV the S_2 bends around into a second throat such that it is closed in the IR on both sides as depicted in Figure 1. This is crucial since we would otherwise generate a UV-mass-term.

Let y be the radial coordinate such that $y = 0$ at the boundary of $S^3_{1/2}$ (see Figure 2). Our starting point is the type-IIB supergravity action in Einstein frame (see [56], ch. 12.1):

$$S_{IIB} \supset \frac{1}{2\kappa_{10}^2}\left(\int d^{10}x\sqrt{-g^{10d}}R - \frac{g_s}{2}\int F_3 \wedge *F_3 - \frac{1}{2g_s}\int H_3 \wedge *H_3\right) , \qquad (4)$$

where $F_3 = dC_2$ and $H_3 = dB_2$ are the three-form field-strengths and we have restricted ourselves to constant dilaton $e^\phi \equiv g_s$ and vanishing C_0. We now expand:

$$C_2 = \phi(x,y)\,\omega_2 \quad , \; x \in \mathbb{R}^{1,3} \, , \qquad (5)$$

where we take ω_2 to be the canonical volume-form of S^2 (normalized to $\omega_2 = (Vol_{S^2})^{-1} *_2 1$).

We now want to derive the effective 5d-action which we will then treat as an effectively 5d Randall-Sundrum-model [57, 58] in Section 4 (see [59] for the 5d-description of the throat).

First we derive the bulk-term. Thus we plug the above into the 10d-(Einstein-frame)-action $S_{IIB} \supset -\frac{g_s}{4\kappa_{10}^2}\int dC_2 \wedge *dC_2$ and get a bulk kinetic term

$$-\frac{g_s}{4\kappa_{10}^2}\int_{M_5} d^5x\sqrt{-g^{5d}}\,\partial_A\phi\,\partial^A\phi\int_{S^3_{KS}} *_3 1\underbrace{\int_{S^2}\omega\wedge *_2\,\omega}_{=(Vol_{S^2})^{-1}\equiv a_1}\quad , A\in\{0,...,4\}\,. \tag{6}$$

Next let us calculate the contribution of the boundary $S^3_{1/2}$. Since S^2 is trivial (e.g. at $y=0$, $S^2=\partial S^3_{1/2}$) we have

$$\phi(x,y)=\int_{\partial S^3_{1/2}} C_2=\int_{S^3_{1/2}} F_3\,, \tag{7}$$

with $F_3 = dC_2$. Neglecting the warping, the lowest energy configuration is where the field-strength F_3 is equally distributed over $S^3_{1/2}$. Hence we make an ansatz

$$F_3=\gamma\,\omega_3\quad , \gamma\in\mathbb{R}\,, \tag{8}$$

where ω_3 is the canonical volume form of the three-sphere (i.e. $\omega_3 = (Vol_{S^3})^{-1} *_3 1$). It follows that

$$\phi(x,0)=\int_{S^3_{1/2}} F_3=\frac{1}{2}\int_{S^3} F_3=\frac{\gamma}{2}\quad\to\quad F_3=2\phi(x,0)\,\omega_3\,. \tag{9}$$

Plugging this into the 10d-action we get a boundary mass term

$$-\frac{g_s}{4\kappa_{10}^2}\int F_3\wedge *\,F_3=-\frac{g_s}{4\kappa_{10}^2}\int_{\mathbb{R}^{1,3}} d^4x\sqrt{-g^{5d}_{y=0}}\,4\,\phi^2(x,0)\int_{S^3_{KS}} *_3 1\underbrace{\int_{S^3_{1/2}}\omega_3\wedge *_3\,\omega_3}_{=\frac{1}{2}(Vol_{S_3})^{-1}\equiv a_2}\,, \tag{10}$$

where we have again neglected the effect of warping on $S^3_{1/2}$. Going over to a canonically normalized $5d$-field ($\frac{g_s}{2\kappa_{10}^2}a_1\,Vol_{S^3_{KS}}\,\phi^2\ \to\ \phi^2$) we get a 5d-action

$$S_5=\int\sqrt{-g^{5d}}\,\{-\frac{1}{2}\partial_A\phi\partial^A\phi-\frac{1}{2}M\delta(y)\,\phi^2\}\quad\text{with}\quad M=4\frac{a_2}{a_1}=\frac{4}{\pi R}\,. \tag{11}$$

Therefore the localized mass-term is essentially $M \sim R^{-1}$ where R is the typical transverse size of the throat which in this case coincides with the length-scale over which the throat contracts.

4 KK-Reduction on the Effective 5d-Throat and the 4d Action

The 5d action derived in the previous section can now be reduced to an effective 4d action containing an infinite tower of 4d-KK-modes. We now treat the throat as an effectively 5-dimensional Randall-Sundrum-model [57–59].

Consider the following 5d-metric [57, 58]:

$$ds^2 = e^{2ky}\eta_{\mu\nu}dx^\mu \otimes dx^\nu + dy^2 \; . \tag{12}$$

The 5d Lagrangian now reads

$$\mathcal{L}_5 = e^{4ky}\{-\frac{1}{2}e^{-2ky}\partial_\mu\phi\partial^\mu\phi - \frac{1}{2}(\partial_y\phi)^2 - \frac{M}{2}\delta(y)\phi^2\} \; , \tag{13}$$

where 4d indices are contracted using $\eta = \mathrm{diag}(-1, 1, 1, 1)$.

We now let y take values on a strip of length L choosing orbifold identification $y \cong -y$ and $y \cong y + 2L$ (note that in this case we need to double M, that is $M = 8/\pi R$, in order to get the physical boundary-condition for the 5d-field).

Inserting a 4d-plane-wave-ansatz $\phi(x, y) = e^{ipx}\chi(y)$ with $p^2 = -m^2$ the equations of motion take the form

$$-\partial_y(e^{4ky}\partial_y\chi) + Me^{4ky}\delta(y)\chi = e^{2ky}m^2\chi \; , \tag{14}$$

which can be brought into the form of a 1d Schrödinger equation [58] (we follow explicitly [60])

$$(-\partial_z^2 + V(z))f(z) = E\,f(z) \; , \tag{15}$$

with $z := e^{-ky}$, $\chi = z^{\frac{3}{2}}f$, $E = \frac{m^2}{k^2}$ and potential

$$V(z) = \frac{15}{4}\frac{1}{z^2} + \frac{3 + M/k}{z_{IR}}\delta(z - z_{IR}) - \frac{3}{z_{UV}}\delta(z - z_{UV}) \; , \tag{16}$$

where $z_{IR} = 1$ and $z_{UV} = e^{-kL}$.

Note that the delta-potentials come from enforcing the appropriate boundary conditions on χ (not on f). The general solution (a special case of the more general situation considered in [61]) now takes the form

$$\begin{aligned} f(z) &= \sqrt{z}\left(A\,J_2\left(\frac{m}{k}z\right) + B\,Y_2\left(\frac{m}{k}z\right)\right) \quad , \; A, B \in \mathbb{C} \\ \Rightarrow \quad \chi(y) &= e^{-2ky}\left(A\,J_2\left(\frac{m}{k}e^{-ky}\right) + B\,Y_2\left(\frac{m}{k}e^{-ky}\right)\right) \; , \end{aligned} \tag{17}$$

where J_n and Y_n are the Bessel functions of first and second kind respectively.

From the form of the potential we immediately deduce the existence of a single (UV-)

bound state and wave solutions of higher energy (mass) that are exponentially suppressed in the UV. Note that the bound state solution can be determined exactly in the case where $M = 0$:

$$f_0(z) = A\, z^{-\frac{3}{2}} + B\, z^{\frac{5}{2}} \;\Rightarrow\; \chi_0(z) = A + B\, z^4 \,, \tag{18}$$

which simplifies to $\chi = const.$ after imposing boundary conditions. This is of course the constant mode of zero mass which can be immediately read of from (14).

The mass-condition follows from the two boundary conditions ($\partial_y \chi(0) = \frac{M}{2}\chi(0)$ and $\partial_y \chi(e^{-kL}) = 0$) and reads

$$J_1\left(\frac{m}{k}\right) + \frac{M}{2m} J_2\left(\frac{m}{k}\right) - \frac{J_1(\frac{m}{k}e^{-kL})}{Y_1(\frac{m}{k}e^{-kL})}\left(Y_1\left(\frac{m}{k}\right) + \frac{M}{2m} Y_2\left(\frac{m}{k}\right)\right) = 0 \,. \tag{19}$$

We will now focus on the case $r_c := \frac{1}{k} \ll L$ (which is the interesting case of strong warping). For the bound-state solution we expect a small mass ($m \ll k$) for which we can use the small argument approximations of the Bessel-functions

$$\begin{aligned} J_1(x) &= \frac{x}{2} + \mathcal{O}(x^3) \quad & J_2(x) &= \frac{x^2}{8} + \mathcal{O}(x^4) \\ Y_1(x) &= -\frac{2}{\pi x} + \mathcal{O}(x) \quad & Y_2(x) &= -\frac{4}{\pi x^2} + \mathcal{O}(x^0) \,, \end{aligned} \tag{20}$$

to arrive at

$$\frac{m_0}{k} = \left(\frac{k}{M} + \frac{1}{8}\right)^{-\frac{1}{2}} e^{-kL} \,. \tag{21}$$

Remarkably this mass is exponentially suppressed by the warp factor (thereby a posteriori justifying our small argument approximation). It is crucial to realize that this is not the usual hierarchy induced by warping in Randall-Sundrum models [57] but is rather a suppression 'on top of that' since our metric conventions are such that $g^{IR}_{\mu\nu} \equiv g_{\mu\nu}(y = 0) = \eta_{\mu\nu}$.

The zero-mode profile takes the following form:

$$\chi_0(y) \propto \left(1 - \frac{1}{8}\frac{1}{k/M + 1/8} e^{-4ky}\right) \,. \tag{22}$$

The higher KK-modes (with $1 \lesssim \frac{m}{k} \ll e^{kL}$) are obtained by noting that $J_1/Y_1(x) = -\frac{\pi}{4}x^2 + \mathcal{O}(x^4)$ such that the mass condition is approximately

$$J_1\left(\frac{m}{k}\right) + \frac{M}{2m} J_2\left(\frac{m}{k}\right) = 0 \,. \tag{23}$$

The solutions interpolate between the zeros of the two Bessel-functions ($j_{1,n}$ and $j_{2,n}$), that

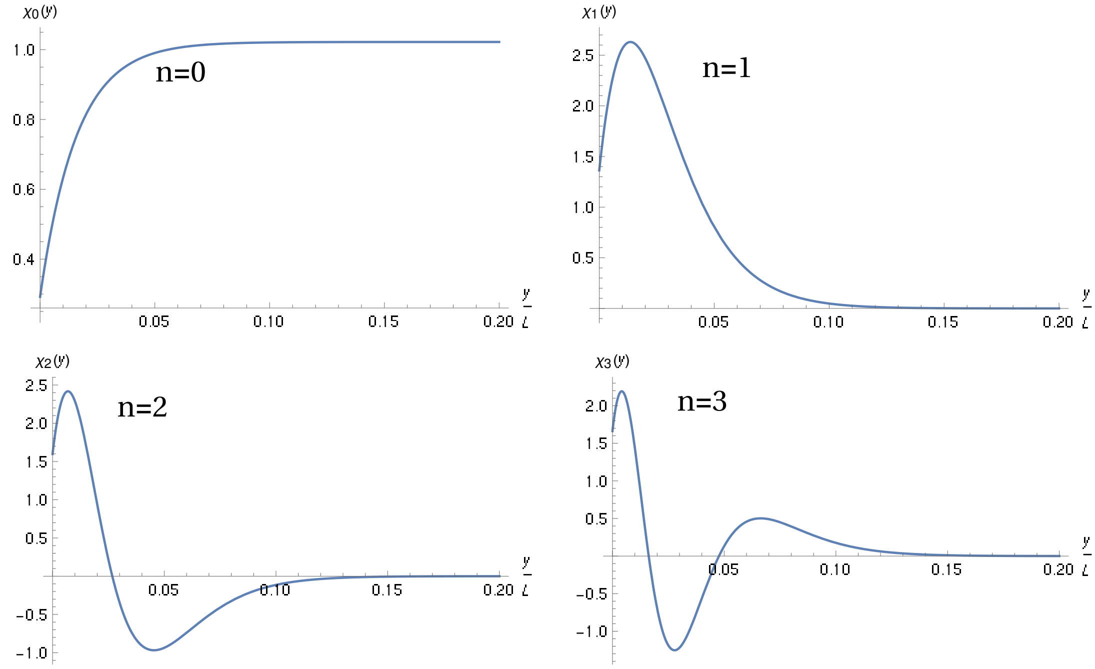

Figure 3: The bound state and first excited solutions with $kL = 5\pi$ and $M = 100\pi/L$ ($r_c/L \approx 0.06$), plotted in the range $0 \leq y \leq 0.2\,L$

is

$$\begin{aligned} m \ll M: &\quad \frac{m_n}{k} \approx j_{2,n} \\ m \gg M: &\quad \frac{m_n}{k} \approx j_{1,n} \end{aligned} \tag{24}$$

and asymptotically (that is $m_n \gg k, M$)

$$m_n = \pi \frac{n}{r_c} \,, \tag{25}$$

which are the usual KK-masses but with L replaced by the curvature radius $r_c \equiv k^{-1}$. The bound-state and the first excited states are plotted in Figure 3.

Using the 4d-Planck mass [57]

$$M_{pl}^2 = \frac{M_5^3}{k}\left(e^{2kL} - 1\right) \overset{kL \gg 1}{\approx} \frac{M_5^3}{k} e^{2kL} \,, \tag{26}$$

one immediately sees the double exponential suppression of the bound-mode:

$$\begin{aligned} \frac{m_0^2}{M_{pl}^2} &\approx \frac{k^3}{M_5^3}(k/M + 1/8)^{-1} e^{-4kL} \propto e^{-4kL} \\ \frac{m_n^2}{M_{pl}^2} &\approx \frac{k^3}{M_5^3}\pi^2 n^2 e^{-2kL} \propto e^{-2kL} \quad \forall\, n \neq 0 \,, \end{aligned} \tag{27}$$

Note that this agrees with the expression for the axion-potential in equation (4.76) of [18] where the potential comes from the NS5-DBI-action.

This behavior could have already been anticipated from the form of the potential (16): The bound-state-solution approaches a constant in the UV while the positive delta-potential in the IR leads to a dip in the IR. It therefore gets its mass from the IR while its kinetic term lives in the whole bulk (concerning the kinetic term arguments along these lines have already been given in [18], Sec. 4.3.2). This leads to the already mentioned 'double'-suppression. The higher KK-modes are the solutions to Schrödinger's equation (15) that oscillate in the IR-region $0 < y \lesssim r_c$ and fall off exponentially towards the UV due to the $\sim 1/z^2$-term in the potential (16). This leads to the modified KK-mass-formula (25).

Note furthermore that m_0 (more precisely its upper bound) is not particularly sensitive to the value of M:

$$0 \leq m_0 \leq \sqrt{8} k e^{-kL} \quad \Rightarrow \quad 0 \leq \frac{m_0}{m_1} \leq (j_{1,1})^{-1} \sqrt{8}\, e^{-kL} \,, \tag{28}$$

where $j_{1,1} \approx 3.8317...$ is the first zero of J_1. The 4d-effective action is

$$S = \int_{\mathbb{R}^{1,3}} d^4x \sum_{n=0}^{\infty} \{-\frac{1}{2}\partial_\mu \phi_n \partial^\mu \phi_n - \frac{1}{2} m_n^2 \phi_n^2\} \,. \tag{29}$$

Let us pause here and highlight what we have found:

The lightest KK-mode of the RR-2-form C_2 on the KS-throat with trivial S^2-cycle is exponentially lighter than the next higher mode in the case of strong warping. This makes it an ideal candidate for single-field chaotic inflation since we can safely ignore the higher modes.

5 Simple Consistency-Checks from Energetics and Mass-Scales

5.1 Energy-Density at the boundary of $S^3_{1/2}$

It is important to check that the energy-density at $y = 0$ on the cylinder is the same (at least up to $\mathcal{O}(1)$-factors) as the one on $S^3_{1/2}$.

On the cylinder we have

$$S_{cyl} = -\int d^5x \int_{S^3_{KS}} *_3 1 \int_{S^2} *_2 1 \sqrt{-g^{5d}} \underbrace{\frac{g_s}{4\kappa_{10}^2} (Vol_{S^2})^{-2} (\partial_y \phi)^2}_{:=\varepsilon_{cyl}} \,, \tag{30}$$

while on $S^3_{1/2}$ we have

$$S_{S^3_{1/2}} = -\int_{\mathbb{R}^{1,3}} d^4x \int_{S^3_{KS}} *_3 1 \int_{S^3_{1/2}} \sqrt{g^{S^3_{1/2}}}\sqrt{-g^{4d}}\, \underbrace{\frac{g_s}{4\kappa_{10}^2}(Vol_{S^3})^{-2}4(\phi)^2}_{:=\varepsilon_{S^3_{1/2}}} , \tag{31}$$

which implies that

$$\frac{\varepsilon_{cyl}}{\varepsilon_{S^3_{1/2}}} = \frac{1}{4}\underbrace{\left(\frac{\partial_y\phi(0)}{\phi(0)}\right)^2}_{\equiv(M/2)^2}\underbrace{\left(\frac{Vol_{S^3}}{Vol_{S^2}}\right)^2}_{\equiv(\frac{\pi}{2}R)^2} = 1\ . \tag{32}$$

Therefore the energy-densities are exactly the same. Note that this were also true if we had chosen any other eigen-mode of the 5d-Laplacian since the identity $M = \partial_y\chi_n(0)/\chi(0)$ is simply the boundary condition for the y-profile of *any* mode χ_n.

5.2 Inflaton Energy-Density vs F_3-Flux-Energy-Density

Since we have a model of single-field large field inflation we have to make sure that the field excursion of the inflaton does not back-react strongly on the geometry thereby destroying the throat.

The Flux-Energy-Density can be calculated from the Type IIB Supergravity action

$$S_{RR} \supset -\frac{g_s}{4\kappa_{10}^2}\int F_3\wedge *F_3 \quad \text{with} \quad F_3 = (2\pi)^2\alpha'\tilde{M}\omega_3 + ...\ , \tag{33}$$

where ω_3 is the appropriately normalized volume form on S^3_{KS} and $\tilde{M}$ is the F_3-flux on S^3_{KS} stabilizing the throat. Ellipsis indicate terms that integrate to zero over S^3_{KS}.

Using $\kappa_{10}^2 = \frac{l_s^8}{4\pi}$ (where $l_s = 2\pi\sqrt{\alpha'}$) and $\int \omega_3\wedge *\omega_3 = (Vol_{S^3})^{-2}\int d^{10}x\sqrt{-G^{10d}}$ this yields the local (10d)-energy-density

$$\varepsilon^{10d}_{KS} = \frac{g_s\tilde{M}^2}{4\pi^3R^6l_s^4} = \frac{2^4\pi^3}{l_s^{10}}(\tilde{M}g_s^2)^{-1}\ , \tag{34}$$

where R is the radius of the S^3_{KS} (which we identify with the $S^3_{1/2}$ radius). In the second step we have used that $R^2 = \tilde{M}g_s(l_s/2\pi)^2$ in the KS-region (see equ. (93) in [42]).

The inflaton energy-density (using equations (10),(11) and the explicit form of the bound mode 22) is given by

$$\begin{aligned}\varepsilon^{10d}_\phi &= \frac{g_s\phi^2}{\kappa_{10}^2(Vol_{S^3})^2} = \frac{2\phi^2_{5d-can.}(Vol_{S^2})}{(Vol_{S^3})^3} = \frac{2\alpha^2}{\pi^5}\frac{M_5^3}{R^7}\frac{k/M}{k/M+1/8} \\ &\leq \frac{2\alpha^2}{\pi^5}\frac{M_5^3}{R^7} = \frac{2^6}{\pi}\frac{1}{R^2l_s^8} = \frac{2^8\pi}{l_s^{10}}\alpha^2(\tilde{M}g_s)^{-1}\ ,\end{aligned} \tag{35}$$

where α measures the 4d field excursion in 4d-Planck units (equivalently the 5d excursion in 5d-Planck units). The ratio of the densities therefore satisfies

$$\varepsilon_\phi/\varepsilon_{KS} \leq \frac{16}{\pi^2}\alpha^2 g_s \sim \alpha^2 g_s \ . \tag{36}$$

Therefore if $g_s \ll 1$ we do not expect the inflaton to back-react strongly on the KS-geometry for field excursions up to $g_s^{-1/2}$.

5.3 The H_3-Energy-Density

Next we will check the energy-density of the Kalb-Ramond-field B_2 and compare it to the inflaton contribution. The action contains a term

$$S \supset -\frac{1}{4g_s\kappa_{10}^2}\int dB_2 \wedge *dB_2 \ . \tag{37}$$

As shown in [62] the B_2 field has a radial dependence parametrized by $N_{eff}(y)$:

$$l_s^4 N_{eff}(y) \sim \int_{S^3_{KS}} F_3 \int_{S^2} B_2 = l_s^2 \tilde{M} \int_{S^2\, at\, y} B_2 \ , \tag{38}$$

where

$$\frac{dN_{eff}}{dy} = \frac{3}{2\pi}\frac{g_s \tilde{M}^2}{y} \ . \tag{39}$$

With the ansatz $B_2 = f(y)\,\omega_2$ this implies that $f(y) = l_s^2 N_{eff}/\tilde{M}$ and hence

$$\begin{aligned} -\frac{1}{4g_s\kappa_{10}^2}\int dB_2 \wedge *dB_2 &= -\frac{1}{4g_s\kappa_{10}^2}\int d^{10}x\sqrt{G}\left(\frac{f'(y)}{Vol_{S^2}}\right)^2 \\ &= -\int d^{10}x\sqrt{G}\frac{1}{4g_s\kappa_{10}^2}\left(\frac{3}{2\pi}\frac{l_s^2\tilde{M}g_s}{Vol_{S^2}y}\right)^2 \ . \end{aligned} \tag{40}$$

Evaluated at $y = R$ this gives

$$\varepsilon_{B_2} = \frac{9\pi^3}{l_s^{10}}(g_s^2\tilde{M})^{-1} = \frac{9}{16}\varepsilon_{KS} \ . \tag{41}$$

Thus $\varepsilon_\phi \ll \varepsilon_{B_2}$ once we have achieved $\varepsilon_\phi \ll \varepsilon_{KS}$.

6 Conclusion

In this paper we presented a new model for axion monodromy inflation in which the inflaton is the lightest Kaluza-Klein mode of the RR-2-form potential C_2 wrapped on a homologically trivial 2-cycle. One of the crucial technical points is that the mass of the lightest Kaluza-Klein mode is exponentially lower than that of the next excited mode, thus making this mode an ideal inflaton candidate. The monodromy arises due to the homological triviality of the 2-cycle similar to models proposed in [31], rather than due to a coupling to branes. Consequently, our construction does not require the presence of brane-antibrane pairs, thus avoiding the associated back-reaction issues [34, 35]. Crucially, the exponential mass-suppression is due to the S^2 shrinking to zero size only in IR regions. This is why we base our model on the 'double-throat' shown in Figure 1. Furthermore, by comparing the energy density in the inflationary sector with the contribution due to fluxes, we find that the inflaton sector does not back-react substantially on the geometry in the weak coupling regime $g_s \ll 1$.

Of course, there are many questions which we left unanswered. For example, any further violation of the shift symmetry of C_2 might introduce potentially dangerous corrections to the inflaton potential. It is thus of great importance to determine to what extent the shift-symmetry of C_2 is preserved in the warped background. Further, while we showed that back-reaction on the geometry can be controlled, it would be desirable to include the effects of back-reaction onto the inflaton potential explicitly. The expectation is that this would lead to a flattening of the inflaton potential as observed in [30, 63]. This flattening might be beneficial for the phenomenology of our construction. In its current form, our model results in a quadratic inflaton potential which is not compatible with experimental data (see e.g. [27, eq. (2.51)]). This situation would improve for a flatter inflaton potential.

Overall, we observe that our proposal realizes axion monodromy inflation inflation for a fairly minimal amount of ingredients. Given this relative simplicity and the high level of sophistication with which throat geometries can be controlled [64, 65], we expect our model to be a promising arena for further investigations into the viability of large field inflation in string theory. Regardless of the phenomenological implications, we would even like to hope that, as a matter of principle, the possibility of large field inflation could be firmly established based on our simple scenario.

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