Towards Universal Axion Inflation and Reheating in String Theory

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Abstract

The recent BICEP2 measurements of B-modes indicate a large tensor-to-scalar ratio in inflationary cosmology, which points towards trans-Planckian evolution of the inflaton. We propose possible string-theory realizations thereof. Schemes for natural and axion monodromy inflation are presented in the framework of the type IIB large volume scenario. The inflaton in both cases is given by the universal axion and its potential is generated by F-terms. Our models are shown to feature a natural mechanism for inflaton decay into predominantly Standard Model particles.

1. Introduction

The recent reports of the BICEP2 collaboration\textsuperscript{1} indicate that for the first time there have been direct measurements of B-modes, which are CMB imprints of primordial gravity waves. Indeed, BICEP2 found the tensor-to-scalar ratio to be rather large, \( r = 0.2 \). Realizing this value in slow-roll inflationary cosmology requires a motion of the inflaton over trans-Planckian distances in field space. This is difficult to achieve in a UV complete theory of quantum gravity, such as string theory, as one has to have control over many possible higher-order operators, which can spoil the slow-roll property. This is known as the \( \eta \)-problem, which is a challenge in general and becomes even stronger for trans-Planckian evolutions.

As a matter of fact, most string or string-inspired models of inflation give much lower values of \( r \) and would, if BICEP2 is confirmed, be ruled out. In order to achieve the required control over higher-order corrections, one can take advantage of the perturbative shift symmetry of axions, which is only broken non-perturbatively, by fluxes, or by the presence of branes. Interestingly, axions are ubiquitously present in string compactifications. In the prototype model, the inflationary potential takes the form

\[ V(\theta) = \Lambda^4(1 - \cos(\theta/f)), \]

where \( f \) denotes the instanton decay constant of the axion \( \theta \). In fact, this simple model is known to lead to a large tensor-to-scalar ratio, consistent with all other cosmological parameters, if \( f > M_{\text{pl}} \). It is a priori unclear whether in this regime the effective field-theory description is still trustable, however, this model of natural inflation\textsuperscript{2, 3} fits the data amazingly well. Expanding the potential to quadratic order, it essentially reproduces Linde’s model of chaotic inflation\textsuperscript{4}, which in fact has been shown to be compatible with large field variations in\textsuperscript{7, 8, 9}.

In order to avoid the regime \( f > M_{\text{pl}} \), extensions of this simple axion model have been proposed. One is N-flation\textsuperscript{8, 9}, where the collective evolution of \( N \) axions in a sort of radial direction leads to the same predictions, but where each individual axion only travels over a sub-Planckian distance and has a decay constant \( \sqrt{\mathcal{N}} f_i > M_{\text{pl}} \). A second variation is so-called axion monodromy inflation\textsuperscript{10–12}, where the shift symmetry is broken by the presence of a brane, inducing an approximately linear (or quadratic) potential for the axion. Thus, the former periodic axion is “unwrapped” and can now move over trans-Planckian distances, increasing the energy by a certain amount each time it goes around the period.

It turns out that realizing N-flation in a concrete string-theory model is not an easy task\textsuperscript{13–14}. One generic problem with N-flation is that for \( N > 1 \) a substantial renormalization of the Planck mass occurs. Another potential problem of many models is that at the end of the inflationary epoch, the inflaton is not guaranteed to predominantly decay into the visible sector, but the decay rates into a visible and a hidden sector degree of freedom tend to be of the same order (see e.g.\textsuperscript{15, 16}).

Clearly, in view of the BICEP2 results and its fairly constraining consequences, it is important to study what possibilities string theory can offer to realize such axion-inflation models. In this letter, we ignore the difficulties the regime \( f > M_{\text{pl}} \) causes and construct two models: a string-theory model of natural inflation, and a new type of monodromy inflation, where the shift symmetry of the axion is broken by fluxes instead of by the presence of branes. For concreteness, we focus on the LVS framework\textsuperscript{17} and consider as the inflaton (a linear combination involving) the axionic component of the complex axio-dilaton field

\[ S = C_0 + i \exp(-\phi). \]

Note that in most approaches before, this universal axion was considered to be fixed at high mass scale by background three-form fluxes.

We argue that type IIB string theory has the necessary ingredients to construct successful models of inflation. In
particular,

- We assume that the (flux) landscape admits points where the masses of the saxions (including the dilaton) are hierarchically different from the mass of $C_0$. In particular, apart from the nearly massless axion of the big four-cycle in a LVS, $C_0$ can be the lightest closed-string modulus, making it a good candidate for the inflaton.

- For natural inflation, the potential of the axion is generated by non-perturbative effects from fluxed E3-instantons, whereas for axion monodromy inflation the axion $C_0$ can appear quadratically in the flux induced scalar potential.

- There exists a mechanism guaranteeing that inflaton decay at the end of inflation predominantly goes into standard model (SM) degrees of freedom.

This last point is one of the very interesting aspects of the models considered in this letter. Note furthermore that the relevant axion potentials are F-terms in an effective field theory approach. Recall some results for the cosmological parameters derived from the simple natural-inflation Lagrangian

\[ \mathcal{L} = \frac{1}{2} \partial_\mu \theta \partial^\mu \theta + \Lambda^4 \left( 1 - \cos \left( \frac{\theta}{7} \right) \right), \]

where after canonical normalization the axion $\theta$ has a period $2\pi f$. The slow-roll parameters

\[ \epsilon = \frac{M^2_{\text{pl}}}{2} \left( \frac{V''}{V} \right)^2, \quad \eta = M^2_{\text{pl}} \left( \frac{V'''}{V''} \right), \]

can be expanded in $\theta/f \ll 1$ as

\[ \epsilon = \eta \approx \frac{2M^2_{\text{pl}}}{f^2} \left( \frac{f}{\theta} \right)^2. \]

For this to be small one needs $M_{\text{pl}} < \theta < f$ during inflation, and inflation ends at $\theta_{\text{end}} = \sqrt{2}M_{\text{pl}}$. The number of e-foldings are expressed as follows

\[ N_e = \frac{1}{M^2_{\text{pl}}} \int_{\theta_{\text{end}}}^{0} \frac{V}{\sqrt{V'}} d\theta \approx \frac{1}{2M^2_{\text{pl}}} \int_{\theta_{\text{end}}}^{0} \theta d\theta = \frac{\theta^2}{4M^2_{\text{pl}}} - \frac{1}{2}. \]

Thus, we write $N_e \sim \frac{1}{\theta f^2}$. Therefore, for the spectral indices and the tensor-to-scalar ratio one obtains the prediction

\[ n_s = 1 + 2\eta - 6\epsilon \sim 1 - 4\epsilon \sim 1 - \frac{2}{N_e}, \]

\[ n_t = -2\epsilon \sim - \frac{1}{N_e}, \quad r = 16\epsilon \sim \frac{8}{N_e}. \]

For 60 e-foldings, this rough estimate gives

\[ n_s \sim 0.967, \quad n_t \sim -0.017, \quad r = 0.133, \]

which is in good agreement with the recent measurements from Planck for $n_s = 0.9624 \pm 0.0075$ and from BICEP2 for $r = 0.20^{+0.07}_{-0.05}$. The amplitude of the scalar power spectrum

\[ \mathcal{P} = 2.2 \cdot 10^{-9} \]

leading to a Hubble constant during inflation of $H_{\text{inf}} \sim 1.1 \cdot 10^{14}$ GeV. Via $V_{\text{inf}} = 3M^2_{\text{pl}}H^2_{\text{inf}}$ we can now extract the mass scale of inflation as $\frac{M^2_{\text{inf}}}{V_{\text{inf}}} = V_{\text{inf}}^{1/4} \sim 2.1 \cdot 10^{16}$ GeV, which is of the order of the GUT scale. Finally, the mass of the axion $m^2_{\theta} = 3\eta H^2_{\text{inf}}$ comes out as $m_{\theta} \sim 1.7 \cdot 10^{13}$ GeV.

Note that since only the quadratic approximation of the cosine function was used, these predictions are the same as for a quadratic potential, which is nothing else than Linde’s model of chaotic inflation. The main problem to realize such a model in string theory is the constraint $f > M_{\text{pl}}$, which for all known cases means that one is outside of the regime of validity of the effective field theory approach. To overcome this restriction, it was proposed to consider axions and break their shift symmetry by branes or fluxes, generating a non-oscillatory potential for the axion and thus allowing it to roll over trans-Planckian distances in field space, without needing $\theta/f < 1$. Hence, the periodic axion unfolds to allow for non-trivial monodromies.

2. Natural inflation

Before we present our string-theory realization, let us recall some results for the cosmological parameters derived from the simple natural-inflation Lagrangian

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3. Universal natural inflation in String theory

Let us present a possible string-theory realization of the above scenario, based on the well-established framework of moduli stabilization in type IIB orientifolds on Calabi-Yau three-folds with 07- and 03-planes. Recall that in the standard LVS one starts with a swiss-cheese type Calabi-Yau three-fold with Kähler potential

\[ K = K_{\text{cal}}(U) - \log(-i(S - \overline{S})) - 2 \log \left( (T_b + \overline{T}_b)^2 - (T_s + \overline{T}_s)^2 + \xi \left( \frac{S + \overline{S}}{2} \right)^2 \right), \]

and fixes the axio-dilaton and complex structure moduli via a tree-level three-form flux-induced superpotential $W = \int \Omega \wedge (F - iSH)$. For $\xi = 0$, the no-scale structure leaves the Kähler moduli $T_b = \tau_b + i\alpha_b, T_s = \tau_s + i\alpha_s$
massless at this level. Note that the masses of the complex structure moduli and the axio-dilaton are generically of the order $m_{\alpha s} \simeq 1/\sqrt{2}$, which is heavier than the Kähler moduli. Therefore, for constructing a model with the universal axion as the inflaton, we need to keep the axion essentially massless at the level of the flux-induced potential.

Now, due to the landscape of flux vacua, there could exist special but still numerous non-supersymmetric vacua where the universal axion remains unfixed with its shift symmetry still intact. Therefore, the superpotential at the minimum can be independent of $C_0$. For this letter, we leave aside the technical issue of stabilizing e.g. the dilaton in such a way that its mass is hierarchical bigger than the mass of the axion. Thus, we assume that the complex structure moduli and the dilaton are fixed either directly by fluxes, or by additional no-scale violating terms, such that the above assumption is satisfied. In this case the value of the superpotential $V_0$ at the minimum is a constant. Let us emphasize that a concrete realization of this moduli stabilization scheme might be technically challenging and deserves a closer investigation [19].

In the LVS one breaks the no-scale structure by a combination of $\alpha'$-corrections leading to $\xi \neq 0$ and an $E3$-instanton induced superpotential

$$W = W_0 + A_s e^{-\alpha_s T_s},$$

where $A_s = O(1)$ denotes the one-loop Pfaffian after integrating out the complex structure moduli and $\alpha_s$ is a numerical factor. Note that $A_s$ is independent of the axio-dilaton $S$. For such an instanton to directly contribute to $W$ it has to have the right number of fermionic zero modes. First, it has to be invariant of type $O(1)$ under the orientifold projection. Second, sufficient conditions are that it carries no deformation or Wilson line modulini and does not intersect any of the potentially present space-time filling D7-branes. Del-Pezzo surfaces are good examples.

The main observation of [17] is that the scalar potential for the Kähler-moduli sector

$$V = e^X \left( \sum_a K^{ab} D_a W D_b W - 3|W|^2 \right),$$

(10)

can be expanded in large volumes $V$ of the compactification space. With the superpotential [19], one then finds

$$V \simeq \sqrt{\tau_s} \alpha_s |A_s|^2 e^{-2\alpha_s \tau_s} \frac{V}{\sqrt{2}} + \tau_s \alpha_s |A_s| W_0 e^{-\alpha_s \tau_s} \frac{|W_0|^2}{g_s^2/2} + \xi \frac{|W_0|^2}{g_s^{3/2} \lambda^3},$$

(11)

admitting a non-supersymmetric AdS minimum in which the moduli are stabilized hierarchically as $\tau_b \gg \tau_s$. Concretely one finds

$$\tau_s \simeq \frac{O(1)}{g_s}, \quad \alpha_s \simeq O(1), \quad \nu \sim \frac{O(1)}{g_s^{1/2}} e^{\alpha_s \tau_s}$$

(12)

and the masses of the moduli around this minimum scale with the overall volume as

$$m_{\tau_s} \sim \frac{M_{pl}}{\sqrt{2}}, \quad m_{\alpha_s} \sim 0, \quad m_{\nu} \sim \frac{M_{pl}}{\sqrt{V}}.$$ 

(13)

From the second line it follows that the small-cycle axion $a_s$ is not a good candidate for axion inflation, as its mass is of the same order as $m_{\tau_s}$ and actually heavier than $m_{\tau_s}$. The big cycle axion $a_b$, after assuming it can get a mass of order $m_{ab} \sim \exp(-\sqrt{2}V)$ via a non-perturbative effect, has been considered for axion inflation in [14]. To avoid the problem with $f > M_{pl}$, more such big cycle axions were introduced and a model of N-flation was successfully built. However, let us mention that from the string-instanton zero-mode perspective it is questionable whether such a non-perturbative effect exist, as the big cycles tend to come with many deformation and Wilson line modulini.

From now on, we investigate the axion $\theta := C_0$ a bit closer. First, computing the axion decay constant from the Kähler potential we find from the log$(S + 3)$ term in [9] that

$$f = \frac{M_{pl}}{\sqrt{2}} g_s,$$ 

(14)

where the term in the second line of [9] gives only a sub-leading correction. Therefore, for $f > M_{pl}$ one needs $g_s > 1$, which actually is outside the perturbative regime of string theory. We will come back to this delicate issue later.

At this stage the universal axion $\theta$ is still massless as it does not appear in the superpotential. However, in addition to standard $E3$-brane instantons also magnetized $E3$-brane instantons generically make contributions to $W$ (see e.g. [20]). Say as in [19], we have an unfluxed $E3$-brane instanton wrapping the small del-Pezzo divisor $D_s$. Now, the del-Pezzo surface $D_s$ has $h^{11}_s = r + 1$ holomorphic two-cycles. Under the orientifold involutions these split into $h^{11}_s$ even and $h^{11}_s$ odd ones. Turning on a background gauge flux $f$ along the odd cycles still preserves the $O(1)$ property, and does not introduce modulini. Therefore, such a configuration also makes a contribution to the superpotential, so that one gets

$$W = W_0 + A_s e^{-\alpha_s T_s} + B_s e^{-\alpha_s (T_s + Sh(f))} + \ldots,$$ 

(15)

where the dots indicate that there can be many such contributions, and where $h(f) = \frac{1}{2\pi} \int_{D_s} f \wedge f \in \mathbb{Z}$. If $h(f)$ is big enough, this second term does not affect the moduli stabilization of e.g. $T_s$, as it is by a factor

$$\lambda_f = \exp \left( -\frac{\alpha_s h(f)}{2g_s} \right)$$

(16)

smaller than the second term in [15].
Next, we want to compute the masses and the mixing of the two axions $\phi^{(1)} = a_\tau$ and $\phi^{(2)} = \theta$. For this purpose we have to diagonalize the action

$$L_{ax} = K_{ab} \partial_a \phi^a \partial_b \phi^b - V_{ab} \phi^a \phi^b$$

for the fluctuations of these axions around their values at the LVS minimum of the scalar potential. Focusing on the order $V$ and $\lambda_f$ dependence, we find

$$V_{ab} = \begin{pmatrix} O(1) & O(1) \lambda_f \sqrt{V} \\ O(1) & O(1) \lambda_f \sqrt{V} \end{pmatrix}$$

and

$$K_{ab} = \begin{pmatrix} O(1) & O(1) \\ O(1) & O(1) \end{pmatrix} .$$

Computing the eigenvalues of the matrix $(K^{-1})^{ab} V_{mb}$ we obtain

$$M_1 \sim \frac{M_{pl}}{\sqrt{V}}, \quad M_2 \sim \frac{M_{pl} \sqrt{\lambda_f}}{\sqrt{V^{3/2}}} .$$

Diagonalizing then the axion sector, we find the relation between the old axions and the new eigensystem as

$$\phi^{(1)} \simeq O(1) V^{1/2} \psi^{(1)} + O(1) \lambda_f \psi^{(2)} ,$$

$$\phi^{(2)} \simeq O(1) \frac{\lambda_f}{V^{1/2}} \psi^{(1)} + O(1) \psi^{(2)} .$$

Therefore, $\phi^{(1)} = a_\tau$ is mostly $\psi^{(1)}$ and $\phi^{(2)} = \theta$ is mostly $\psi^{(2)}$. These relations will become important in the discussion of reheating in section 8.

Before closing this section, let us comment on the validity of the regime $f > M_{pl}$. We have seen that $f > M_{pl}$ leads to $g_s > 1$, so that one is actually outside the perturbative regime where the effective field theory was computed for. This clearly is the most severe issue of the natural inflation model presented in this letter. Let us make two comments:

- The perturbative axionic couplings are protected by a shift symmetry, and hence the main concern is about the description prior to the stabilization of the saxionic directions.

- The actual order parameter in the LVS is $V^{-1}$, so it is not completely obvious that for say $1 < g_s < 10$ we immediately get unstable results.

Concerning the last point, let us have a look at the corrections to the scalar potential originating from one-loop corrections to the Kähler potential. In [22] an extended no-scale structure was observed, for which the corrections to the scalar potential become

$$V = V_{LVS} + V_{1-loop}$$

$$\simeq \sqrt{\tau_s} e^{-2a_\tau \tau_s} \frac{V^2}{V} - \tau_s e^{-a_\tau \tau_s} + \frac{1}{g_s^{1/2} \sqrt{V}}$$

$$+ \frac{g_s^2}{V^3 / \sqrt{V_s} \sqrt{\theta}} + \frac{g_s^2}{V^3 / \sqrt{\theta}} .$$

In the LVS minimum these terms scale as

$$V \simeq \frac{1}{g_s^3 \sqrt{V}} + \frac{g_s^2}{V^{3/2}} + \frac{g_s^2}{V} \sqrt{\log \left( \frac{g_s^{1/2} \sqrt{V}}{V} \right)} .$$

Therefore, higher orders in $g_s$ are accompanied by further volume suppression factors. Of course, we cannot control all higher loop-corrections, but as long as they still come with suppressions in $V$, the LVS scalar potential could still be trusted for $g_s > 1$ and $V$ sufficiently large.

4. Universal axion monodromy inflation

One approach to avoid the regime $f > M_{pl}$ is to consider models where the axion shift symmetry is broken either by D-branes or by fluxes. Here, we want to construct such a model for the universal axion, whose shift symmetry is broken by the three-form flux. In this case a non-oscillatory potential is generated, allowing the axion to roll over trans-Planckian distances in field space without needing $\theta / f < 1$. Thus, the periodic axion unwraps to allow for non-trivial monodromies.

We consider a model in the flux landscape, where the fluxes break the shift symmetry of the axion slightly by giving it a parametrically small mass

$$m_\theta \simeq \frac{\lambda_0 M_{pl}}{V} .$$

Again, we assume that the complex-structure moduli and the dilaton can be fixed at a hierarchically bigger mass scale by a combination of fluxes and contributions to the scalar potential violating its no-scale structure [19]. In this case, the value of the superpotential in the minimum will have the following simple form

$$W_0 = w_0 + \lambda_0 \theta ,$$

guaranteeing that for the axion becoming massless the shift symmetry of the superpotential is restored. For the mass of axion to be smaller than the masses of Kähler moduli we require $\lambda_0 \ll \sqrt{V_0 / 2}$. The effective potential for the axion can then be written as

$$V_{eff} \simeq \frac{\lambda_0^2}{V^2} \theta^2 .$$

Interestingly, due to the flux-breaking of the axionic shift symmetry we eventually get an effective axion potential which is of quadratic order and therefore can be considered a candidate for a stringy realization of chaotic inflation. Moreover, this model is of the type of axion monodromy inflation, as the compact interval $\theta = \theta + 2 \pi f$ gets unwrapped.

For computing the masses and the mixing of the two axions $\phi^{(1)} = a_\tau$ and $\phi^{(2)} = \theta$, we now have the mass matrix

$$V_{ab} = \begin{pmatrix} O(1) & 0 \\ 0 & O(1) \lambda_f \sqrt{V} \end{pmatrix} .$$
For the eigenvalues of the matrix $\mathcal{K}^{-1}_{ab} V^m b$ we find
\[
M_1 \simeq \frac{M_{pl}}{\sqrt{\lambda}}, \quad M_2 \simeq \frac{M_{pl} \lambda_0}{\sqrt{\lambda}},
\]
with eigenvectors
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\phi^{(1)} \simeq O(1) V^a \psi^{(1)} + O(1) \frac{\lambda_0^2}{\sqrt{\lambda}} \psi^{(2)}, \quad \phi^{(2)} \simeq O(1) \frac{\lambda_0^2}{\sqrt{\lambda}} \psi^{(1)} + O(1) \psi^{(2)}.
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Therefore, as for universal natural inflation, $\phi^{(1)} = a_s$ is mostly $\psi^{(1)}$ and $\phi^{(2)} = \theta$ is mostly $\psi^{(2)}$.

It is important to know how stable this model is in the UV complete theory, i.e. whether it suffers from an $\eta$-problem due to other corrections depending on the inflaton. Shift symmetry breaking effects should either be proportional to the fluxes or come from non-perturbative effects like $E3$-brane instantons. The effect of the latter should be harmless as they are exponentially suppressed relative to the leading instanton contribution to $W$. In string theory, higher-order flux induced corrections to the Kähler potential are not well understood so that their analysis is beyond the scope of this letter. One should keep in mind that they could be potentially dangerous.

5. Reheating

At the end of inflation the axion $\theta$ oscillates around its minimum and decays into the various modes it couples to, thereby reheating the universe. It is often a difficult problem to guarantee that the inflaton mostly decays into the visible Standard Model sector $[15, 16]$. Too many decays into the hidden sectors produces dark matter and can over-close the universe. In order to estimate this, we consider the coupling of the inflaton to gauge fields localized to $W$. In string theory, higher-order flux induced corrections to the Kähler potential are not well understood so that their analysis is beyond the scope of this letter. One should keep in mind that they could be potentially dangerous.

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However, in $[26]$ the authors pointed out that there exist a tension between the chirality of intersecting $D7$-branes

from the magnetic flux on the $D7$-branes and $\chi(D_a)$ is the Euler characteristic of the complex surface $D_a$.

For simplicity, we first assume that the $D7$-branes wrap the small cycle, even though this is against the instanton zero-mode arguments of $[23]$. We come back to this point below. In order to obtain chirality, the Standard-Model ($SM$) $D7$-branes are also equipped with a non-vanishing background gauge flux $f$. Therefore, the axionic coupling to a SM brane becomes
\[
\mathcal{L}_{SM} \simeq \frac{1}{M_{pl}} \left( a_s + h(f) \theta \right) F \wedge F \bigg|_{SM}.
\]

To satisfy the $D7$-brane tadpole cancellation conditions, generically additional hidden $D7$-branes have to be present. However, these branes can usually be chosen non-chiral so that they just serve as “filler” branes. Their axionic couplings are therefore
\[
\mathcal{L}_{hid} \simeq \frac{1}{M_{pl}} a_s F \wedge F \bigg|_{hid}.
\]

Thus, there is a clear distinction in couplings to the axions between these two kinds of $D7$-branes. Note that if there are also $D3$-branes present, they would couple to the axion in a direct way
\[
\mathcal{L}_{D3} \simeq \frac{1}{M_{pl}} \theta F \wedge F \bigg|_{D3}.
\]

Reheating for natural inflation

For natural inflation, we recall from equation (21) that the inflaton is mostly $\psi^{(2)}$ which couples to the SM sector dominantly as
\[
\mathcal{L}_{SM} \simeq \frac{1}{M_{pl}} h(f_{SM}) \psi^{(2)} F \wedge F \bigg|_{SM},
\]
while it couples to the hidden sector as
\[
\mathcal{L}_{hid} \simeq \frac{1}{M_{pl}} \lambda_f \psi^{(2)} F \wedge F \bigg|_{hid}.
\]

Assuming that we can saturate the $D3$-brane tadpole condition (31) by the flux contributions $N_{\text{flux}}$ and $N_{\text{gauge}}$, i.e. without $D3$ branes, the hidden sector decays are suppressed relative to SM decays as
\[
\frac{\Gamma(\theta \rightarrow \gamma_{\text{hid}})}{\Gamma(\theta \rightarrow \gamma_{\text{SM}})} = \lambda_f^2,
\]
and reheating predominantly occurs into the SM degrees of freedom. For the decay rate induced by a dimension-five operator $\frac{\lambda_f^2}{M_{pl}^6} \theta F \wedge F$, here we have used the formula
\[
\frac{\Gamma(\theta \rightarrow \gamma)}{\Gamma(\theta \rightarrow \gamma_{\text{SM}})} = \frac{g^2 m_f^2}{64 \pi M_{pl}^2}.
\]

However, in $[26]$ the authors pointed out that there exist a tension between the chirality of intersecting $D7$-branes

\[
\mathcal{L}_{SM} \simeq \frac{1}{M_{pl}} \left( a_s + h(f) \theta \right) F \wedge F \bigg|_{SM},
\]

To satisfy the $D7$-brane tadpole cancellation conditions, generically additional hidden $D7$-branes have to be present. However, these branes can usually be chosen non-chiral so that they just serve as “filler” branes. Their axionic couplings are therefore
\[
\mathcal{L}_{hid} \simeq \frac{1}{M_{pl}} a_s F \wedge F \bigg|_{hid}.
\]

Thus, there is a clear distinction in couplings to the axions between these two kinds of $D7$-branes. Note that if there are also $D3$-branes present, they would couple to the axion in a direct way
\[
\mathcal{L}_{D3} \simeq \frac{1}{M_{pl}} \theta F \wedge F \bigg|_{D3}.
\]
and stabilizing the corresponding four-cycles via $E3$-brane instantons. This can be reconciled by sequestering the SM and moduli-stabilization sectors, that is placing the chiral Standard Model on a different del-Pezzo surface. Such a scenario was analyzed in detail in [27]. In this case, the complex Kähler moduli controlling the SM sector are stabilized via D-terms and the generalized Green-Schwarz mechanism. Performing the same computation as before with

$$V_{ab} = \begin{pmatrix} \frac{O(1)}{V^2} & 0 \\ 0 & \frac{O(1) \lambda_f}{V^2} \end{pmatrix},$$

one finds

$$\phi^{(1)} \simeq O(1) \frac{V^2}{\lambda_f} \psi^{(1)} + O(1) \frac{\lambda_f}{V^2} \psi^{(2)},$$

$$\phi^{(2)} \simeq O(1) \frac{V^2}{\lambda_f} \psi^{(1)} + O(1) \psi^{(2)}.$$ (40)

Therefore, compared to (37) the hidden sector decays are further volume suppressed

$$\frac{\Gamma(\theta \rightarrow \gamma_{hid}^2)}{\Gamma(\theta \rightarrow \gamma_{SM})} = \left(\frac{\lambda_f}{V^2}\right)^2.$$ (41)

**Reheating for axion monodromy inflation**

Performing the same analysis for our model of axion-monodromy inflation using (29), we find that the hidden sector decays are also suppressed relative to SM decays

$$\frac{\Gamma(\theta \rightarrow \gamma_{hid}^2)}{\Gamma(\theta \rightarrow \gamma_{SM})} = \frac{\lambda_0^4}{V^2}.$$ (42)

Localizing the SM on a sequestered $D7$-brane, the suppression is even stronger

$$\frac{\Gamma(\theta \rightarrow \gamma_{hid}^2)}{\Gamma(\theta \rightarrow \gamma_{SM})} = \lambda_0^4 V^2.$$ (43)

Thus, both universal natural inflation and universal axion monodromy inflation provide a natural mechanism for reheating mainly into SM degrees of freedom. The reheating temperature comes out as

$$T_R \simeq \sqrt{V M_{pl}} \simeq h(f_{SM}) \sqrt{\frac{m_0^2}{M_{pl}}} \simeq 10^{10} \text{GeV},$$

which is much higher than the big bang nucleosynthesis temperature $T_{BBN} \sim 1 \text{ MeV}$.

**Inflationary scales**

Finally, let us estimate the relevant scales. Choosing as a reference value $\lambda_0 = 4 \cdot 10^{-3}$ and comparing the inflaton mass $m_0 \simeq 10^{13} \text{GeV}$ with $M_0$ of our sequestered axion monodromy inflation model in (28), we find for the volume $V = 580$. The suppression factor in (43) is then $(\lambda_0/V)^4 \simeq 10^{-21}$. The string scale $M_s \simeq \frac{M_0}{V^{3/2}}$ comes out as $M_s \simeq 10^{17} \text{GeV}$ and the gravitino mass $M_{3/2} = \frac{M_s}{\sqrt{h}}$ as $M_{3/2} \simeq 4 \cdot 10^{13} \text{GeV}$. In the sequestered LVS scenario [27], the soft terms are suppressed as

$$m_{soft} \simeq \frac{M_0}{V^{3/2}} \simeq 7 \cdot 10^{12} \text{GeV},$$

pointing to an intermediate scale of supersymmetry breaking, as it was also found recently in [28].

**6. Conclusions**

In this letter we have proposed a type IIB, LVS-like string-realization of both natural inflation and axion monodromy inflation, where the role of the inflaton is played by the universal axion, whose scalar partner is the dilaton.

Concerning natural inflation, under the assumption that the universal axion is not already fixed by a leading-order flux-induced potential, i.e. its shift symmetry is intact, we have argued that a non-perturbative contribution to the superpotential coming from a magnetized $E3$-brane instanton gives rise to a leading-order potential. The resulting mass of the axion turned out to be smaller than that of the Kähler moduli and than the small-cycle axion.

Concerning axion monodromy inflation, we have argued that shift-symmetry breaking fluxes can still allow for an axion of parametrically small mass. Interestingly, this shift-symmetry breaking potential could be quadratic in the axion, thus providing a string-derived candidate of chaotic inflation.

In both cases, the big cycle axion was still massless at this stage and could still lead to dark radiation [29, 30] via the large Kähler modulus decay.

As one of main results, we have described a natural mechanism guaranteeing that at the end of inflation the inflaton predominantly decays into SM particles. This is achieved by having only the SM branes carry chirality-inducing gauge flux, which leads to a direct coupling of the inflaton to the SM degrees of freedom. Note that this mechanism works for both natural and axion monodromy inflation.

The predictions of universal axion inflation could be met by choosing the overall volume $V$ of the compactification space to be of the order $10^2 - 10^3$. This led to soft masses in the sequestered LVS of the order of $10^{12} \text{GeV}$, implying a high-scale susy breaking. The other moduli are very heavy in this model, so that there is no cosmological moduli problem.

Having the axion decay constant larger than the Planck scale required the string coupling constant to be larger than one. We have presented one argument why this (F-theoretic) regime might still be under control in the LVS. This is of course the weakest point of our model; nevertheless, we think that it shows some new and interesting features. Moreover, in this letter we have simply assumed that the dilaton can be stabilized by either fluxes or no-scale breaking effects such that its mass is hierarchically
bigger than the axion mass scale. However, the corresponding stabilization mechanism of the axio-dilaton deserves a closer technical investigation [19].

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