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**NON-STANDARD PRIMORDIAL FLUCTUATIONS
AND NONGAUSSIANTY IN STRING INFLATION**

(Dedicated to Lev Kofman, whose untimely passing makes the Universe a little bit darker)

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Abstract

Inflationary scenarios in string theory often involve a large number of light scalar fields, whose presence can enrich the post-inflationary evolution of primordial fluctuations generated during the inflationary epoch. We provide a simple example of such post-inflationary processing within an explicit string-inflationary construction, using a Kähler modulus as the inflaton within the framework of LARGE Volume Type-IIB string flux compactifications. We argue that inflationary models within this broad category often have a selection of scalars that are light enough to be cosmologically relevant, whose contributions to the primordial fluctuation spectrum can compete with those generated in the standard way by the inflaton. These models consequently often predict nongaussianity at a level, $f_{NL} \simeq \mathcal{O}(10)$, potentially observable by the Planck satellite, with a bi-spectrum maximized by triangles with squeezed shape in a string realization of the curvaton scenario. We argue that the observation of such a signal would robustly prefer string cosmologies such as these that predict a multi-field dynamics during the very early universe.

1 Introduction

Standard Hot Big Bang cosmology provides a good description of the great wealth of large-scale observations [1] that have recently revolutionized our understanding of cosmology, but it only does so if the universe is started off with a particular kind of initial conditions. Cosmic inflation [2] was initially proposed as an elegant way of obtaining these conditions as the outcome of still-earlier dynamics. But this initial promise was subsequently reinforced by the observation that curvature perturbations generated by quantum fluctuations of the inflaton field can get imprinted on the temperature distribution of the Cosmic Microwave Background (CMB) in the much later universe, in good agreement with the almost-scale-invariant and Gaussian spectrum that is observed.

Obtaining the desired inflationary expansion within a realistic picture of the at-present ill-understood dynamics appropriate to the very high energies required proved to be much harder than expected, however. The last decade has seen some progress, sparked by the understanding of modulus stabilization within string theory. This allows the construction of calculable inflationary configurations within string theory, with the role of the inflaton played either by an open-string degree of freedom — such as the relative positions of BPS branes [3], or of a brane and anti-brane [4], or Wilson lines [5] — or a field from the closed-string sector — such as a geometrical modulus [6] (see [7, 8] for reviews).

Nowadays, various string inflationary models are under reasonably good theoretical control, and developed to a level that can be compared meaningfully to cosmological data. In particular, because mechanisms now exist to stabilize moduli, it is possible to understand the cosmological evolution of *all* of the relevant fields, and therefore to be sure that the motion of fields other than the inflaton do not ruin the simplest single-field inflationary predictions for the evolution of curvature perturbations. It is largely the removal of this potential theoretical error that now makes the predictions of string inflationary scenarios sufficiently reliable for comparisons with observations.

One feature common to the string inflationary models explored so far is the effective absence in them of isocurvature fluctuations in the predictions for CMB observables. This despite the fact that most scenarios involve more than one potentially cosmologically active scalar field during the inflationary epoch. Indeed models are usually designed this way, with all of the non-inflaton moduli sitting in their local minima as the inflaton rolls. Such constructions greatly simplify the calculation of late-time perturbations, because they predict only adiabatic fluctuations, which can be evolved forward to the present time with minimal sensitivity to the details of the poorly-understood cosmological history between inflation and now. It is because of this that the implications of these models are usually well-captured, *ex post facto*, by simple single-field inflationary models [8, 9].

An unfortunate consequence of the focus for convenience on such models is the misconception that string inflation *must* agree in its predictions with single-field models, including in particular a

prediction of vanishingly small nongaussianity. This prediction is sometimes held up as a potential observational way to discriminate [10] between string inflation and alternatives to inflation within string theory [11].

In order to investigate the robustness of such predictions, in this paper we take the first steps towards exploring other mechanisms for generating primordial fluctuations within a concrete string inflationary model based on a LARGE Volume (LV) scenario. (For other discussions of nongaussianities in string inspired scenarios see [12].) We find we are able to construct such string inflationary frameworks by making nontrivial use of the presence of the large number of scalar fields that are generically present during and after the inflationary epoch. If these fields are sufficiently light during inflation they can acquire significant isocurvature fluctuations which post-inflationary evolution can robustly convert into adiabatic perturbations, swamping those contributions coming from the inflaton field itself. Although somewhat more history-dependent than is the standard mechanism, the subsequent evolution of the resulting adiabatic fluctuations remains plausibly independent of the details of cosmic evolution provided only that the universe comes to thermal equilibrium shortly after adiabatic perturbations with the desired features are produced.

Our search for models uses two generic mechanisms for achieving post-inflationary isocurvature to adiabatic conversion: the curvaton mechanism [13, 14, 15]; and the modulation mechanism [16, 17, 18, 19] scenarios. In particular, the main models we present can be regarded as explicit realizations of the curvaton mechanism within a string-inflationary framework. The idea that such modulus dependent effects could contribute to curvature perturbations is not in itself new. What we accomplish in this work is to achieve it for the first time in a fully calculable string set-up, where the required properties are subject to a myriad of constraints imposed by the underlying UV consistency. The precision of this kind of setup is a necessary preliminary for asking more detailed questions about reheating and the ultimate transfer of energy from the inflaton to observable degrees of freedom (d.o.f.). Similar studies for brane-antibrane inflation allowed the identification of cosmic strings as a potential late-epoch signature [20], as well as the utility of warping for channeling energy into the observed low-energy sector [21]. Recent studies of reheating at the end of closed string inflation similarly reveal the need to set severe constraints on the hidden sector dynamics in order to allow an efficient reheating of the visible sector [22].

Because the observed primordial fluctuations are not directly generated by the inflaton their properties in general depend differently on the various underlying parameters, and are in particular not tied to the slow-roll parameters in the way familiar from single-field models. This could ultimately allow string inflationary models for which the string scale is not in conflict with the demands of particle physics during the present epoch (such as the supersymmetry breaking scale) [23], although we do not yet have an explicit example which does so.

The most interesting such difference is the generic prediction of a sizeable level of nongaussianity, $f_{NL} \simeq \mathcal{O}(10)$; a level detectable by the Planck satellite. The underlying imprinting of

the adiabatic fluctuations takes place in the post-inflationary epoch. It is characterized by a non-linear relation between scalar and curvature fluctuations, that generates nongaussianities of local form. The corresponding bi-spectrum, consequently, is robustly predicted to be dominated by triplets of momenta that form long, thin triangles: the so-called squeezed limit.¹

In the models studied here the size of the nongaussianity is a consequence of the properties of the geometry of the extra dimensions in string theory. But if such nongaussianity should really be observed with these properties, they will not tell us about microscopic physics in this much detail. What they most likely would tell us is that the epoch of fluctuation generation and its aftermath are described by some sort of multi-field system similar to the ones we describe.

We perform our search for these mechanisms within the LARGE Volume (LV) scenario of modulus stabilization for Type IIB string vacua [25]. These models are convenient for this purpose for several reasons. First, they predict the existence of a suite of moduli, whose masses naturally come with a hierarchical suppression in different ways by powers of the extra-dimensional volume, $\mathcal{V} = \text{Vol}/\ell_s^6$, in string units [26, 27, 28, 29]. In particular, Kähler moduli for small cycles tend to arise with masses of order M_p/\mathcal{V} while those for large cycles tend to get masses of the order $M_p/\mathcal{V}^{3/2}$ or smaller. Second, the couplings of these moduli to observable fields at late times can be plausibly estimated provided these fields are assumed to reside on a brane (or branes) that wrap the cycles whose volumes are measured by the various moduli [22, 27, 30]. Finally, inflationary mechanisms are already known using these models, with the inflaton being either a small cycle [31] or a large one [32].

To construct our models we splice the frameworks developed in [31] and [32], using a modulus of a small (blow-up) cycle as the inflaton, keeping the modulus for a larger cycle as the (curvaton) field that acquires isocurvature fluctuations. This construction exploits the fact that these moduli like to be light relative to the inflaton, and so would plausibly have extra-Hubble fluctuations imprinted on their profiles. Moreover, after inflation its decay rate to radiation has the right value to convert isocurvature modes into adiabatic fluctuations, with the correct amplitude (and a sizeable level of nongaussianity). The spirit of our construction is to provide an existence proof for mechanisms of this type within a well-developed, modern string set-up, in which issues associated with moduli dynamics and stabilization can be analysed. Although at first sight the model may seem contrived, it actually uses the minimal amount of ingredients that are needed in order to exhibit the effects we are interested in.

The paper is organized as follows. In §2, we briefly review the field content and framework of the LV compactifications. §3 then describes the inflationary setup in these models, which minimally involve 4 moduli: $\mathcal{V} = \mathcal{V}(\tau_1, \tau_2, \tau_3, \tau_4)$. These are: a curvaton field, τ_1 ; the volume modulus, τ_2 together with a blow-up mode, τ_3 , that provides the standard LV stabilization mechanism for the volume \mathcal{V} ; and an inflaton τ_4 .² These have the desired hierarchy of masses if the fluxes are

¹For a recent comprehensive review on nongaussianities see [24].

²We apologize for the slightly opaque notation, which is designed to follow ref. [32] as closely as possible.

adjusted so that the volumes are stabilized with the hierarchy $\tau_2 > \tau_1 \gg \tau_4 > \tau_3$. Because of the LV ‘magic’ this can be done using hierarchies among the input fluxes that are at most $\mathcal{O}(10)$. §4 then gives the \mathcal{V} -dependence of the couplings of these fields to observable d.o.f., which we take to be localized on a brane wrapping either the curvaton cycle or one of the small blow-up cycles. The curvaton mechanism in this framework is explored in §5, where it is shown that the \mathcal{V} -dependence of the masses and couplings can be such as to produce acceptable adiabatic fluctuations. §6 then explores several choices for underlying parameters to get a feel for the range that is possible for observables. One of the models presented in this section predicts $f_{NL} \simeq 56$. Our conclusions are briefly summarized in §7.

2 The system under consideration

We start with a discussion of the system whose inflationary dynamics is of interest. We follow throughout the conventions of [32].

2.1 The field content

The model requires us to choose a compactification based on a Calabi-Yau manifold having at least the following 4 Kähler moduli, whose dynamics are of interest:

- i) A fiber modulus, τ_1 , playing the role of curvaton field and wrapped by a stack of $D7$ -branes.³ The low-energy scalar potential first acquires a dependence on this field through string loop contributions sourced by the $D7$ -branes [28, 34], and for this reason it likes to remain light during inflation.
- ii) A base modulus, τ_2 , that mainly controls the overall extra-dimensional volume and which is wrapped by a stack of $D7$ -branes needed to generate the string loop potential for τ_1 .³ This modulus is heavy during inflation, and remains well-stabilized at its minimum throughout inflation.
- iii) A blow-up mode τ_3 , that is an ‘assisting field’ required to stabilize the volume \mathcal{V} at its minimum in the usual LV way (as in [31]). The potential depends on it through non-perturbative contributions generated by a stack of $D7$ -branes wrapping τ_3 and supporting a hidden sector that undergoes gaugino condensation.⁴ It is heavy during inflation, and its VEV is proportional to the logarithm of the volume.
- iv) A second blow-up mode, τ_4 , that plays the role of the inflaton field, as in [31]. Its non-perturbative potential is again generated by gaugino condensation on the hidden sector

³These $D7$ -branes can support either a visible or a hidden sector in a very model-dependent way.

⁴ τ_3 cannot support a visible sector due to the tension between non-perturbative effects and chirality [33].

supported by a stack of $D7$ -branes wrapping this cycle.⁵ During inflation its VEV is few times the logarithm of the volume.

We finally point out that we shall present two explicit scenarios:

1. Visible sector wrapped around the curvaton cycle τ_1 (and τ_2 since these two cycles intersect each other): this is the case with the minimal number of 4 Kähler moduli and, due to the location of the visible sector on τ_1 , it maximizes the strength of the coupling of the curvaton to visible d.o.f., so yielding the largest amount of nongaussianities. However, this is a non-standard realization of the visible sector supported on a non-rigid 4-cycle which tends to be stabilized large (giving rise to a tiny gauge coupling).
2. Visible sector wrapped around a blow-up mode τ_5 which is heavy during inflation: in this case we need to make the system a bit more involved including a fifth cycle which can be stabilized small either in the geometric regime (by string loop effects as in [28]) or at the quiver locus (by D -terms as in [35]). The advantage is that now we have a standard realization of the visible sector on a rigid 4-cycle whose VEV reproduces the correct order of magnitude of the gauge coupling. However now the geometric separation between τ_1 and τ_5 reduces the strength of the coupling of the curvaton to visible d.o.f., so yielding a smaller amount of nongaussianities.

The compactification

To have these four moduli we consider a Calabi-Yau three-fold with a K3 fibration structure controlled by two moduli, τ_1 and τ_2 , together with two additional blow-up modes, τ_3 and τ_4 . We assume the Calabi-Yau volume when expressed as a function of these moduli has the form [36]:

$$\mathcal{V} = \alpha \left(\sqrt{\tau_1} \tau_2 - \sum_{i=3}^4 \gamma_i \tau_i^{3/2} \right). \quad (2.1)$$

The Kähler potential (including the leading α' corrections) for the effective low-energy 4D supergravity in this case is (we work throughout in the 4D Einstein frame):

$$K \simeq K_0 + \delta K_{\alpha'} = -2 \ln \left[\mathcal{V} + \frac{\hat{\xi}}{2} \right], \quad (2.2)$$

where the α' corrections are controlled by the quantity

$$\hat{\xi} \equiv \frac{\xi}{g_s^{3/2}} = -\frac{\zeta(3) \chi(M)}{2 g_s^{3/2} (2\pi)^3} \quad (2.3)$$

where $\chi(M)$ is the Euler number of the compact manifold. In applications we take the quantity ξ to lie in the interval $(0.1, 1)$.

⁵The potential for τ_4 cannot be generated by an instanton since after inflation it would lead τ_4 to a regime, $\langle \tau_4 \rangle < 1$, where we cannot trust the effective field theory [22].

In the superpotential we neglect non-perturbative contributions associated with the large cycles, τ_1 and τ_2 , relative to those of the small cycles,

$$W \simeq W_0 + A_3 e^{-a_3 T_3} + A_4 e^{-a_4 T_4}, \quad (2.4)$$

since these are negligible relative to those explicitly written and are likely to be absent since τ_1 and τ_2 are non-rigid cycles. The superpotential is characterized by the constant W_0 and the non-perturbative corrections are weighted by constants A_i . We choose the quantity W_0 – as usual in LV models – to be order one, and the parameters a_i satisfy $a_i = 2\pi/N$ since they arise due to gaugino condensation on $D7$ branes (with N being the rank of the associated gauge group).

Following the LV program, our interest is in the form of the resulting potential in a regime where $\ln \mathcal{V} \simeq \mathcal{O}(\tau_3)$, so that terms in the α' expansion compete with the leading non-perturbative contributions from W [26]. However, for the inflationary analysis our interest is not in the local LV minimum. Instead we seek nearby flat regions of the potential along which the potential is shallow as a function of τ_1 and τ_4 , with \mathcal{V} and τ_3 heavy enough to sit at their local minima. Following the reasoning of refs. [31] and [32], we expect such a regime to arise in the region of field space where the fields are hierarchically different: $\tau_2 > \tau_1 \gg \tau_4 > \tau_3$, since in this region the potential likes to become independent of τ_1 and τ_4 , at least before string-loop contributions are included.

We now use these expressions to compute the scalar potential and kinetic terms in the desired regime.

2.2 The kinetic terms

In this section we investigate the field redefinitions needed to put the kinetic terms into canonical form. The starting point in the regime of interest is the Kähler metric for the moduli, which is given by the following symmetric matrix:

$$K_{ij}^0 = \frac{1}{4\tau_2^2} \begin{pmatrix} \frac{\tau_2^2}{\tau_1^2} & \frac{\gamma_3 \tau_3^{3/2} + \gamma_4 \tau_4^{3/2}}{\tau_1^{3/2}} & -\frac{3\gamma_3}{2} \frac{\sqrt{\tau_3}}{\tau_1^{3/2}} \tau_2 & -\frac{3\gamma_4}{2} \frac{\sqrt{\tau_4}}{\tau_1^{3/2}} \tau_2 \\ " & 2 & -3\gamma_3 \frac{\sqrt{\tau_3}}{\sqrt{\tau_1}} & -3\gamma_4 \frac{\sqrt{\tau_4}}{\sqrt{\tau_1}} \\ " & " & \frac{3\alpha\gamma_3}{2} \frac{\tau_2^2}{\mathcal{V}\sqrt{\tau_3}} & \frac{9\gamma_3\gamma_4}{2} \frac{\sqrt{\tau_3\tau_4}}{\tau_1} \\ " & " & " & \frac{3\alpha\gamma_4}{2} \frac{\tau_2^2}{\mathcal{V}\sqrt{\tau_4}} \end{pmatrix}, \quad (2.5)$$

where (as in [32]) we systematically drop terms that are suppressed relative to the ones shown by factors $\sqrt{\tau_i/\tau_2} \forall i = 3, 4$.

The kinetic Lagrangian to leading order therefore becomes⁶

$$\begin{aligned}
-\frac{\mathcal{L}_{kin}}{\sqrt{-g}} &= \frac{1}{4\tau_1^2} (\partial\tau_1)^2 + \frac{1}{2\tau_2^2} (\partial\tau_2)^2 + \sum_{i=3}^4 \frac{3\alpha\gamma_i}{8\mathcal{V}\sqrt{\tau_i}} (\partial\tau_i)^2 + \sum_{i=3}^4 \frac{\gamma_i\tau_i^{3/2}}{2\tau_2^2\tau_1^{3/2}} \partial\tau_1\partial\tau_2 \\
&\quad - \sum_{i=1}^4 \frac{3\alpha\gamma_i\sqrt{\tau_i}}{2\mathcal{V}} \left(\frac{\partial\tau_1}{2\tau_1} + \frac{\partial\tau_2}{\tau_2} \right) \partial\tau_i + \frac{9\alpha^2\gamma_3\gamma_4}{4} \frac{\sqrt{\tau_3\tau_4}}{\mathcal{V}^2} \partial\tau_3\partial\tau_4 \\
&= \frac{3}{8\tau_1^2} (\partial\tau_1)^2 + \frac{1}{2\mathcal{V}^2} (\partial\mathcal{V})^2 + \sum_{i=3}^4 \frac{3\alpha\gamma_i}{8\mathcal{V}\sqrt{\tau_i}} (\partial\tau_i)^2 - \frac{1}{2\tau_1\mathcal{V}} \partial\tau_1\partial\mathcal{V} \\
&\quad - \sum_{i=3}^4 \frac{3\alpha\gamma_i\sqrt{\tau_i}}{2\mathcal{V}^2} \partial\mathcal{V}\partial\tau_i + \frac{9\alpha^2\gamma_3\gamma_4}{4} \frac{\sqrt{\tau_3\tau_4}}{\mathcal{V}^2} \partial\tau_3\partial\tau_4, \tag{2.6}
\end{aligned}$$

where the last equality trades τ_2 for \mathcal{V} , in the limit in which τ_1, τ_2 are much larger than τ_3, τ_4 . It is convenient to canonically normalize order by order in $1/\mathcal{V}$, and so we rewrite (2.6) as:

$$\mathcal{L}_{kin} = \mathcal{L}_{kin}^{\mathcal{O}(1)} + \mathcal{L}_{kin}^{\mathcal{O}(\mathcal{V}^{-1})} + \mathcal{L}_{kin}^{\mathcal{O}(\mathcal{V}^{-2})}, \tag{2.7}$$

where the leading term is

$$-\frac{\mathcal{L}_{kin}^{\mathcal{O}(1)}}{\sqrt{-g}} = \frac{3}{8\tau_1^2} (\partial\tau_1)^2 + \frac{1}{2\mathcal{V}^2} (\partial\mathcal{V})^2 - \frac{1}{2\tau_1\mathcal{V}} \partial\tau_1\partial\mathcal{V}, \tag{2.8}$$

while the subleading terms are

$$-\frac{\mathcal{L}_{kin}^{\mathcal{O}(\mathcal{V}^{-1})}}{\sqrt{-g}} = \sum_{i=3}^4 \frac{3\alpha\gamma_i}{8\mathcal{V}\sqrt{\tau_i}} (\partial\tau_i)^2 - \sum_{i=3}^4 \frac{3\alpha\gamma_i\sqrt{\tau_i}}{2\mathcal{V}^2} \partial\mathcal{V}\partial\tau_i, \tag{2.9}$$

at $\mathcal{O}(1/\mathcal{V})$ and

$$-\frac{\mathcal{L}_{kin}^{\mathcal{O}(\mathcal{V}^{-2})}}{\sqrt{-g}} = \frac{9\alpha^2\gamma_3\gamma_4}{4} \frac{\sqrt{\tau_3\tau_4}}{\mathcal{V}^2} \partial\tau_3\partial\tau_4, \tag{2.10}$$

at $\mathcal{O}(\mathcal{V}^{-2})$. At $\mathcal{O}(1)$ the transformation

$$\tau_1 = \exp(a\chi_1 + b\chi_2), \tag{2.11}$$

$$\mathcal{V} = \exp(c\chi_2), \tag{2.12}$$

puts expression (2.8) into canonical form

$$-\frac{\mathcal{L}_{kin}^{\mathcal{O}(1)}}{\sqrt{-g}} = \frac{1}{2} \left[(\partial\chi_1)^2 + (\partial\chi_2)^2 \right], \tag{2.13}$$

where the coefficients a, b and c are obtained from the condition that the matrix $M = \begin{pmatrix} a & b \\ 0 & c \end{pmatrix}$ satisfies

$$M^T \cdot \begin{pmatrix} \frac{3}{4} & -\frac{1}{2} \\ -\frac{1}{2} & 1 \end{pmatrix} \cdot M = I. \tag{2.14}$$

⁶We use units with $8\pi M_p = 1$ unless otherwise stated.

This has four solutions: (a, b, c) , $(a, -b, -c)$, $(-a, b, c)$, $(-a, -b, -c)$, where:

$$a = \frac{2}{\sqrt{3}}, \quad b = \sqrt{\frac{2}{3}}, \quad c = \sqrt{\frac{3}{2}}, \quad (2.15)$$

and for concreteness we shall choose the first one with all plus signs. As is shown in [22], the fields χ_1 and χ_2 turn out to also diagonalize the mass-squared matrix, $M_{ij}^2 = \sum_k K_{ik}^{-1} V_{kj}$ in the limit where string-loop corrections to the potential are neglected. Once string-loop corrections are included a subdominant dependence of \mathcal{V} on χ_1 also arises that is not important for our purposes.

Next we diagonalize the next-order kinetic term, $\mathcal{L}_{kin}^{\mathcal{O}(\mathcal{V}^{-1})}$. The first term in (2.9) becomes diagonal once we rescale the two small moduli as follows

$$\tau_j = \left(\frac{3\mathcal{V}}{4\alpha\gamma_j} \right)^{2/3} \phi_j^{4/3}, \quad \forall j = 3, 4 \quad (2.16)$$

where we use the notation ϕ_j with $j = 3, 4$ to distinguish these from the large fields, χ_1 and χ_2 . The second term in (2.9) is similarly diagonalized by mixing \mathcal{V} with $\tau_j \forall j = 3, 4$. Explicitly, introducing the following subleading corrections to (2.11) and (2.12):

$$\tau_1 = \exp \left(\frac{2}{\sqrt{3}} \chi_1 + \sqrt{\frac{2}{3}} \chi_2 + \frac{3}{2} \sum_{j=3}^4 \phi_j^2 \right), \quad (2.17)$$

$$\mathcal{V} = \exp \left(\sqrt{\frac{3}{2}} \chi_2 + \frac{9}{4} \sum_{j=3}^4 \phi_j^2 \right), \quad (2.18)$$

gives to this order

$$\mathcal{L}_{kin}^{\mathcal{O}(1)} + \mathcal{L}_{kin}^{\mathcal{O}(\mathcal{V}^{-1})} = \frac{1}{2} \sum_{i=1}^2 (\partial\chi_i)^2 + \frac{1}{2} \sum_{j=3}^4 (\partial\phi_j)^2. \quad (2.19)$$

Notice that the last term in eqs. (2.17) and (2.18) is subleading because $\phi_j \sim \mathcal{O}(\mathcal{V}^{-1/2}) \ll 1$ for $j = 3, 4$, while from (2.11) and (2.12), we have $\chi_i \sim \mathcal{O}(\ln \mathcal{V})$, for $i = 1, 2$. We can now substitute (2.18) in (2.16) to eliminate \mathcal{V} and directly express τ_i in terms of ϕ_i , for $i = 3, 4$, obtaining

$$\begin{aligned} \tau_i &= \left(\frac{3}{4\alpha\gamma_i} \right)^{2/3} \left[\exp \left(\sqrt{\frac{3}{2}} \chi_2 + \frac{9}{4} \sum_{j=3}^4 \phi_j^2 \right) \right]^{2/3} \phi_i^{4/3} \\ &\simeq \left(\frac{3}{4\alpha\gamma_i} \right)^{2/3} \exp \left[\sqrt{\frac{2}{3}} \chi_2 \right] \left(1 + \frac{3}{2} \sum_{i \neq j=3}^4 \phi_j^2 \right) \phi_i^{4/3}, \quad \forall i = 3, 4. \end{aligned} \quad (2.20)$$

Notice that passage from the first to the second line neglects subleading contributions controlled by higher order powers of ϕ_i .

Finally, the off-diagonal term in $\mathcal{L}_{kin}^{\mathcal{O}(\mathcal{V}^{-2})}$ is removed by modifying (2.20) slightly, into:

$$\tau_i \simeq \left(\frac{3}{4\alpha\gamma_i} \right)^{2/3} \exp \left[\sqrt{\frac{2}{3}} \chi_2 \right] \left(1 - \frac{3}{4} \sum_{i \neq j=3}^4 \phi_j^2 \right) \phi_i^{4/3}, \quad \forall i = 3, 4, \quad (2.21)$$

$$\simeq \left(\frac{3}{4\alpha\gamma_i} \right)^{2/3} \mathcal{V}^{2/3} \left(1 - \frac{9}{4} \sum_{i \neq j=3}^4 \phi_j^2 \right) \phi_i^{4/3}, \quad \forall i = 3, 4. \quad (2.22)$$

The field redefinitions we have determined render canonical the form of the kinetic terms.

2.3 The potential

We next chase these field redefinitions through the definition of the scalar potential, again following the discussion of ref. [32].

The potential without loop corrections

After minimizing the axion directions, the scalar potential constructed using the Kähler potential and superpotential of eqs. (2.2) and (2.4) (and neglecting subleading powers of large moduli) is

$$V = \frac{g_s e^{K_{cs}}}{8\pi} \left[\sum_{i=3}^4 \frac{8 a_i^2 A_i^2}{3\alpha\gamma_i} \left(\frac{\sqrt{\tau_i}}{\mathcal{V}} \right) e^{-2a_i\tau_i} - 4 \sum_{i=3}^4 W_0 a_i A_i \left(\frac{\tau_i}{\mathcal{V}^2} \right) e^{-a_i\tau_i} + \frac{3\beta \hat{\xi} W_0^2}{4\mathcal{V}^3} \right]. \quad (2.23)$$

The overall factor of $g_s e^{K_{cs}}/(8\pi)$ in front of the potential is a consequence of an overall normalization of the superpotential, that is needed to express all quantities in the 4D Einstein frame [30], as is explained in detail in Appendix A.⁷ The constant β appearing in the last, τ -independent, term,

$$V_0 \equiv \frac{3g_s \beta \hat{\xi} W_0^2}{32\pi \mathcal{V}^3}, \quad (2.24)$$

includes contributions due to the stabilization of the massive field τ_2 , and due to uplifting terms. It is tuned in such a way that, at the minimum for τ_4 , the potential vanishes.

This potential completely stabilizes τ_3 , τ_4 and the volume \mathcal{V} , at the following values (where we assume $a_i\tau_i \gg 1$):

$$a_i \langle \tau_i \rangle = \left(\frac{\hat{\xi}}{2\alpha J} \right)^{\frac{2}{3}}, \quad \langle \mathcal{V} \rangle = \left(\frac{3\alpha\gamma_i}{4a_i A_i} \right) W_0 \sqrt{\langle \tau_i \rangle} e^{a_i \langle \tau_i \rangle}, \quad \forall i = 3, 4, \quad (2.25)$$

where $J = \sum_{i=3}^4 \gamma_i/a_i^{3/2}$. What is noteworthy is that eq. (2.23) does not depend at all on the fibre modulus, τ_1 [32]. It does not do so because the dominant contribution to the potential of large moduli such as these first arises at the string loop level [28], whose size we now estimate.

The potential with loop corrections

Each cycle wrapped by a stack of $D7$ -branes receives 1-loop open string corrections [28, 34] which, as pointed out in [32], spoil the flatness of the inflationary potential for τ_4 . However it is possible to fine-tune the coefficient of the τ_4 -dependent loop correction in order to render it negligible (the amount of fine-tuning needed has been estimated in [22]). Hence we shall focus only on the τ_1 and τ_2 -dependent loop corrections which can be estimated using a procedure identical to [32]:

$$V = V_0 + \frac{g_s a_4^2 A_4^2}{3\pi \alpha \gamma_4} \left(\frac{\sqrt{\tau_4}}{\mathcal{V}} \right) e^{-2a_4\tau_4} - \frac{g_s W_0 a_4 A_4}{2\pi} \left(\frac{\tau_4}{\mathcal{V}^2} \right) e^{-a_4\tau_4} + \left(\frac{A}{\tau_1^2} - \frac{B}{\mathcal{V}\sqrt{\tau_1}} + \frac{C\tau_1}{\mathcal{V}^2} \right) \frac{g_s W_0^2}{8\pi \mathcal{V}^2}, \quad (2.26)$$

⁷From now on we shall set $e^{K_{cs}} = 1$.

where A , B , C are given by

$$A = (g_s C_1^{KK})^2 \quad (2.27)$$

$$B = 4\alpha C_{12}^W \quad (2.28)$$

$$C = 2 (\alpha g_s C_2^{KK})^2, \quad (2.29)$$

where C_1^{KK} , C_{12}^W , and C_2^{WW} are constants that depend on the details of the string loop corrections (see [32] for more details). In what follows we regard these constants as free to be fixed using phenomenological requirements.

The minimum for τ_1 is at:

$$\langle \tau_1 \rangle \simeq \left(-\frac{B\mathcal{V}}{2C} \right)^{2/3} \quad \text{if } B < 0 \quad \text{or} \quad \langle \tau_1 \rangle \simeq \left(\frac{4A\mathcal{V}}{B} \right)^{2/3} \quad \text{if } B > 0. \quad (2.30)$$

In the following, for definiteness, we consider the case $B > 0$. It is important to notice that $\langle \tau_1 \rangle$ does not depend on τ_4 , and so τ_4 and τ_1 can evolve independently in field space. String loop corrections also shift the minimum for τ_3 , with respect to its value in eq. (2.25), but this small correction does not modify the discussion that follows.

The canonically normalized potential

We next identify that part of the potential relevant to inflation. We set \mathcal{V} and τ_3 to their minima, and follow the dependence of the rest of the potential on the remaining two fields. This adiabatic approximation is valid in the large- \mathcal{V} limit because the masses of these fields are parametrically larger than those of the fields whose motion we consider.

Recall that the fields τ_1 and τ_4 are given in terms of their canonically normalized counterparts by

$$\tau_1 = \mathcal{V}^{2/3} \exp\left(\frac{2}{\sqrt{3}} \langle \chi_1 \rangle\right) e^{\frac{2}{\sqrt{3}} \hat{\chi}_1} \quad (2.31)$$

$$\tau_4 = \left(\frac{3\mathcal{V}}{4\alpha\gamma_4} \phi_4^2 \right)^{2/3} \left(1 - \frac{9}{4} \phi_3^2 \right) \simeq \left(\frac{3\mathcal{V}}{4\alpha\gamma_4} \phi_4^2 \right)^{2/3} \quad (2.32)$$

where we define

$$\chi_1 = \langle \chi_1 \rangle + \hat{\chi}_1, \quad (2.33)$$

and the approximate equality for τ_4 neglects the subleading dependence on the modulus ϕ_3 .

Keeping in mind the factors of g_s appearing in the constants A , B and C , we expect

$$32AC \ll B^2, \quad (2.34)$$

in weak coupling, and in this case one finds

$$\langle \chi_1 \rangle = 1/\sqrt{3} \ln(q\mathcal{V}), \quad \text{with} \quad q \equiv 4A/B. \quad (2.35)$$

With this information, the leading contribution to the inflationary potential breaks into a sum of terms for the would-be inflaton and curvaton

$$V(\phi_4, \hat{\chi}_1) = V_{inf}(\phi_4) + V_{cur}(\hat{\chi}_1) \quad (2.36)$$

where

$$V_{inf}(\phi_4) \simeq V_0 - \frac{g_s W_0 a_4 A_4}{2\pi \mathcal{V}^2} \left(\frac{3\mathcal{V}}{4\alpha\gamma_4} \right)^{2/3} \phi_4^{4/3} \exp \left\{ - \left[a_4 \left(\frac{3\mathcal{V}}{4\alpha\gamma_4} \right)^{2/3} \phi_4^{4/3} \right] \right\}, \quad (2.37)$$

and

$$V_{cur}(\hat{\chi}_1) = \frac{g_s W_0^2}{8\pi \mathcal{V}^{10/3}} \left[C_0 e^{\frac{2}{\sqrt{3}}\hat{\chi}_1} - C_1 e^{-\frac{1}{\sqrt{3}}\hat{\chi}_1} + C_2 e^{-\frac{4}{\sqrt{3}}\hat{\chi}_1} \right], \quad (2.38)$$

with (see [32])

$$C_0 = C q^{2/3} \quad (2.39)$$

$$C_1 = B q^{-1/3} \quad (2.40)$$

$$C_2 = A q^{-4/3}. \quad (2.41)$$

We call χ_1 the curvaton and ϕ_4 the inflaton because the potential for χ_1 is parametrically suppressed by powers of $1/\mathcal{V}$ relative to that for ϕ_4 , thereby ensuring that it is ϕ_4 whose energy dominates the cosmic expansion.

Since $\hat{\chi}_1$ has been defined such that $\hat{\chi}_1 = 0$ at the minimum of the potential, it follows that the dependence of the constants on $\langle \chi_1 \rangle$ ensures, within the limit (2.34), that they satisfy

$$\left(\frac{\partial V_{cur}}{\partial \hat{\chi}_1} \right) \Big|_{\hat{\chi}_1=0} = 0 \Rightarrow C_0 + \frac{C_1}{2} - 2C_2 = 0. \quad (2.42)$$

In the following we work in regimes with $\hat{\chi}_1$ very small, for which the exponentials in eq. (2.38) can be expanded up to quadratic order,

$$V_{cur}(\hat{\chi}_1) \simeq V_{cur,0} + \frac{m_{\chi_1}^2}{2} \hat{\chi}_1^2 \quad \text{with} \quad m_{\chi_1}^2 = \frac{g_s W_0^2}{12\pi \mathcal{V}^{10/3}} [4C_0 - C_1 + 16C_2] \equiv \frac{g_s C_t W_0^2}{8\pi \mathcal{V}^{10/3}} \quad (2.43)$$

which defines the new constant

$$C_t \equiv \frac{2}{3} [4C_0 - C_1 + 16C_2]. \quad (2.44)$$

The constant piece, $V_{cur,0}$, is absorbable into a subdominant contribution to the constant V_0 in formula (2.37). We check in our later applications that this quadratic expansion of the potential suffices in the regime of interest.

Field masses

For inflationary applications our interest is whether the masses of the various fields are larger or smaller than the Hubble scale. Considering that the inflaton potential is of order $V_{inf} \propto \mathcal{V}^{-2} e^{-a_4 \tau_4} \sim \mathcal{O}(\mathcal{V}^{-3})$, our benchmark during inflation is $H \sim M_p^2 \mathcal{V}^{-3/2}$. Relative to this consider the following masses, evaluated at the potential's minimum:

- If all the fields sit at their minima, the mass spectrum is (we temporarily reintroduce the dependence on the Planck mass):

$$m_{\phi_i}^2 \sim \frac{g_s}{4\pi} \left(\frac{W_0}{\mathcal{V}} \right)^2 M_p^2, \quad \forall i = 3, 4 \quad (2.45)$$

$$m_{\chi_2}^2 \sim \frac{g_s W_0^2}{4\pi \mathcal{V}^3} M_p^2, \quad m_{\chi_1}^2 \sim \frac{g_s W_0^2}{4\pi \mathcal{V}^3 \sqrt{\tau_1}} M_p^2 \sim \frac{g_s C_t W_0^2}{4\pi \mathcal{V}^{10/3}} M_p^2. \quad (2.46)$$

- If the inflation and curvaton fields, ϕ_4 and χ_1 , are moved away from their minima then their masses are potentially modified. Inspection of the above formulae shows that the χ_1 mass remains of the same order in $1/\mathcal{V}$ as it is at its minimum, eq. (2.46), while the ϕ_4 mass changes and becomes smaller for larger ϕ_4 . Considering the regime $a_4 \tau_4 > (2+n) \ln \mathcal{V}$, with $n > 0$ one finds

$$m_{\phi_4}^2 \simeq \frac{g_s W_0^2}{4\pi \mathcal{V}^{3+n}} M_p^2, \quad (2.47)$$

so the inflaton mass is reduced relative to eq. (2.45) as its field moves away from its minimum (as in ref. [31]).

We see from these estimates that for large \mathcal{V} , the fields χ_2 and ϕ_3 have masses that are much larger than H , while χ_1 and ϕ_4 have masses that are smaller, justifying the picture wherein \mathcal{V} (which is mostly given by χ_2) and τ_3 (which is mostly ϕ_3) can be set to their minima while both χ_1 and ϕ_4 remain light enough to have cosmic fluctuations imprinted on them. Since it is the potential for ϕ_4 that is the largest, this is the field whose evolution dictates the end of inflation and so earns the name inflaton.

3 Dynamics during inflation

We next discuss the properties of slow-roll inflation in the above regime, together with a discussion of whether χ_1 has the properties required for it to realize the curvaton scenario in this system. We find these two fields combine the results of [31] and [32].

We start with the hypothesis that the massive moduli χ_2 and ϕ_3 are already at their minima, while ϕ_4 and χ_1 need not be. We then consider the evolution of the moduli χ_1 and ϕ_4 . As we pointed out before, the analysis is comparatively simple because these two fields evolve almost independently: see the potential in eqs. (2.37) and (2.38). If the field ϕ_4 acquires a large value, the dominant term in the inflaton potential is V_0 . Within this regime, both ϕ_4 and χ_1 are lighter than the Hubble parameter. The former plays the role of inflaton field, while the latter is the candidate curvaton.

In this section we recap how ϕ_4 drives inflation, and how the field χ_1 acquires a scale independent spectrum of isocurvature fluctuations, of calculable amplitude, during this inflationary epoch. The next sections discuss how to convert the isocurvature fluctuations of χ_1 into adiabatic perturbations after inflation ends.

3.1 Dynamics of the inflaton field ϕ_4

In the scenario just described it is ϕ_4 that drives inflation, as in the model of [31]. The inflationary potential is

$$V_{inf}(\phi_4) = \frac{3g_s \beta \hat{\xi} W_0^2}{32\pi \mathcal{V}^3} - \frac{g_s W_0 a_4 A_4}{2\pi \mathcal{V}^2} \left(\frac{3\mathcal{V}}{4\alpha\gamma_4} \right)^{2/3} \phi_4^{4/3} \exp \left\{ - \left[a_4 \left(\frac{3\mathcal{V}}{4\alpha\gamma_4} \right)^{2/3} \phi_4^{4/3} \right] \right\}, \quad (3.1)$$

showing again that the scale of inflation is mainly controlled by the value of the volume, being given by V_0 in eq. (2.24).

The corresponding slow-roll parameters, expressed in terms of the field τ_4 , become

$$\epsilon = \frac{512\mathcal{V}^3}{27\gamma_4 \alpha \hat{\xi}^2 \beta^2 W_0^2} a_4^2 A_4^2 \sqrt{\tau_4} (1 - a_4\tau_4)^2 e^{-2a_4\tau_4} \quad (3.2)$$

$$\eta = -\frac{16a_4 A_4 \mathcal{V}^2}{9\alpha \hat{\xi} \gamma_4 \beta W_0 \sqrt{\tau_4}} (1 - 9a_4\tau_4 + 4a_4^2\tau_4^2) e^{-a_4\tau_4}, \quad (3.3)$$

so in the limit of large volume, in order to have ϵ and η small one must choose, at horizon exit,

$$a_4\tau_4 \simeq (2 + n) \ln \mathcal{V}, \quad (3.4)$$

with $n > 0$. The number of e -foldings is given by the integral

$$N_e = \int_{\phi_4^{end}}^{\phi_4} \frac{V_{inf}}{V'_{inf}} d\tilde{\phi} = \frac{-27\beta \hat{\xi} W_0 \gamma_4}{256 \mathcal{V}^2 a_4 A_4} \int_{\tau_4^{end}}^{\tau_4} \frac{e^{a_4\tau_4}}{\sqrt{\tau_4} (1 - a_4\tau_4)} d\tau_4 \geq 60 \quad (3.5)$$

In our case, the field range for τ_4 during inflation is quite limited: from eq. (3.4) we find

$$2 \ln \mathcal{V} \leq a_4\tau_4 \leq (2 + n) \ln \mathcal{V} \quad (3.6)$$

with $n > 0$. Using this information, the number of e -foldings can be re-expressed as

$$N_e = \frac{-27\beta \hat{\xi} W_0 \gamma_4}{256 \mathcal{V}^2 a_4^{3/2} A_4} \int_{\frac{2}{a_4} \log \mathcal{V}}^{\frac{(2+n)}{a_4} \log \mathcal{V}} \frac{e^y}{\sqrt{y} (1 - y)} dy \geq 60 \quad (3.7)$$

Because we seek the dominant contribution elsewhere, we demand that the inflaton contribution to the power spectrum of curvature perturbations is *much lower* than the amplitude measured by the COBE satellite. This gives the following constraint:

$$\frac{V_{inf}^{3/2}}{M_p^3 V'_{inf}} \ll 5.2 \times 10^{-4}. \quad (3.8)$$

Substituting the potential, and using eq. (3.4) at horizon exit in the limit $a_4\tau_4 \gg 1$, we find

$$\left(\frac{g_s}{8\pi} \right) \frac{3^4 \alpha \gamma_4 (\beta \hat{\xi})^3 W_0^2}{4^6 a_4^{3/2} [(2+n) \ln \mathcal{V}]^{5/2}} \left(\frac{W_0}{A_4} \right)^2 \mathcal{V}^{2(n-1)} \ll 2.7 \times 10^{-7}, \quad (3.9)$$

which uses expression (3.4) for the inflaton field at horizon exit.

A successful model must satisfy both of the constraints (3.7) and (3.9). This imposes conditions on some of the parameters of the model, which must be supplemented by the constraints derived in the following sections coming from the successful realization of the curvaton mechanism. We discuss in §6 explicit scenarios that satisfy all the conditions to have a successful curvaton model.

3.2 Dynamics of the curvaton field χ_1

The curvaton field,⁸ χ_1 , is lighter than the Hubble parameter during inflation, since at large volume

$$m_{\chi_1}^2 \simeq \frac{g_s C_t W_0^2}{4\pi \mathcal{V}^{\frac{10}{3}}} \ll H_\star^2 \simeq \frac{3g_s \beta \hat{\xi} W_0^2}{32\pi \mathcal{V}^3}, \quad (3.10)$$

where the ‘ \star ’ indicates a quantity evaluated at horizon exit. During inflation the field χ_1 slowly rolls classically towards its minimum at zero, but because it is so light it also undergoes quantum fluctuations that in some circumstances can dominate the classical motion.

We now estimate when fluctuations dominate, following [37]. In one Hubble time H_\star^{-1} , the light field χ_1 can fluctuate by an amount $\delta\chi_1 \sim H_\star/2\pi$. On the other hand, during the same time interval a classical slow roll would change the field value by $\Delta\chi_1 \sim -V'_{cur}/(3H_\star)\Delta t_\star = -V'_{cur}/(3H_\star^2)$. Fluctuations dominate classical evolution⁹ when $\delta\chi_1 \sim \Delta\chi_1$, which occurs when $\chi_1 = \chi_Q$, given by

$$V'_{cur}(\chi_Q) \simeq H_\star^3. \quad (3.11)$$

During inflation quantum fluctuations cause the field χ_1 to lie in the interval $0 \leq \chi_1 \leq \chi_Q$ with uniform probability. Then, its typical value is of order $\chi_1 \sim \chi_Q$.

In the present case, approximating the curvaton potential as quadratic, as in eq. (2.43), one finds

$$\chi_Q \simeq \left(\frac{g_s}{8\pi}\right)^{1/2} \frac{(\beta\hat{\xi})^{3/2} W_0}{3C_t \mathcal{V}^{7/6}}, \quad (3.12)$$

and so $\chi_Q \gg H_\star$. But because χ_Q is suppressed by $1/\mathcal{V}^{7/6}$ these fluctuations are nevertheless very small at large volume. A posteriori, it is these powers of $1/\mathcal{V}$ that justify the expansion of the curvaton potential up to second order in χ_1 .

We now estimate in more detail the amplitude of the power spectrum for the curvaton fluctuations, following [39]. The classical evolution equation for the curvaton field is

$$\ddot{\chi}_1 + 3H\dot{\chi}_1 + V'_{cur} = 0. \quad (3.13)$$

Making the first order expansion $\delta V'_{cur} \simeq V''_{cur} \delta\chi_1$, one finds the following equation for the inhomogeneous curvaton fluctuation at superhorizon scales

$$\delta\ddot{\chi}_1 + 3H\delta\dot{\chi}_1 + V''_{cur} \delta\chi_1 = 0. \quad (3.14)$$

Since, for a quadratic potential, $\delta\chi_1$ and χ_1 satisfy the same equation, their ratio does not evolve in time. This means that this ratio keeps the same value it has at horizon exit:

$$\left(\frac{\delta\chi_1}{\chi_1}\right) = \left(\frac{\delta\chi_1}{\chi_1}\right)_\star. \quad (3.15)$$

⁸From now on we drop the *hat* from the field $\hat{\chi}_1$ parameterizing the displacement from the minimum, in eq. (2.33).

⁹This estimate has been debated in the literature, in particular the value of the power of H in the right hand side of (3.11). See for example [38]. In this work, we follow the prescription of [37], but our approach can be adapted to different possibilities. We thank Sami Nurmi for discussions on this point.

Then the power spectrum of fractional field perturbations reads

$$\mathcal{P}_{\delta\chi_1/\chi_1}^{1/2} = \frac{H_\star}{2\pi\chi_\star} \simeq \frac{4C_t}{3\pi\beta\hat{\xi}\mathcal{V}^{1/3}} \quad (3.16)$$

where in the last approximate equality we suppose $\chi_\star \simeq \chi_Q$ (see the previous discussion).

In the next section we discuss how to convert these isocurvature fluctuations into adiabatic curvature fluctuations when the curvaton decays after inflation and reheating have already taken place.

4 Moduli Couplings to Visible Sector Fields

An important feature of the LV framework is that it is possible to directly compute the couplings between the moduli (among which the inflaton and the curvaton) and all the other visible or hidden d.o.f. localized on $D7$ -branes wrapped on internal 4-cycles [22, 27, 30]. This is a necessary ingredient for calculating the inflaton and curvaton decay rates into visible d.o.f. allowing us to understand reheating at the end of inflation [22], and to determine whether a curvaton mechanism can be successfully developed.

In the case of the curvaton, we have to focus only on its decay rate to visible gauge bosons. In fact χ_1 is so light that it cannot decay to any supersymmetric particle or even to the Higgs since this receives a large SUSY breaking contribution to its mass. Thus χ_1 can only decay to gauge bosons g and fermions ψ which are massless before the EW phase transition. However it has been shown in [22] that since the fermions are massless, there is no direct decay $\chi_1 \rightarrow \psi\bar{\psi}$, but only a 3-body decay $\chi_1 \rightarrow \psi\bar{\psi}g$ which is suppressed with respect to the 2-body decay $\chi_1 \rightarrow gg$ by a phase space factor. In addition χ_1 cannot decay to light hidden d.o.f. since the requirement of a viable reheating forces to have for each hidden sector a pure $N = 1$ SYM theory that develops a mass gap [22].

In order to analyse the coupling of moduli to the gauge bosons of the field theory living on a stack of $D7$ -branes, we proceed as follows. The $D7$ s of interest wrap a 4-cycle whose volume is given by τ (which can be any of our moduli): the couplings with the moduli can be worked out from the moduli dependence of the tree-level gauge kinetic function $4\pi/g^2 = \tau$ (see [27]). In full generality, the kinetic terms read:

$$\mathcal{L}_{gauge} = -\frac{\tau}{M_p} F_{\mu\nu} F^{\mu\nu}, \quad (4.1)$$

and we must expand τ around its minimum $\tau \rightarrow \langle\tau\rangle + \hat{\tau}$, and go to the canonically normalized field strength $G_{\mu\nu}$ defined by

$$G_{\mu\nu} = 2\sqrt{\langle\tau\rangle} F_{\mu\nu}. \quad (4.2)$$

Doing so, we obtain:

$$\mathcal{L}_{gauge} = -\frac{1}{4} G_{\mu\nu} G^{\mu\nu} - \frac{\hat{\tau}}{4M_p\langle\tau\rangle} G_{\mu\nu} G^{\mu\nu}. \quad (4.3)$$

4.1 First scenario

As explained in section 2.1 we imagine the observable sector to be localized on a stack of $D7$ -branes wrapped on the τ_1 and τ_2 cycle, and analyze the decay of the inflaton and curvaton fields into visible gauge bosons. This set-up has the following advantages:

1. It represents the simplest example of multi-field curvaton scenario with the smallest number of Kähler moduli which is 4;
2. The geometric localization of the visible sector on τ_1 maximizes the strength of the coupling of the curvaton to visible gauge bosons. As we shall see in section 5, this will yield the largest amount of nongaussianities.

However there are also some shortcomings:

1. The K3 fiber is not a rigid cycle and so one has to worry about how to fix the $D7$ -brane deformation moduli that would give rise to unwanted matter in the adjoint representation. Here we shall assume that these moduli can be fixed by the use of background fluxes.
2. There is a constraint on the volume of τ_1 coming from constraints on the size that is expected for the observed gauge coupling. Denoting the gauge coupling as g , using eq. (2.30), we have

$$\frac{4\pi}{g^2} = \tau_1 \simeq \left(\frac{4A\mathcal{V}}{B} \right)^{2/3}. \quad (4.4)$$

Focusing for definiteness on a GUT theory,¹⁰ $4\pi/g^2 \simeq 25$, we find constraints on the parameters that characterize the string loop contributions. Indeed, the previous relation implies

$$\frac{4A}{B} = \frac{125}{\mathcal{V}} \quad (4.5)$$

from which, using the definitions of A and B in eqs (2.27), (2.28), we obtain

$$(C_1^{KK})^2 = \frac{125\alpha}{g_s^2} \frac{C_{12}^W}{\mathcal{V}} \quad (4.6)$$

As we see in the following, when discussing explicit examples, this condition is relatively easy to satisfy. We do not have to choose unnaturally large hierarchies between the parameters C_1^{KK} and C_{12}^W .

As studied in [43], at the end of inflation, due to the steepness of the potential, the inflaton τ_4 stops oscillating just after two or three oscillations due to an extremely efficient non-perturbative particle production of τ_4 fluctuations. Expanding the canonical normalization (2.22) around the global minimum ($\tau_i = \langle \tau_i \rangle + \hat{\tau}_i \forall i$) we find [22]:¹¹

$$\hat{\tau}_4 \sim \mathcal{O}(\mathcal{V}^{-1/3})\hat{\chi}_1 + \mathcal{O}(1)\hat{\chi}_2 + \mathcal{O}(\mathcal{V}^{-1/2})\hat{\phi}_3 + \mathcal{O}(\mathcal{V}^{1/2})\hat{\phi}_4, \quad (4.7)$$

¹⁰Assuming that the gauge bosons on τ_2 decouple from the EFT getting an $\mathcal{O}(M_s)$ mass.

¹¹The subleading dependence on $\hat{\chi}_1$ is introduced once string loop corrections are included.

realising that the Universe is mostly filled with $\hat{\phi}_4$ -particles plus some $\hat{\chi}_2$ and fewer $\hat{\chi}_1$ and $\hat{\phi}_3$ -particles. Therefore the energy density of the Universe is dominated by $\hat{\phi}_4$ whose decay to visible d.o.f. is responsible for reheating.

The following table summarizes the moduli couplings to visible gauge bosons living on τ_1 (denoting the corresponding field strength as $F_{\mu\nu}^{(1)}$):

| | $\hat{\chi}_1$ | $\hat{\chi}_2$ | $\hat{\phi}_i, \forall i = 3, 4$ |
|----------------------------------|--------------------------|------------------------------------|--|
| $F_{\mu\nu}^{(1)} F^{(1)\mu\nu}$ | $\frac{2}{\sqrt{3} M_p}$ | $\sqrt{\frac{2}{3}} \frac{1}{M_p}$ | $\frac{3(\ln \mathcal{V})^{\frac{3}{4}}}{2 a_i \mathcal{V}^{1/2} M_p}$ |

Table 1: Couplings between moduli and gauge bosons for a field theory on the τ_1 cycle.

Because the light curvaton field mixes through its kinetic terms with both τ_1 and \mathcal{V} , one might hope to use the χ_1 -dependence of couplings and masses to use the modulation mechanism [18, 16] to generate the primordial fluctuations. Although in the present instance the couplings do not depend on the fluctuations of χ_1 , the masses of the fields do. However, it turns out that in all cases we investigated the amplitude of modulation-generated fluctuations is too small to have interesting cosmological consequences. It is for this reason that we focus on the curvaton mechanism in the following.

We can now derive the total decay rate of a generic modulus φ into gauge bosons g :

$$\Gamma_{\varphi \rightarrow gg} = \frac{N_g \lambda^2 m_\varphi^3}{64\pi}, \quad (4.8)$$

where λ is the coupling listed in Table 1 and N_g is the total number of gauge bosons: for definiteness we choose $N_g = 12$ as in the MSSM. We obtain, for our set of fields,

$$\Gamma_{\hat{\chi}_1 \rightarrow gg} = \frac{1}{4\pi} \frac{m_{\chi_1}^3}{M_p^2} \simeq \frac{M_p}{\mathcal{V}^5}, \quad (4.9)$$

$$\Gamma_{\hat{\chi}_2 \rightarrow gg} = \frac{1}{8\pi} \frac{m_{\chi_2}^3}{M_p^2} \simeq \frac{M_p}{\mathcal{V}^{9/2}}, \quad (4.10)$$

$$\Gamma_{\hat{\phi}_j \rightarrow gg} = \frac{27(\ln \mathcal{V})^{\frac{3}{2}}}{64\pi} \frac{m_{\phi_j}^3}{\mathcal{V} M_p^2} \simeq \frac{M_p}{\mathcal{V}^4}, \quad \forall j = 3, 4. \quad (4.11)$$

where we have emphasized, in the extreme right, the dominant volume dependence. Notice that the curvaton decay rate is suppressed with respect to the inflaton decay rate, in the limit of large volume. This observation plays an important role in the viability of the mechanism. The reheating temperature in the approximation of sudden thermalization turns out to be [22]:

$$T_{RH} \simeq \left(\Gamma_{\hat{\phi}_4 \rightarrow gg} M_p \right)^{1/2} \simeq \frac{M_p}{\mathcal{V}^2}. \quad (4.12)$$

4.2 Second scenario

In this section we briefly present a different brane set-up with the visible sector localized on a small blow-up cycle, showing that it is possible to build a curvaton scenario with a standard

realization of the visible sector on a rigid del-Pezzo 4-cycle without any constraint on the overall volume to keep the gauge coupling from getting too small (given that the VEV of blow-up moduli does not depend on \mathcal{V}). However, the blow-up mode supporting the visible sector cannot be either τ_3 or τ_4 due to the tension between chirality and non-perturbative effects [33]. Hence we need to introduce a fifth modulus τ_5 with the following three possible brane set-ups [22]:

1. Visible sector built with a stack of $D7$ -branes wrapped around τ_5 which is stabilized in the geometric regime (for example by string loop effects as in [28]). In this case the inflaton τ_4 is not wrapped by the visible sector. The inflaton and curvaton total decay rates to gauge bosons scale as:

$$\Gamma_{\hat{\chi}_1 \rightarrow gg} \simeq \frac{M_p}{\mathcal{V}^{17/3}}, \quad \Gamma_{\hat{\phi}_4 \rightarrow gg} \simeq \frac{M_p}{\mathcal{V}^4}, \quad \Rightarrow \quad T_{RH} \simeq \frac{M_p}{\mathcal{V}^2}. \quad (4.13)$$

2. Visible sector built with a stack of $D7$ -branes wrapped around a combination of τ_4 and τ_5 with chiral intersections only on τ_5 so that the non-perturbative corrections in τ_4 are not destroyed. In this case the inflaton τ_4 is wrapped by the visible sector. The inflaton and curvaton total decay rates to gauge bosons scale as:

$$\Gamma_{\hat{\chi}_1 \rightarrow gg} \simeq \frac{M_p}{\mathcal{V}^{17/3}}, \quad \Gamma_{\hat{\phi}_4 \rightarrow gg} \simeq \frac{M_p}{\mathcal{V}^2}, \quad \Rightarrow \quad T_{RH} \simeq \frac{M_p}{\mathcal{V}}. \quad (4.14)$$

3. Visible sector built via fractional branes at the quiver locus $\tau_5 \rightarrow 0$ (τ_5 can shrink to zero size by D -terms as in [35]). The inflaton and curvaton total decay rates to gauge bosons scale as:

$$\Gamma_{\hat{\chi}_1 \rightarrow gg} \simeq \frac{M_p}{\mathcal{V}^{20/3}}, \quad \Gamma_{\hat{\phi}_4 \rightarrow gg} \simeq \frac{M_p}{\mathcal{V}^5}, \quad \Rightarrow \quad T_{RH} \simeq \frac{M_p}{\mathcal{V}^{5/2}}. \quad (4.15)$$

It is interesting to notice that, due to the geometric separation between τ_1 and τ_5 , the coupling of the curvaton to visible gauge bosons is weaker than in the first scenario. This yields a lower level of nongaussianities since, as we shall see in section 5, $f_{NL} \propto \Gamma_{\chi_1}^{1/2}$.

5 Dynamics after inflation: the curvaton mechanism

In this section, we summarize the curvaton mechanism for converting isocurvature fluctuations into adiabatic, curvature fluctuations in the above inflationary model. Moreover, we estimate the resulting level of nongaussianity produced in this process focusing on the first scenario with the visible sector localized on τ_1 . However it is easy to re-formulate our analysis for the second scenario.

5.1 Amplitude of adiabatic fluctuations

In our scenario, at the end of inflation the energy density of the Universe is dominated by the energy in the inflaton field, ϕ_4 . Because of the \mathcal{V} dependence of the decay rates found above, the small moduli, ϕ_i , have the largest decay rate – see eq. (4.11) – and so these moduli are also

the first of the moduli to decay. This decay converts the inflaton energy density into radiation, after which its energy density falls with the scale factor like $\rho_\gamma \propto a^{-4}$. Since this is the dominant component of the energy, after this point the Hubble parameter falls like a^{-2} .

Energy tied up in the curvaton field, on the other hand, need not fall this fast. For instance, once H falls below the curvaton field's mass this field starts to oscillate coherently around its minimum, during which its energy density scales like non-relativistic matter: $\rho_{\chi_1} \sim a^{-3}$. Since this is much slower than the energy density of radiation the relative proportion of curvaton energy to radiation energy can grow while the curvaton oscillates.

This continues until the curvaton field starts to decay, which happens once the Hubble parameter becomes comparable to the curvaton decay rate (that, recall, in our set-up is the most suppressed: see eq. (4.9)). At this point the curvaton energy density also converts into radiation, bringing with it any isocurvature fluctuations that had been stored in the curvaton field. This converts the curvaton fluctuations into the adiabatic fluctuations of the radiation energy density.

The total size of the adiabatic fluctuations inherited by such a conversion depends on the size of the curvaton energy density relative to the radiation at the point where the curvaton decays. Denoting this fraction by $\Omega = \rho_{cur}/\rho_\gamma$, then in a sudden decay approximation, we find:

$$\Omega \simeq \left[\frac{1}{6} \left(\frac{\chi_\star}{M_p} \right)^2 \left(\frac{m_\chi}{\Gamma_{\chi_1}} \right)^{\frac{1}{2}} \right]^a \simeq \left[\left(\frac{3}{32} \right)^2 \frac{g_s^{1/2} (\beta \hat{\xi})^3 W_0}{C_t^{5/2} \mathcal{V}^{2/3}} \right]^a, \quad (5.1)$$

with:

$$\begin{cases} a = 1 & \text{for radiation dominance} & \Leftrightarrow & \Omega \ll 1, \\ a = 4/3 & \text{for curvaton dominance} & \Leftrightarrow & \Omega \gg 1. \end{cases} \quad (5.2)$$

The last equality in (5.1) substitutes the value of the various quantities in the present scenario. The resulting expression for the power spectrum of curvature fluctuations, in the limit¹² in which $\Omega \ll 1$, is [39]:

$$\mathcal{P}_\zeta^{\frac{1}{2}} = \frac{2}{3} \Omega \mathcal{P}_{\delta\chi_1/\chi_1}^{\frac{1}{2}} \simeq \frac{1}{128\pi} \frac{g_s^{1/2} (\beta \hat{\xi})^2 W_0}{C_t^{3/2} \mathcal{V}}. \quad (5.3)$$

Demanding this converted amplitude agree with the amplitude measured by COBE then gives $\mathcal{P}_\zeta^{\frac{1}{2}} = 4.8 \times 10^{-5}$, which imposes the constraint

$$C_t^{\frac{3}{2}} \simeq 50 \frac{g_s^{1/2} (\beta \hat{\xi})^2 W_0}{\mathcal{V}}. \quad (5.4)$$

5.2 Nongaussianities

Following [39], it is not difficult to provide an estimate for the amount of nongaussianity predicted in this scenario. We focus on nongaussianities of local form

$$\zeta = \zeta_G + \frac{3}{5} f_{NL} \zeta_G^2, \quad (5.5)$$

¹²It is easy to re-express each quantity in the curvaton dominance case.

where ζ_G is a Gaussian curvature fluctuation. This Ansatz is particularly well-suited to the present context, since there is a non-linear relation between scalar fluctuations, and curvature perturbations produced after inflation ends. In writing eq (5.3), indeed, we implicitly express the curvature fluctuation as a first order expansion in the fluctuation of χ_1 . The complete expression, generalizing the linear order relation given in eq. (5.3), allows to exhibit the non-linear connection between scalar and curvature fluctuations. Indeed, it reads

$$\zeta = \frac{\Omega}{3} \frac{\delta\rho_{\chi_1}}{\rho_{\chi_1}}. \quad (5.6)$$

In our case, since we work with a quadratic potential, one finds that

$$\frac{\delta\rho_{\chi_1}}{\rho_{\chi_1}} = 2 \frac{\delta\chi_1}{\chi_1} + \frac{(\delta\chi_1)^2}{\chi_1^2} \quad (5.7)$$

Consequently, including this second order expansion in the definition of ζ of eq. (5.6), and comparing with the Ansatz in (5.5), one can read the following expression for f_{NL} :

$$f_{NL} = \frac{5}{4\Omega} \simeq \frac{140 C_t^{5/2} g_s^4 \mathcal{V}^{2/3}}{W_0 \beta^3 \xi^3} = 10^5 \frac{(\beta\xi W_0^2)^{1/3}}{g_s^{1/6} \mathcal{V}}, \quad (5.8)$$

where in the last step we use relation (5.4). This expression quantifies the amount of nongaussianity in this set-up. Notice that the size of f_{NL} is inversely proportional to the conversion factor Ω . This is expected: if we decrease the efficiency of the conversion process, by decreasing Ω , we have at the same time to increase the ration $\delta\chi_1/\chi_1$ in order to account for the observed amplitude of fluctuations (see eqs. (5.3) and (5.6)-(5.7)). But in this case, the quadratic contribution in $\delta\chi_1/\chi_1$, in formula eq. (5.7), becomes important in comparison with the linear term, implying an increase of nongaussianity.

It is also possible to analyse nongaussianity beyond the parameter f_{NL} , for example discussing the parameters τ_{NL} and g_{NL} that characterize the trispectrum. Expressions for these parameters, in curvaton scenarios, have been provided in the literature: see for example [40] for a recent review. For our curvaton model, with quadratic potential, small decay rate Ω and in the sudden decay approximation, one finds

$$\tau_{NL} = \frac{36}{25} f_{NL}^2 \quad , \quad g_{NL} \simeq -\frac{10}{3} f_{NL} \quad (5.9)$$

with f_{NL} given in eq. (5.8). The expression for τ_{NL} is the typical one for models where only one species contributes to the generation of curvature perturbations. The value of g_{NL} , being proportional to f_{NL} , turns out to be too low for being detectable by Planck, given the already stringent bounds on f_{NL} from WMAP7 [41]. It would be interesting to extend the model above such as to find set-ups in which τ_{NL} and g_{NL} turn out to be large, or in which one obtains a sizeable running of nongaussianity, as analysed in [42].

5.3 Constraints from Big-Bang nucleosynthesis

Besides the requirements of providing the correct amplitude for curvature perturbations, Big-Bang nucleosynthesis (BBN) imposes further constraints on the curvaton model. This since we must ensure that the curvaton field decays by the time BBN takes place, at around $T_{BBN} \sim 1$ MeV. In order to satisfy this constraint, we impose the following inequality

$$\Gamma_{\hat{\chi}_1 \rightarrow gg} > H_{BBN} \sim 10^{-24} \text{ GeV} \quad (5.10)$$

Using the expression for Γ given in (4.9), we obtain

$$\Gamma_{\hat{\chi}_1 \rightarrow gg} \simeq \left(\frac{g_s^{3/2} C_t^{3/2} W_0^3}{32\pi^{5/2}} \right) \frac{M_p}{\mathcal{V}^5} \quad (5.11)$$

Now, recalling that $C_t \simeq 2B q^{-\frac{1}{3}}$ within the approximation we are considering, we get an upper bound on the volume,

$$\mathcal{V} < 5.5 \times 10^7 \left(g_s^{3/2} B^2 A^{-1/2} W_0^3 \right)^{1/5}. \quad (5.12)$$

For standard values of the parameters, this imposes a bound on the volume of order $\mathcal{V} \leq 10^8$.

6 Explicit set-ups

The previous sections present the conditions that our system must satisfy in order to furnish a realization of a curvaton scenario. In this section, we present two representative parameter choices that satisfy all the constraints, to get a preliminary sense of how much observable quantities vary.

There is a simple first observation. The results of the previous sections suggest that once volumes are too large (and so the inflationary Hubble scale becomes too low) then it becomes difficult to obtain adequately large primordial fluctuations using the curvaton mechanism. Indeed, eq. (5.3) cannot be satisfied for volumes that are too large without requiring other parameters to acquire unnatural values. For typical values of the parameters a curvaton scenario has a chance for volumes in the range $10^3 \leq \mathcal{V} \leq 10^8$. Also, eq. (5.8) shows that very large volumes are usually associated with nongaussianities of small size. Obtaining a large f_{NL} is therefore easiest when choosing relatively small volumes. Because the underlying expansion is in powers of $\alpha'/\ell_s^2 \propto 1/\mathcal{V}^{1/3}$ we never allow ourselves to consider volumes smaller than $\mathcal{V}_{\min} \simeq 10^3$.

6.1 First example: small volume, large f_{NL}

Consider the following representative choice of parameters:

| \mathcal{V} | a_4 | ξ | g_s | n | W_0 | α | A_4 | γ_4 |
|---------------|---------------|-------|-------|----------------|-------|----------|-------|------------|
| 10^3 | $\frac{1}{2}$ | 0.1 | 0.3 | $\frac{1}{16}$ | 1 | 0.1 | 1 | 0.1 |

This example is characterized by not-too-large a volume, $\mathcal{V} = 10^3$ in Planck units, and by a relatively large string coupling, $g_s \simeq 0.3$. Also $a_4 = 1/2$ corresponds to a gauge group with

large rank in the non-perturbative contribution to the inflaton superpotential. Plugging these parameters in eqs. (3.7) and (3.9), we find a large number of e-foldings in this model ($N_e \simeq 170$) and a very small inflaton contribution to the amplitude of adiabatic fluctuations ($\mathcal{P}_\zeta^{inf} \simeq 10^{-8} \mathcal{P}_\zeta^{COBE}$). Since the volume is relatively small, the scale of inflation is fairly high in this example. Next, the conditions of having an acceptable size for the gauge coupling theory, discussed in section 4, imposes the condition $C_{12}^W = 10 (C_1^{KK})^2$, which in turn implies $C_t \simeq C_{12}^W$. The COBE normalization condition for the curvaton fluctuations (5.4) then fixes $C_t \sim 0.05$.

The most important feature of this model is the high level of nongaussianity it predicts: using the previous results we find

$$f_{NL} \simeq 56. \quad (6.1)$$

This value can be slightly changed by tuning the choice of parameters, but the requirement of satisfying all the constraints does not leave much freedom in this regard. Consequently the order of magnitude for f_{NL} is fairly robust in this scenario with not too large volume ($\mathcal{V} = 10^3$) and high rank gauge group ($a_4 = 1/2$).

6.2 Second example: larger volume, smaller f_{NL}

Choosing a different set of parameters shows how the results change as the volume grows. Consider the following choice

| \mathcal{V} | a_4 | ξ | g_s | n | W_0 | α | A_4 | γ_4 |
|---------------|-------|-------|-----------|----------------|-------|-----------|-------|------------|
| 10^6 | 0.85 | 1 | 10^{-2} | $\frac{1}{16}$ | 40 | 10^{-2} | 1 | 10^{-2} |

In this example, the volume is larger with respect to the previous example, the string coupling small, and $a_4 \sim 1$. Plugging these parameters in eqs. (3.7) and (3.9), we find a sufficient number of e-foldings in this model ($N_e \simeq 67$) and a small contribution of the inflaton sector to the COBE amplitude of adiabatic fluctuations ($\mathcal{P}_\zeta^{inf} \simeq 10^{-2} \mathcal{P}_\zeta^{COBE}$). After requiring to have an acceptable gauge coupling, as discussed in section 4, and imposing COBE normalization condition (5.4), we find that $C_t \sim 34$. The amount of nongaussianity in this case is small:

$$f_{NL} \simeq 2.5, \quad (6.2)$$

showing that the value of f_{NL} strongly depends on the choice of underlying parameters. Different models characterized by different volumes, although providing the same amplitude for the spectrum of adiabatic fluctuations, nevertheless give very different values for f_{NL} .

7 Conclusions

In this paper we use LARGE Volume string compactifications to construct a controlled string-inflation model that does not use the inflaton to generate primordial fluctuations. Because the

dynamics cannot be captured by a simple single-field slow roll, it becomes possible to generate observably large non-gaussianities. These tend to have the local form in the model examined because they are generated well after inflation ends.

The key ingredients for any such a scenario are twofold. There must be other fields, besides the inflaton, with masses $m \ll H$ during the inflationary epoch in order to have isocurvature fluctuations be generated over extra-Hubble distances. The second ingredient is a mechanism for converting these isocurvature fluctuations into adiabatic fluctuations.

We find that both ingredients are possible in the LV scenario. The hierarchy of volume-suppressed modulus masses enjoyed by this scenario allows some moduli to have masses that are parametrically suppressed relative to the Hubble scale during inflation, thereby providing a source of isocurvature fluctuations.

These states also plausibly have a hierarchy of decay rates into ordinary matter, assuming that ordinary matter is localized on a brane that wraps one of the cycles whose moduli appear in the low-energy theory. This allows the isocurvature mode to first accumulate as an overall fraction of the total energy density, by oscillating after the inflaton has decayed to radiation. It can then itself decay at much later times, converting its fluctuations into adiabatic perturbations. The resulting picture provides a realization of the curvaton mechanism for string inflationary models. The fraction of the energy density carried by the curvaton is suppressed by powers of $1/\mathcal{V}$, naturally leading this fraction to be a small (and nongaussianities to be comparatively large – $\mathcal{O}(10)$ – if the amplitude is the one observed).

Ultimately, the reason such a construction is possible is because of the potentially large number of fields that can be cosmologically active during LV inflation. Indeed, should local nongaussianity be observed, this is probably what it would be telling us: the dynamics generating primordial fluctuations likely involves several cosmologically active fields rather than just one.

Because additional light fields are present these models can be expected also to manifest other nonstandard mechanisms for generating fluctuations, such as the modulation mechanism, although we do not yet have explicit working examples of this type. A potential benefit of these kinds of models might be the ability to lower the string scale while still obtaining acceptably large primordial fluctuations, since this makes it easier to have a lower supersymmetry-breaking scale, as seems to be preferred by particle phenomenology in the later universe. As ever, it would be useful to know how common such models might be in the string landscape.

It is worth noticing that even though these scenarios require many moduli to work, this is the generic case in string compactifications. The perspective taken in this article is that simplicity arguments using the minimum number of fields are usually good starting points but may not capture the dynamics of the generic case. Furthermore, contrary to most models of string cosmology, we also consider phenomenological constraints, such as the location of the standard-model brane, the value of the present-day gauge coupling, efficient reheating, and so on. We believe this to be crucial because string theory asks to be more than just a model of inflation:

string scenarios must therefore address all observable issues and not only a subset of them. Even with these constraints, we find it encouraging that non-inflaton generation of primordial perturbations appears possible, consistent with having the right amount of inflation required by later cosmology, agreement with current CMB measurements, with potentially observable features like nongaussianity for future experiments. The imminent start of Planck observations makes these questions timely and worth pursuing.

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Appendix

A String versus Einstein frame

The correct prefactor of the scalar potential in 4D Einstein frame has been explicitly shown in [30]. Given that the Kähler potential that reproduces the kinetic terms for the moduli in 4D Einstein frame is known to be (with $S = e^{-\phi} + iC_0$):

$$\frac{K_E}{M_p^2} = -2 \ln \mathcal{V}_E - \ln(S + \bar{S}) - \ln \left(-i \int \Omega \wedge \bar{\Omega} \right), \quad (\text{A.1})$$

here we shall briefly review just the derivation of the prefactor of the superpotential starting from the 10D type IIB supergravity action in string frame (showing only the relevant terms):

$$S_{10D}^{(s)} \supset \frac{1}{(2\pi)^7 \alpha'^4} \int d^{10}x \sqrt{-g_{10}^{(s)}} \left(e^{-2\phi} \mathcal{R}_{10}^{(s)} - \frac{G_3 \cdot \bar{G}_3}{2 \cdot 3!} \right). \quad (\text{A.2})$$

The action in Einstein frame is obtained via a Weyl rescaling of the metric of the form $g_{MN}^{(s)} = e^{\phi/2} g_{MN}^{(E)}$:

$$S_{10D}^{(E)} \supset \frac{2\pi}{l_s^8} \int d^{10}x \sqrt{-g_{10}^{(E)}} \left(\mathcal{R}_{10}^{(E)} - \frac{G_3 \cdot \bar{G}_3}{12 \text{Re } S} \right), \quad (\text{A.3})$$

where $l_s = 2\pi\sqrt{\alpha'}$. The dimensional reduction of (A.3) from 10D to 4D then yields:

$$S_{4D}^{(E)} \supset \frac{2\pi}{l_s^8} \left(\int d^4x \sqrt{-g_4^{(E)}} \mathcal{R}_4^{(E)} Vol_E - \overbrace{\int d^4x \sqrt{-g_4^{(E)}} \left(\int d^6x \sqrt{g_6^{(E)}} \frac{G_3 \cdot \bar{G}_3}{12 \text{Re } S} \right)}^{V_{flux}} \right), \quad (\text{A.4})$$

where $Vol_E = \int d^6x \sqrt{g_6^{(E)}} \equiv \mathcal{V}_E l_s^6$. Comparing the first term in (A.4) with the Einstein-Hilbert action $S_{EH} = (M_p^2/2) \int d^4x \sqrt{-g^{(E)}} \mathcal{R}^{(E)}$, we find:

$$M_p^2 = \frac{4\pi \mathcal{V}_E}{l_s^2} \quad \text{and} \quad M_s \equiv \frac{1}{l_s} = \frac{M_p}{\sqrt{4\pi \mathcal{V}_E}}. \quad (\text{A.5})$$

Writing the superpotential in 4D Einstein frame as:

$$W_E = \frac{p}{l_s^2} \int G_3 \wedge \Omega, \quad (\text{A.6})$$

the correct prefactor p can be found from requiring that V_{flux} is reproduced by:

$$V = \int d^4x \sqrt{-g_4^{(E)}} e^{K_E/M_p^2} \left[K_E^{i\bar{j}} D_i W_E D_{\bar{j}} \bar{W}_E - \frac{3}{M_p^2} W_E \bar{W}_E \right], \quad (\text{A.7})$$

obtaining $p = M_p^3/(\sqrt{4\pi})$. Therefore, including the leading order α' corrections to K_E and non-perturbative corrections to W_E , the F -term scalar potential in 4D Einstein frame can be derived from:

$$\frac{K_E}{M_p^2} = -2 \ln \left[\mathcal{V}_E + \frac{\xi}{2} \left(\frac{S + \bar{S}}{2} \right)^{3/2} \right] - \ln(S + \bar{S}) - \ln \left(-i \int \Omega \wedge \bar{\Omega} \right), \quad (\text{A.8})$$

$$W_E = \frac{M_p^3}{\sqrt{4\pi}} \left(\frac{1}{l_s^2} \int G_3 \wedge \Omega + \sum_i A_i e^{-a_i T_i^{(E)}} \right). \quad (\text{A.9})$$

Stabilising the dilaton $\langle \text{Re}(S) \rangle = g_s^{-1}$ and the complex structure moduli via background fluxes at tree-level, (A.8) and (A.9) reduce to:

$$\frac{K_E}{M_p^2} = -2 \ln \left(\mathcal{V}_E + \frac{\xi}{2g_s^{3/2}} \right) + \ln \left(\frac{g_s}{2} \right) + K_{cs}, \quad (\text{A.10})$$

$$W_E = \frac{M_p^3}{\sqrt{4\pi}} \left(W_0 + \sum_i A_i e^{-a_i T_i^{(E)}} \right), \quad (\text{A.11})$$

where:

$$K_{cs} = -\ln \left(-i \int \langle \Omega \wedge \bar{\Omega} \rangle \right), \quad \text{and} \quad W_0 = \frac{1}{l_s^2} \int \langle G_3 \wedge \Omega \rangle. \quad (\text{A.12})$$

Hence the prefactor of the scalar potential in 4D Einstein frame can be worked out from:

$$e^{K_E/M_p^2} \frac{|W_E|^2}{M_p^2} \implies \left(\frac{g_s e^{K_{cs}}}{8\pi} \right) M_p^4. \quad (\text{A.13})$$

The expressions for K_s and W_s in 4D string frame can be derived by transforming the scalar potential (recalling that $T_i^{(E)} = T_i^{(s)}/g_s$), and then working out the form of K_s and W_s that reproduce such a potential. We obtain:

$$\frac{K_s}{M_p^2} = -2 \ln \left(\mathcal{V}_s + \frac{\xi}{2} \right) + \ln \left(\frac{g_s}{2} \right) + K_{cs}, \quad (\text{A.14})$$

$$W_s = \frac{g_s^{3/2} M_p^3}{\sqrt{4\pi}} \left(W_0 + \sum_i A_i e^{-a_i T_i^{(s)}/g_s} \right). \quad (\text{A.15})$$

Thus the prefactor of the scalar potential in 4D string frame can be worked out from:

$$e^{K_s/M_p^2} \frac{|W_s|^2}{M_p^2} \implies \left(\frac{g_s^4 e^{K_{cs}}}{8\pi} \right) M_p^4. \quad (\text{A.16})$$

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