We present a curvaton model from type IIB string theory compactified on a warped throat with approximate isometries. Considering an (anti-)D3-brane sitting at the throat tip as a prototype standard model brane, we show that the brane’s position in the isometry directions can play the role of curvatons. The basic picture is that the fluctuations of the (anti-)D3-brane in the angular isometry directions during inflation eventually turns into the primordial curvature perturbations, and subsequently the brane’s oscillation excites other open string modes on the brane and reheat the universe. We find in the explicit case of the KS throat that a wide range of parameters allows a consistent curvaton scenario. It is also shown that the oscillations of branes at throat tips are capable of producing large non-Gaussianity, either through curvature or isocurvature perturbations. Since such setups naturally arise in warped (multi-)throat compactifications and are constrained by observational data, the model can provide tests for compactification scenarios. This work gives an explicit example of string theory providing light fields for generating curvature perturbations. Such mechanisms free the inflaton from being responsible for the perturbations, thus open up new possibilities for inflation models.
Contents

1 Introduction

2 Requirements for the Curvaton

3 Curvatons in Warped Throats
   3.1 Warped Deformed Conifold
   3.2 Effective Action of the Curvaton
   3.3 Consistency Conditions

4 Generating the Cosmological Perturbations
   4.1 Curvaton Scenario
   4.2 Non-Gaussianity through Curvature Perturbations
   4.3 Non-Gaussianity through Isocurvature Perturbations

5 Conclusions

1 Introduction

The inflationary era in the early universe not only sets appropriate initial conditions for the subsequent Hot Big Bang cosmology, but also provides initial inhomogeneities of the universe and seeds the formation of large scale structures. However, since inflation likely takes place at energy scales way beyond our current experimental reach, a detailed understanding of the origin of inflationary cosmology has remained a major theoretical challenge. In order to connect the underlying high-energy physics and current cosmological observations, it is essential to come up with explicit models in a UV-complete framework. This has stimulated numerous attempts to embed inflationary cosmology within string theory which is so far our best candidate for a UV-completion of quantum field theory and gravity. (For a review, see e.g., [4, 5, 6, 7, 8].)

An important feature of string cosmology is its high sensitivity to the physics of string compactification. This imposes stringent restrictions on string inflationary models, and it is nontrivial whether the inflaton can generate the primordial curvature perturbations. In order to avoid decompactification, low-scale inflation is favored, which often leads to the amplitudes of the resulting curvature perturbations being too small compared to the observed COBE normalization value [9]. Recent studies [10, 11, 12, 13, 14] have shown that the inflation scale can be disentangled from the low energy scale of supersymmetry phenomenology, but still the scale of inflation is required to lie within the energy range compatible with stable compactifications. Furthermore, moduli stabilization effects generally introduce steep potentials for the inflaton, which was shown explicitly for the case of warped...
D-brane inflation models [15, 16, 17, 18, 19] in the form of the $\eta$-problem. Although there exist rapid-roll inflationary attractors for such potentials [20, 21, 22, 23], the inflaton itself cannot produce a scale-invariant perturbation spectrum.

Thus, one can expect some field(s) other than the inflaton to have generated the curvature perturbations, as in the case of the curvaton scenario [24, 25, 26, 27]. Realization of such mechanisms can drastically relax the constraints on the inflaton potential, setting inflation models free from the standard slow-roll type [28].

In this paper, we present a simple curvaton model from string theory compactified on a warped throat with approximate isometries. A good example of such throats is the deformed conifold [37] in type IIB string theory, where it has been shown that fluxes and nonperturbative effects can stabilize all its moduli [38, 39]. Considering an (anti-)D3-brane sitting at the throat tip as a prototype standard model brane\(^4\), we show that the brane’s position in the isometry directions can play the role of the curvaton. The basic picture is that the fluctuations of the (anti-)D3-brane in the angular isometry directions during inflation eventually turns into the primordial curvature perturbations, and subsequently the brane’s oscillation excites other open string modes on the brane and reheat our universe\(^5\). We find in the explicit case of the Klebanov-Strassler (KS) throat [37] that a wide range of parameters allows a consistent curvaton scenario.

An important feature of the model is that the generated perturbation spectrum can also have large non-Gaussianity. In this regard, later in this paper, we also consider effects on the perturbation spectrum when the internal manifold contains multiple throats with (anti-)D3-branes at their tips. Even if the branes’ fluctuations in the isometry directions produce only a negligible amount of curvature perturbations, they turn out to be capable of contributing substantially to the non-Gaussianity either through adiabatic or isocurvature perturbations. These features provide tests not only for the curvaton model, but also for multi-throat compactifications.

The paper is organized as follows. In Section 2 we lay out the requirements for a general curvaton scenario to successfully produce the primordial perturbations. In Section 3 we discuss the dynamics of an $\overline{D3}$-brane at the tip of a warped deformed conifold, and show that its oscillation can serve as a curvaton. Then in Section 4 we explore the parameter space and study several cases where the curvaton produces curvature perturbations and/or non-Gaussianity. Finally in Section 5 we conclude and give discussions on future prospects.

\(^3\)Some earlier discussions on curvatons in string theory appeared in [29]. There are also other known mechanisms for generating curvature perturbations, such as modulated reheating [30, 31] and the Lyth effect [32, 33]. Stringy realizations of the latter have been attempted in [34, 35, 36].

\(^4\)For an attempt to realize the standard model on $\overline{D3}$-branes at throat tips, see e.g. [40].

\(^5\)Reheating the universe from D-brane oscillation was also considered in [41, 42], however, the details differ from what we are going to discuss in this paper.
2 Requirements for the Curvaton

A curvaton \([24, 25, 26, 27, 43]\) is a light field, hence during inflation it acquires fluctuations that are nearly scale invariant and Gaussian. Although its energy density is initially negligible, after inflation the curvaton starts oscillating around its potential minimum and behaves like nonrelativistic matter. As its energy density grows relative to radiation, its fluctuations increasingly contribute to the curvature perturbation until the curvaton decays or dominates the universe. Let us begin by laying out the requirements for a consistent curvaton scenario.

We assume the inflation energy to be converted to radiation right after inflation ends (a stringy realization of such a setup is discussed in Section [4]), and also the curvaton to be out of equilibrium until it decays into SM particles, or particles that eventually turn into SM particles. We denote the curvaton mass as \(m\), its decay rate \(\Gamma\), and its field value during inflation \(\sigma^*\). When the curvaton potential is quadratic, then its energy density is negligible during inflation if

\[
\frac{m^2 \sigma^*}{V_{\text{inf}}} \sim \frac{m^2 \sigma^*_*}{M_{\text{Pl}}^2 H_{\text{inf}}^2} \ll 1, \tag{2.1}
\]

where \(V_{\text{inf}}\) and \(H_{\text{inf}}\) are respectively the inflation energy scale and Hubble parameter, and \(M_{\text{Pl}}\) is the reduced Planck mass. We also require the curvaton mass to be sufficiently smaller than the Hubble parameter,

\[
\frac{m^2}{H_{\text{inf}}^2} \ll 1. \tag{2.2}
\]

Setting the cosmic time \(t\) by \(\rho \sim M_{\text{Pl}}^2 / t^2\), then inflation ends at \(t_{\text{end}} \sim H_{\text{inf}}^{-1}\). Due to the Hubble damping, the curvaton remains at \(\sigma^*\) until the Hubble parameter drops below its mass at \(t_{\text{osc}} \sim m^{-1}\) and then the curvaton starts to oscillate. If the curvaton’s lifetime is long enough, then it dominates the energy density of the universe at time \(t_{\text{dom}}\) satisfying

\[
V_{\text{inf}} (t_{\text{end}}/t_{\text{dom}})^2 \sim m^2 \sigma^* (t_{\text{osc}}/t_{\text{dom}})^{3/2},
\]

which gives

\[
t_{\text{dom}} \sim \frac{M_{\text{Pl}}^4}{m \sigma^*_*}. \tag{2.3}
\]

This will be after the onset of the curvaton oscillation if and only if

\[
\frac{\sigma^*}{M_{\text{Pl}}} < 1 \tag{2.4}
\]

is satisfied. Later in (3.28) we will see that the condition (2.4) is automatically satisfied in our model. It should also be noted that (2.1) is trivially satisfied under (2.2) and (2.4).

We further require the curvaton decay at \(t_{\text{dec}} \sim \Gamma^{-1}\) to be after \(t_{\text{osc}}\), hence we impose

\[
\frac{\Gamma}{m} < 1. \tag{2.5}
\]

This requirement turns out to be inevitable for our model, when we introduce additional microscopic constraints in Subsection [3.3](cf. (3.34)). If the curvaton contributes a significant
fraction of the total energy density, then it should decay before the Big Bang Nucleosynthesis (BBN) at \( \rho_{\text{BBN}} \sim (1 \text{ MeV})^4 \sim 10^{-84} M_{\text{Pl}}^4 \), which requires
\[
10^{42} \frac{\Gamma}{M_{\text{Pl}}} > 1. \tag{2.6}
\]

The curvature perturbation produced by the curvaton stops growing when the curvaton decays or dominates the universe. For the generated perturbation to be nearly Gaussian, the field perturbation of the curvaton should be smaller than the unperturbed value,
\[
\frac{H_{\text{inf}}}{\sigma_*} < 1. \tag{2.7}
\]
This requirement is also inevitable for our model, as we will see in \( \text{(3.31)} \). Then, the final curvature perturbation is given by
\[
\zeta \sim r \frac{\delta \rho_{\sigma}}{\rho_{\sigma}} \sim r \left[ 2 \frac{\delta \sigma_*}{\sigma_*} + \left( \frac{\delta \sigma_*}{\sigma_*} \right)^2 \right]. \tag{2.8}
\]
Here, \( r \) is the energy density fraction of the curvaton at the decay epoch, which one can estimate as
\[
\begin{align*}
    r &\equiv \frac{\rho_{\sigma}}{\rho} \bigg|_{\text{dec}} \\
    &\sim \min \left\{ \left( \frac{t_{\text{dec}}}{t_{\text{dom}}} \right)^{1/2}, 1 \right\} \\
    &\sim \min \left\{ \frac{m^{1/2} \sigma_*^2}{\Gamma^{1/2} M_{\text{Pl}}^2}, 1 \right\}. \tag{2.9}
\end{align*}
\]
Therefore the power spectrum of the generated curvature perturbation is
\[
P_{\zeta}^{1/2} \sim r \frac{H_{\text{inf}}}{\sigma_*} \lesssim P_{\text{COBE}}^{1/2} \sim 10^{-5}, \tag{2.10}
\]
where the upper bound on the amplitude is set by the COBE normalization \cite{9}. (If the curvaton is the only source for the perturbations, then the inequality should be saturated.) The second term in the far right hand side of \( \text{(2.8)} \) contributes to non-Gaussian perturbations. In terms of the non-linearity parameter \( f_{\text{NL}} \) \cite{14}, the non-Gaussianity in the curvature perturbation\footnote{We discuss non-Gaussianity produced through isocurvature perturbation in Subsection 4.3} can be estimated as
\[
f_{\text{NL}} \sim \frac{r^3}{P_{\text{COBE}}^2} \left( \frac{H_{\text{inf}}}{\sigma_*} \right)^4 \lesssim 100, \tag{2.11}
\]
where we have also presented the observational upper bound from WMAP \cite{9}. It is clear from \( \text{(2.10)} \) and \( \text{(2.11)} \) that if the curvaton dominates the universe before decay, i.e. \( r \sim 1 \), then the produced \( f_{\text{NL}} \) is at most of order one.

\footnote{We have assumed that the fluctuation of the curvaton is not correlated with that of any other field which generates curvature perturbations.}
3 Curvatons in Warped Throats

We now show an explicit realization of the curvaton in the context of warped type IIB compactifications. We consider the six-dimensional internal space to be compactified to a (conformally) Calabi-Yau (CY) space which includes warped deformed conifold throat regions. Assuming that we are living on an $\mathbb{D}^3$-brane sitting at the tip of the conifold throat, the angular position of the $\mathbb{D}^3$ in the isometry directions plays the role of curvatons. (The reason why we choose an $\mathbb{D}^3$ and not a $\mathbb{D}^3$ will be given at the end of Subsection 3.1.)

However, since the throat is glued to the bulk CY which in general do not preserve such isometries, the throat isometries are broken. Also, nonperturbative effects that stabilize the Kähler moduli confine the $\mathbb{D}^3$ to certain loci on the tip. The warping of the throat suppresses all such effects at the tip, and consequently the angular position of the $\mathbb{D}^3$ receives small mass and couplings to other open string modes on the brane.

3.1 Warped Deformed Conifold

Let us begin by discussing the geometry of the warped deformed conifold. Away from the tip, the deformed conifold \cite{45} is well approximated by a cone whose five-dimensional base space is $S^2 \times S^3$, with the isometry group $SU(2) \times SU(2) \times U(1)$. As one approaches the tip, the $S^2$ shrinks to zero size, while the $S^3$ remains finite.

The supergravity solution for the warped deformed conifold was found by KS \cite{37}, and is also included in the class of flux compactifications considered by Giddings, Kachru, and Polchinski (GKP) \cite{38}. There it was shown that the KS solution can be constructed by turning on $M$ units of R-R flux $F_3$ on the $S^3$ at the tip, and also $-K$ units of NS-NS flux $H_3$ on the dual cycle,

\begin{equation}
\frac{1}{2\pi\alpha'} \int_A F_3 = 2\pi M, \quad \frac{1}{2\pi\alpha'} \int_B H_3 = -2\pi K.
\end{equation}

The product of $M$ and $K$ produces the background D3-charge $N$. By lifting the KS solution to F-theory, the tadpole-cancellation condition relates $N$ to the Euler number $\chi$ of the corresponding CY four-fold. The known maximal value for the Euler number $\chi_* = 1820448$ \cite{46} implies an upper bound on the background charge number,

\begin{equation}
N \equiv MK \leq \frac{\chi_*}{24} \sim 10^5.
\end{equation}

It should also be noted that the supergravity description is valid for $g_s M \gg 1$, where $g_s$ is the string coupling.

We write down the leading order background geometry in the following form,

\begin{equation}
ds^2 = h(r)^2 g^{(4)}_{\mu\nu} dx^\mu dx^\nu + h(r)^{-2} (dr^2 + r^2 g^{(5)}_{mn} d\theta^m d\theta^n),
\end{equation}

where $x^\mu$ ($\mu = 0, 1, 2, 3$) are the external four-dimensional coordinates, $r$ is the radial coordinate which decreases as one approaches the tip of the throat, and $\theta^m$ ($m = 5, 6, 7, 8, 9$) are
the five-dimensional angular coordinates. \( h(r) \) is the warp factor, which is to leading order independent of the angular coordinates. The warping at the tip of the throat is determined by the flux numbers in [3.1],

\[
h_0 \sim \exp \left( -\frac{2\pi K}{3g_s M} \right). \tag{3.4}
\]

The fluxes also stably deform the conifold, therefore setting an IR cutoff for the radial coordinate \( r \),

\[
r_0 \sim (g_s M \alpha')^{1/2} h_0. \tag{3.5}
\]

Hence the \( S^3 \) at the throat tip has radius of order \( r_0/h_0 \).

Away from the tip, the warp factor is approximated by \( h(r) \sim r/R \), where the curvature radius of \( AdS_5 \) space is given by

\[
R^4 \sim g_s N \alpha'^2. \tag{3.6}
\]

We assume that the bulk CY dominates the six-dimensional volume of the internal space, i.e. the typical length scale of the internal bulk \( L \) satisfies \( L \gtrsim R \). Hence the four-dimensional Planck mass is

\[
M_{Pl}^2 \sim \frac{L^6}{g_s^2 \alpha'^4}. \tag{3.7}
\]

Before proceeding, let us pause for a moment to explain why we consider an \( \overline{D3} \)-brane and not a \( D3 \)-brane for the curvaton. This is due to the \( D3 \) being free from force in a KS throat, which can be seen as follows: The form of the self-dual five-form flux is determined by the Bianchi identity,

\[
\tilde{F}_5 = (1 + \ast) d\alpha(r, \theta^a) \wedge \sqrt{-g^{(4)}} dx^0 \wedge dx^1 \wedge dx^2 \wedge dx^3, \tag{3.8}
\]

where \( \alpha \) is a function of the internal coordinates. Then from the Dirac-Born-Infeld (DBI) and the Chern-Simons (CS) terms in the D-brane action (cf. (3.10)), one can see that a probe \( D3 \) feels a potential proportional to \( h^4 - \alpha \), whereas \( h^4 + \alpha \) for an \( \overline{D3} \). Since

\[
h^4 = \alpha \tag{3.9}
\]

is satisfied in a KS throat (and more generally in GKP-type compactifications with imaginary self-dual fluxes), \( D3 \)-branes are free from force. As long as the compactification of the bulk CY is of the GKP-type, the no-force condition for a \( D3 \)-brane is preserved, i.e., isometry-breaking bulk effects do not give any potential for a \( D3 \), unless the bulk effects drastically distort the throat geometry from the KS solution. (Note that, by contrast, nonperturbative effects for moduli stabilization correct (3.9) and give similar forces to a \( D3 \) and an \( \overline{D3} \) [47]. Also, the relation (3.9) is violated by \( D3 \)-branes.) Since the interplay of bulk effects and nonperturbative effects plays an important role in our model, we consider probe \( \overline{D3} \)-branes at the tip of throats.
### 3.2 Effective Action of the Curvaton

As was explained at the end of the previous subsection, we consider a probe $\overline{D3}$-brane inside a GKP-type warped throat. Assuming the $\overline{D3}$ to stretch out along the external space, the $\overline{D3}$ is driven towards the tip of the throat where the warp factor is minimized [48]. In this subsection we study the dynamics of an $\overline{D3}$-brane in the angular directions at the throat tip. From these angular directions does the curvaton(s) arise.

The effective action for the $\overline{D3}$’s angular degrees of freedom is given by the DBI and the CS terms,

$$S = -T_3 \int d^4 \xi \sqrt{-\det G_{\mu\nu}} \left(1 - \bar{\Psi} i D \Psi\right) - T_3 \int C_4,$$

where $T_3 \sim 1/(g_s \alpha'^2)$ is the D3 ($\overline{D3}$) tension. Note that in the DBI action we have included the leading order term for the world-volume fermions $\Psi$ [49, 50], into which the angular degrees of freedom (i.e. the curvaton) decay first. (We assume that (some of) the world-volume fermions are significantly lighter than the curvaton.) We dropped gauge fields on the brane since their interactions with the curvaton show up from four-point interaction terms and can be neglected.

When the $\overline{D3}$ is restricted to the throat tip, the determinant of the induced metric is

$$\det G_{\mu\nu} = \det \left(h_0^2 g^{(4)}_{\mu\nu} + h_0 r_0^2 g^{(3)}_{mn} \partial_\mu \theta^m \partial_\nu \theta^n\right) \approx h_0^8 g^{(4)} \left(1 + h_0^{-4} r_0^2 g^{(3)}_{mn} g^{(4)}_{\mu\nu} \partial_\mu \theta^m \partial_\nu \theta^n\right),$$

where we have set the world-volume coordinates to coincide with the external coordinates, and $g^{(3)}_{mn}$ denotes the metric of the $S^3$ at the tip. Then by using (3.8), the four-dimensional Lagrangian from (3.10) is

$$\frac{\mathcal{L}}{-g^{(4)}} \approx -T_3 (h_0^4 + \alpha_0) - T_3 r_0^2 (\partial \theta)^2 + T_3 h_0^4 \bar{\Psi} i D \Psi,$$

where we dropped numerical factors and used the abbreviation: $(\partial \theta)^2 \equiv g^{(3)}_{mn} g^{(4)\mu\nu} \partial_\mu \theta^m \partial_\nu \theta^n$. We have also ignored terms which includes more than two derivatives in deriving (3.11) and (3.12). For simplicity, in most part of this paper we drop numerical factors and carry out order-of-magnitude estimates.

So far, the Lagrangian does not include any angular dependent potential and the $\overline{D3}$ enjoys the isometry. However, as was briefly mentioned at the beginning of this section, bulk effects break the throat’s isometries and perturb the background geometry in the following form,

$$\delta (h^4) \sim h^\Delta f (\theta^m), \quad \delta \alpha \sim h^\Delta f (\theta^m).$$

---

8 One can check that the decay of the curvaton into world-volume fermions induced by the three-point term in (3.20) is the most important interaction. The curvaton also interacts with world-volume gauge fields, gravitons, and Kaluza-Klein modes in the throat, but those interactions have no major effect on the curvaton dynamics, perturbatively nor through parametric resonance [51, 52].
Note that from the discussion at the end of Subsection 3.1, the perturbation for \( h^4 \) and \( \alpha \) are correlated, especially, both have the same angular dependence \( f(\theta^m) \). Such effects can be analysed via gauge/gravity duality, then the power \( \Delta \) of the warp factor suppressing the perturbations in (3.13) corresponds to the dimension of the irrelevant operator deforming the dual gauge theory [53] (see also [54, 55, 56]). In the KS throat, the leading perturbation has \( \Delta = \sqrt{28} \approx 5.3 \), giving a mass (cf. (3.21))

\[
m_{\text{bulk}}^2 \sim \frac{h_0^{3.3}}{g_s M_{\text{A'}}}
\]

to the canonically normalized angular directions.

Nonperturbative effects which stabilize the Kähler moduli also give potentials in the angular directions. In [47], explicit examples of Kähler moduli stabilization due to D7-branes wrapping four-cycles of conifold throats were studied. For example, in the simple case of the Kuperstein embedding [58] of the D7-brane, then under the assumption that an \( \overline{\text{D}3} \) at the throat tip lifts the stabilized vacuum energy to a metastable \( dS \), the angular directions receive mass of order

\[
m_{\text{np}}^2 \sim \frac{h_0^2 \epsilon}{g_s M_{\text{A'}} \mu}.
\]

Here, \( \epsilon \sim h_0^{3/2}(g_s M_{\text{A'}})^{3/4} \) is the deformation parameter of the conifold, and \( \mu \) measures the minimal radial location reached by the D7. However, it should be noted that (3.15) is not necessarily the typical value for our curvaton model, where multiple throats (e.g., the inflation throat) may exist and various uplifting mechanisms are possible (see e.g. [36]).

Taking into account the isometry breaking bulk effects and nonperturbative effects, the Lagrangian (3.12) is corrected to

\[
\mathcal{L} \sqrt{-g^{(4)}} \simeq -T_3 h_0^4 - T_3 r_0^2 (\partial \theta)^2 + T_3 h_0^4 \Psi i \bar{\Psi} \Psi

- T_3 h_0^4 f(\theta^m) - \tilde{V}_{\text{np}}(\theta^m) + T_3 h_0^4 f(\theta^m) \bar{\Psi} i \Psi,
\]

where we dropped numerical factors of order unity. Here \( \Delta \) indicates the suppression of the bulk effect by the warp factor (see (3.13)), and we have represented the angular potential from nonperturbative effects by \( \tilde{V}_{\text{np}}(\theta^m) \). When there are further effects which give potentials to the angular position (e.g., the violation of (3.9) due to \( \overline{\text{D}3} \)-branes), one can formally include them in \( \tilde{V}_{\text{np}}(\theta^m) \). The important point here is that \( f(\theta^m) \) and \( \tilde{V}_{\text{np}}(\theta^m) \) generally have different angular dependence.

Now we see that various effects induce a gradual but rather bumpy potential for an \( \overline{\text{D}3} \)-brane at the throat tip. Among the three angular directions of the \( S^3 \) tip, henceforth we
focus on the flattest one, and assume that the $D3$ is stabilized in the other directions. We refer to this flattest direction as $\theta$, and set the (local) minimum of $\tilde{V}_{np}(\theta)$ as its origin. Also, we consider the region where the bulk and nonperturbative effects are well approximated by expanding the Lagrangian up to quadratic order in $\theta$, hence (3.16) can be rewritten as

$$\mathcal{L} \sqrt{-g^{(4)}} \simeq -T_3 r_0^2 (\partial \theta)^2 + T_3 h_0^4 \bar{\psi} i \mathcal{D} \Psi - m_{\text{bulk}}^2 (\theta + \theta^2) - m_{np}^2 \theta^2 + m_{\text{bulk}}^2 (\theta + \theta^2) \bar{\psi} i \mathcal{D} \Psi, \quad (3.17)$$

where $m_{\text{bulk}}^2 \equiv T_3 h_0^\Delta$, $m_{np}^2 \equiv \partial_\theta \partial_\theta \tilde{V}_{np}(0)$. Upon expanding the bulk effects, we neglected numerical factors and wrote down schematically $f(\theta) \sim \theta^0 + \theta + \theta^2$, which suffices for our order-of-magnitude estimations. Furthermore, we have omitted constant terms which are irrelevant for us. It should also be noted that as long as $\Delta > 4$, terms of the form $T_3 h_0^\Delta \bar{\psi} i \mathcal{D} \Psi$ can well be ignored.

Let us again reparameterize $\theta$ so that the minimum of its total potential (i.e. the first two terms in the second line of (3.17)) becomes the origin. The shift is roughly

$$\theta \rightarrow \theta - \frac{m_{\text{bulk}}^2}{2(m_{np}^2 + m_{np}^2)}, \quad (3.18)$$

Canonically normalizing the fields,

$$\sigma \equiv \sqrt{T_3 r_0} \theta, \quad \psi \equiv \sqrt{T_3 h_0^2} \Psi, \quad (3.19)$$

we arrive at

$$\mathcal{L} \sqrt{-g^{(4)}} \simeq -(\partial \sigma)^2 + \bar{\psi} i \mathcal{D} \psi - (m_{\text{bulk}}^2 + m_{np}^2) \sigma^2 + g_s M_{1/2}^{1} \alpha'^{3/2} \frac{m_{\text{bulk}}^2 m_{np}^2}{m_{np}^2 + m_{np}^2} \sigma \bar{\psi} i \mathcal{D} \psi, \quad (3.20)$$

where we have neglected four-point and higher interaction terms. The masses arising from bulk and nonperturbative effects are defined as follows, respectively,

$$m_{\text{bulk}}^2 \equiv \frac{m_{\text{bulk}}^2}{T_3 r_0^2} = \frac{h_0^\Delta}{r_0^2} \sim \frac{h_0^{\Delta - 2}}{g_s M_\alpha'}, \quad (3.21)$$

$$m_{np}^2 \equiv \frac{m_{np}^2}{T_3 r_0^2} \equiv \frac{h_0^{\lambda - 2}}{g_s M_\alpha'}. \quad (3.22)$$

In the far right hand side of (3.22), we have additionally introduced the symbol $\lambda$ in order to denote the suppression of the nonperturbative effects (or more generally, the sum of all effects except from the bulk). By comparing $\Delta$ and $\lambda$, one can measure the relative strength of the bulk and nonperturbative effects at the throat tip, e.g., in the specific case of (3.14) and (3.15) with $\epsilon/\mu \sim h_0^{3/2}$, then $\Delta \sim 5.3$ and $\lambda \sim 5.5$. The point we would like to emphasize is that various effects with different angular dependence misalign the (local) minima of the potential and interaction terms, hence providing decay channels to the curvaton.
Let us summarize what we have obtained. The curvaton candidate $\sigma$ has mass

$$m^2 \sim m^2_{\text{bulk}} + m^2_{\text{np}}, \quad (3.23)$$

and its decay rate into world-volume fermions is

$$\Gamma \sim \left( \frac{g_s M^{1/2} \alpha^3/2}{\hbar_0^3} \frac{m^2_{\text{bulk}} m^2_{\text{np}}}{m^2_{\text{bulk}} + m^2_{\text{np}}} \right)^2 m^3 \sim \frac{g_s^2 M \alpha^3}{\hbar_0^6} \frac{m^4_{\text{bulk}} m^4_{\text{np}}}{(m^2_{\text{bulk}} + m^2_{\text{np}})^{1/2}}. \quad (3.24)$$

Specifically, when bulk effects are dominant over nonperturbative effects, i.e. $\Delta < \lambda$, then

$$m^2 \sim \frac{\hbar_0^{\Delta-2}}{g_s M \alpha}, \quad \Gamma \sim \frac{h_0^{2\lambda+\frac{3}{2}\Delta-13}}{g_s^{3/2} M^{5/2} \alpha^{1/2}}. \quad (3.25)$$

The opposite case where nonperturbative effects are stronger can be treated simply by exchanging $\Delta$ and $\lambda$.

The field range of $\sigma$ is restricted by the radius of the $S^3$ tip. From (3.19),

$$\sigma \lesssim \sqrt{T_{3R_0}} \sim \frac{h_0 M^{1/2}}{\alpha^{1/2}}. \quad (3.26)$$

Then statistically one expects that the field value of the curvaton during inflation is

$$\sigma_s \sim \frac{h_0 M^{1/2}}{\alpha^{1/2}}. \quad (3.27)$$

We assume that the number of peaks and valleys of the periodic angular potential to be of order one, so that the approximation of the quadratic potential in (3.20) be valid. Here, from (3.6) and (3.7) one finds that the condition (2.4) is always satisfied under $R \lesssim L$ and the weak string coupling $g_s < 1$,

$$\frac{\sigma_s}{M_{\text{Pl}}} \sim \frac{h_0 g_4^{1/4}}{N^{1/4} K^{1/2}} \left( \frac{R}{L} \right)^3 < 1. \quad (3.28)$$

### 3.3 Consistency Conditions

Before moving on to discuss cosmology, we should mention some consistency conditions for the above analyses. In addition to the requirements discussed in Section 2, the following microscopic constraints will further restrict the parameter space.

**Speed Limit of the Curvaton**

We have expanded the DBI action up to two derivatives, hence for the validness of this procedure the curvaton has to be moving nonrelativistically. This requirement is also important for avoiding significant back reaction of the curvaton brane to the throat.
The light speed at the tip is given by \( \dot{\theta}_{\text{rel}} \sim \frac{h_0^2}{r_0} \), whereas the maximal speed of a harmonically oscillating curvaton is \( \dot{\theta}_{\text{max}} \sim m \). Therefore we require

\[
\frac{mr_0}{h_0^2} \sim \frac{m(g_* M_{\alpha'})^{1/2}}{h_0} < 1. \tag{3.29}
\]

From (3.21) and (3.22) it is clear that as long as \( \Delta, \lambda > 4 \), the curvaton brane is nonrelativistic.

**Stringy Corrections to the Throat**

One should be aware of stringy corrections during inflation. When the Hubble parameter during inflation is larger than the local string scale of the throat, such corrections can become significant and may even shorten the throat \[59\], therefore we require

\[
H_{\text{inf}} \frac{\alpha'^{1/2}}{h_0} < 1. \tag{3.30}
\]

Note that this condition can be rewritten in terms of \( \sigma_* \) (3.27) as

\[
M^{1/2} H_{\text{inf}} \frac{\alpha'}{\sigma_*} < 1, \tag{3.31}
\]

which guarantees the density perturbation from the curvaton to be nearly Gaussian (2.7).

**Curvaton’s Energy Scale and the Local String Scale**

The curvaton’s oscillation energy also should be smaller than the local string scale for avoiding serious stringy corrections in the throat, hence

\[
m^2 \sigma_*^2 \left( \frac{\alpha'^{1/2}}{h_0} \right)^4 \sim \frac{m^2 M_{\alpha'}}{h_0^2} < 1. \tag{3.32}
\]

One sees that this requirement contains the speed limit condition (3.29) provided \( g_* < 1 \).

Before ending this section, let us mention that the condition

\[
\Delta, \lambda > 4 \tag{3.33}
\]

following from (3.29) and (3.32) can also be understood as the requirement that the bulk and nonperturbative effects be weak corrections to the original angular independent background.

\[\text{10Our observable universe might have happened to be in a place which is less likely than the statistically preferred } \sigma_*, \text{ (3.27), e.g., extremely near the (local) minimum/maximum of the periodic angular potential. In such case, (2.7) can be violated and the curvaton may contribute to producing large non-Gaussianity.}\]
One way of understanding this is by looking at the induced mass (3.21) and (3.22): One expects from naive power counting that \( m^2 \propto h_0^2 \), whereas the actual mass is \( m^2 \propto h_0^{\Delta - 2} h^{\lambda - 2} \).

We also note that under (3.33), the condition (2.5) becomes trivial. It suffices to show this in the bulk effect dominant case (3.25), where

\[
\frac{\Gamma}{m} \sim \frac{h_0^{2\lambda + \Delta - 12}}{g_s M^2} < 1. \tag{3.34}
\]

4 Generating the Cosmological Perturbations

Equipped with the information on the angular position of the \( \overline{D3} \)-brane and the consistency conditions discussed in the previous sections, let us look into the parameter space and show that the angular oscillation of the \( D3 \) at the throat tip actually plays the role of the curvaton. In this section we investigate several scenarios of interest. First we consider the case where the curvaton generates the observed curvature perturbation (and in some situations large non-Gaussianity) in Subsection 4.1. Then, cases where the curvaton contributes in generating large non-Gaussianity either through curvature or isocurvature perturbations are discussed, respectively, in Subsection 4.2 and 4.3.

We assume that the SM particles are realized on the curvaton \( \overline{D3} \)-brane in Subsection 4.1 and 4.2. How the SM particles are realized on the world-volume of the D-brane is out of the scope of this paper, though one naively expects that the interactions among the open string modes on the brane is suppressed by the local string scale. As can be seen from (3.24), or more explicitly from (3.25) and (3.33), the decay rate from the curvaton to the world-volume fermions is further suppressed by the warp factor. Therefore, after the curvaton decay into world-volume fermions, we can expect the fermions to soon decay into or thermalize with the SM particles (or particles that eventually turn into SM particles). In Subsection 4.3, the SM brane can be far apart from the curvaton brane.

For the cosmic inflation, we do not specify its details. However, throughout this paper we have assumed that the inflaton energy is transferred to radiation right after the end of inflation. For instance, one can have in mind brane-antibrane inflation in a warped throat (possibly, a throat different from the throat containing the curvaton brane). From the brane-antibrane annihilation which ends inflation, heavy closed strings are produced, which further decay into lighter states. Such processes happen right after inflation ends, and if there are branes left, the inflation energy will soon be converted to radiation on the world-volume of the residual brane. (However, one may have to be careful with stable angular KK modes, see e.g. [60, 61]. Of course, such problems do not arise in other inflation mechanisms, such as modular inflation models.)

We should also comment on the “overshooting problem” which states that the Hubble parameter during inflation is bounded by the gravitino mass, i.e. \( H_{\text{inf}} \lesssim m_{3/2} \) \([62]\), in the simplest inflationary models based on the moduli stabilization scenario of \([39]\). Requiring a gravitino mass of \( \mathcal{O}(1) \text{TeV} \), the inflation energy scale is restricted to be very low. If the
curvaton generates the observed curvature perturbations, i.e. the inequality \( (2.10) \) being saturated, then one can show from the requirements laid out in Section 2 that the Hubble parameter during inflation cannot be that small. However, as we briefly mentioned in Section 1, recent studies have demonstrated several ways to cure this problem, e.g., by introducing more general class of Kähler moduli stabilization with superpotential of the racetrack form, with positive exponents, and so on [62, 10, 11, 12, 13, 14]. We expect such mechanisms to be operating and the inflation energy scale be disentangled from the gravitino mass.

### 4.1 Curvaton Scenario

Let us consider the case where the angular fluctuation of the \( \overline{D3} \)-brane at the throat tip generates the observed primordial curvature perturbation. We assume the SM particles to be realized on the curvaton \( \overline{D3} \)-brane.

In cases where the curvaton dominates the universe before it decays, the reheating of the universe is sourced by the decay of the curvaton. We require the curvaton to decay before BBN, but if one also wants to incorporate baryogenesis, then the condition \( (2.6) \) should be corrected according to the baryogenesis scenario.

On the other hand, if the curvaton is subdominant at decay, then reheating should rely on the remnants of inflation, e.g., one can imagine the source that drives inflation to be also in the SM throat (i.e. the curvaton throat) and that the inflation energy directly flows into the SM brane, or it could also be that the inflation energy tunnels into the SM throat from the bulk or from some other throat [60, 63, 64, 65, 66, 67, 68, 69]. Although we keep our discussion general and do not go into details beyond the curvaton model, let us address some issues which may rise in the curvaton-subdominant case: Depending on how the inflation energy is transported to the SM brane, curvaton production after inflation may happen and the curvaton contribution to the curvature perturbation may be reduced [70]. Also, large isocurvature perturbation in cold dark matter (CDM) or baryons can be produced if such particles are created significantly before the curvaton decay, and/or if the curvaton itself creates them [71, 72, 73, 74, 75, 76].

Recall from the discussions in Section 2 and Section 3 that after dropping the conditions that trivially follow from the other ones, we arrive at the following six conditions: \( (2.2), (2.6), (2.10), (2.11), (3.30), (3.32) \). Here, especially, \( (2.10) \) should be saturated since we require the curvaton to generate the observed curvature perturbation. The \( f_{\text{NL}} \) constraint is now simply

\[
f_{\text{NL}} \sim \frac{1}{r} \lesssim 100, \tag{4.1}
\]

and the inflation energy scale and the produced non-Gaussianity are related by

\[
\frac{H_{\text{inf}}}{M_{\text{Pl}}} \sim 10^{-5} h_0 g_s M_{\text{Pl}}^{1/2} \left( \frac{\alpha'^{1/2}}{L} \right)^3 f_{\text{NL}}. \tag{4.2}
\]
Under (2.10) and (4.1), the condition (3.30) is automatically satisfied,

\[ H_{\text{inf}} \frac{\alpha'^{1/2}}{h_0} \lesssim \frac{M^{1/2}}{10^3} < 1, \]  

where we have used the known maximal value of the CY four-fold Euler number (3.2) in obtaining the second inequality.

Therefore, after fixing the inflation energy scale from the saturated (2.10), the independent conditions that restrict the parameter space are (2.2), (2.6), (2.11), and (3.32). Especially in the bulk effect dominant case (i.e. \( \Delta < \lambda \)) (3.25), the four conditions take the following form, respectively,

\[ r^2 \frac{10^{10} h_0^{\Delta - 4}}{g_s M^2} \ll 1, \quad 10^{42} h_0^{2\lambda + \frac{2}{3} \Delta - 13} \left( \frac{\alpha'^{1/2}}{L} \right)^3 > 1, \quad \frac{1}{r} \lesssim 10^2, \quad \frac{h_0^{\Delta - 4}}{g_s} < 1, \]  

where

\[ r \sim \min \left\{ h_0^{8 - \lambda - \Delta + \frac{\Delta}{2} g_s^{5/2} M^2 \left( \frac{\alpha'^{1/2}}{L} \right)^6}, 1 \right\}. \]  

We illustrate the parameter space on the \( \Delta - \lambda \) plane in Figure 1\textsuperscript{11}. One sees that as \( \Delta \) or \( \lambda \) increase, i.e. the bulk or nonperturbative effects weaken, the decay rate of the curvaton is suppressed and the curvaton comes closer to dominating the universe before decay. Also, the produced non-Gaussianity (displayed in the figure as contour lines) decreases as \( r \to 1 \). The figure clearly shows that when either the bulk or nonperturbative effect is absent, i.e. \( \Delta \) or \( \lambda \to \infty \), the curvaton cannot decay\textsuperscript{12} and one crosses the orange line which is the BBN constraint (2.6).

As we take different values for the four parameters \( g_s, M, L/\alpha'^{1/2}, \) and \( h_0 \), the consistent region (the yellow region in the figure) deforms and shifts in the \( \Delta - \lambda \) plane. It can also split into \( r \gtrsim 1 \) regions when \( h_0 \) becomes large. Though we have shown only a single example, one can check that consistent curvaton scenarios are allowed for broad ranges of the parameters.

### 4.2 Non-Gaussianity through Curvature Perturbations

Next we consider the case where the curvaton produces large non-Gaussianity, but generates only negligible curvature perturbations. Hence in this subsection we assume some other mechanism (e.g. the inflaton) to be responsible for generating the observed curvature perturbations. As in the previous subsection, the SM particles are assumed to be realized on the curvaton D3-brane.

\textsuperscript{11}As we carry out only an order-of-magnitude estimate, the explicit values of the parameters for each figure should not be taken so seriously, e.g., the actual \( L/\alpha'^{1/2} \) is larger than the set values since we dropped numerical factors in (3.7).

\textsuperscript{12}Strictly speaking, the curvaton does decay due to the interaction terms we have ignored in deriving (3.20), but such effects are negligibly small.
Figure 1: The consistency conditions for the curvaton scenario on the \( \Delta \)-\( \lambda \) plane, where the other parameters are set to \( g_s = 0.1 \), \( M = 300 \), \( L/\alpha^{1/2} = 3 \), \( h_0 = 10^{-5} \) (hence \( N \sim 50000 \), \( K \sim 200 \), \( M_{\text{Pl}}(\alpha^{1/2} \sim 300) \). The lines denote where each condition is saturated, blue: masslessness (2.2), orange: BBN (2.6), red: \( f_{\text{NL}} \) (2.11), purple: curvaton oscillation energy (3.32). The yellow region satisfies all four conditions. On the right (left) side of the dashed line, the curvaton dominates (subdominates) the universe at the decay epoch. The produced non-Gaussianities are also shown as contour lines for \( f_{\text{NL}} \). Here the inflation scale can be estimated from \( H_{\text{inf}}/M_{\text{Pl}} \sim 10^{-11} f_{\text{NL}} \).

Since we consider curvatons generating large non-Gaussianity, we focus on cases where the curvaton decays before dominating the universe. (One clearly sees from (2.10) and (2.11) that when \( r \sim 1 \), the amplitude constraint is severer than the \( f_{\text{NL}} \) constraint.) If the energy fraction of the curvaton is sufficiently small, then it can survive beyond the BBN epoch without ruining nucleosynthesis.

The independent consistency conditions are (2.2), (2.10), (2.11), (3.30), (3.32), and depending on \( r \), the BBN constraint (2.6) is also required. In Figure 2 we show the conditions in the \( \Delta \)-\( \lambda \) plane under a certain set of parameters. Note that we now have an additional parameter \( H_{\text{inf}}/M_{\text{Pl}} \) since (2.10) need not be saturated. On the red line in the figure, the curvaton produces the observationally allowed maximum non-Gaussianity of order \( f_{\text{NL}} \sim 100 \). In contrast to the previous subsection, the non-Gaussianity increases for larger \( \Delta \) or \( \lambda \), since then \( r \) becomes larger and the non-Gaussian curvature perturbations generated by the curvaton increase.
4.3 Non-Gaussianity through Isocurvature Perturbations

In this subsection, we consider the case where the curvaton survives until now and contribute
to dark matter. Here the SM particles need not be on the curvaton brane. Specifically, in
cases where the throat containing the curvaton brane is geometrically separated from where
the SM particles are realized, then the curvaton dark matter interacts with SM particles
only through graviton mediation and thus behaves as hidden dark matter (see e.g. [60]).

Since the curvaton does not decay, here we need not require both the bulk and nonper-
turbative effects, and one of them can be totally absent, i.e. $\Delta$ or $\lambda \rightarrow \infty$. D3-branes which
do not feel bulk effects are thus always of this type (unless there are additional effects giving
angular dependencies.)

As for the cosmological perturbations, the effects of having such curvatons are similar
to that of axions. Since they are decoupled from (most of) the thermal history of the SM
particles, their fluctuations entirely contribute to isocurvature perturbations. Here we discuss
non-Gaussianity produced from the curvaton through isocurvature perturbations. As in the
previous subsection, we assume the universe to be reheated by the remnants of inflation, and
the observed curvature perturbations to be generated by some mechanism other than the
curvaton.

Here, additional constraints are required. Instead of the BBN constraint (2.6), we require
the curvaton to be stable until now, therefore

$$10^{60} \frac{\Gamma}{M_{Pl}} < 1.$$  \hspace{1cm} (4.6)

Also, the present energy density of the curvaton should not exceed that of dark matter, $\Omega_{\sigma 0} \lesssim \Omega_{CDM 0}$. Assuming radiation domination from the end of inflation to the time of matter-
radiation equality, the present curvaton energy density is

$$\rho_{\sigma 0} \sim m^2 \sigma^2 (t_{osc}/t_{eq})^{3/2} (a_{eq}/a_0)^3,$$

and one arrives at the requirement

$$R \equiv \frac{\Omega_{\sigma 0}}{\Omega_{CDM 0}} \sim 10^{28} \frac{m^{1/2} \sigma^2}{M_{Pl}^{3/2}} \lesssim 1.$$  \hspace{1cm} (4.7)

Fluctuations of the curvaton dark matter give rise to isocurvature perturbations between
CDM and radiation,

$$S_{CDM} \sim R \left[ \frac{2 \delta \sigma_s}{\sigma_s} + \left( \frac{\delta \sigma_s}{\sigma_s} \right)^2 \right].$$  \hspace{1cm} (4.8)

Here the isocurvature and curvature perturbations are expected to be uncorrelated, hence
the observational bound on isocurvature perturbation [9] imposes

$$\frac{\mathcal{P}_S}{\mathcal{P}_\zeta} \sim \frac{R^2}{\mathcal{P}_{COBE}} \left( \frac{H_{inf}}{\sigma_s} \right)^2 \lesssim \frac{1}{10^4}.$$  \hspace{1cm} (4.9)
where $P_{\text{COBE}} \sim 10^{-10}$. Also, the non-Gaussianity from isocurvature perturbation in terms of the non-linearity parameter is

$$f_{\text{NL}} \sim \frac{R^3}{P_{\text{COBE}}^2} \left( \frac{H_{\text{inf}}}{\sigma_*} \right)^4 \lesssim 100. \quad (4.10)$$

One should note that the non-Gaussianity from isocurvature and curvature perturbations manifest themselves differently in the CMB temperature fluctuations. Especially, non-Gaussian effects from isocurvature perturbations are enhanced in the large scales [73]. For detailed analyses on this issue, see [73, 74, 75, 76, 77].

Now (2.6), (2.10), and (2.11) are replaced by the above conditions, and we arrive at the following seven independent consistency conditions: (2.2), (3.30), (3.32), (4.6), (4.7), (4.9), and (4.10). In Figure 3 we show an example where the KS throat with $\Delta = \sqrt{28}$ bulk effect can produce large non-Gaussianity. In the example, the consistent region is bounded by the isocurvature bound (4.9) rather than the $f_{\text{NL}}$ bound (4.10). Larger curvaton mass leads to larger $R$, therefore larger isocurvature perturbation and non-Gaussianity.

In this subsection we have studied curvatons that do not decay. However, in cases where the curvaton brane is geometrically separated from the SM sector, the curvaton can decay and still generate large non-Gaussianity. The light open string modes excited from the curvaton decay will serve as hidden dark radiation and/or matter and can have significant effects through isocurvature perturbations, as in the case of the non-decaying curvatons. It would be interesting to examine systematically the range of possibilities arising from oscillating D-branes in multi-throat compactifications.

5 Conclusions

We have proposed a model for generating the primordial perturbations and reheating our universe from angular oscillations of D-branes at the tip of throats. The geometrical features of throats – warping and (approximate) isometries – yielded curvaton scenarios. We have also seen that effects that break the force-free condition of the D-brane in the isometry directions, such as the isometry breaking bulk effects and moduli stabilizing nonperturbative effects played an important role in our model. Depending on the (un)balance between the various features of the background geometry, the curvaton model showed different behaviours. We have studied cases where the curvaton generated the observed curvature perturbations (Subsection 4.1), the curvaton contributed mainly to non-Gaussianity through curvature perturbations (Subsection 4.2), and through isocurvature perturbations (Subsection 4.3). Each scenario could be realized in a wide range of parameter space. In other words, our model may be considered as generally arising from compactification scenarios containing warped throats with isometries. Especially the case considered in Subsection 4.3 arise no matter where the SM particles are realized, and therefore may serve as a test for discussing the
validness of (multi-)throat compactification scenarios. (See also [61] for discussions in this direction.) The curvaton model is capable of producing large non-Gaussianity, and upcoming CMB experiments are expected to allow us to give more rigorous arguments.

In this paper we tried to keep our discussions general and did not go into details of individual moduli stabilization mechanisms. We have treated the strengths of the bulk and nonperturbative effects as free parameters $\Delta$ and $\lambda$, and only made order-of-magnitude estimations. However, based on previous studies, one can carry out more elaborate analyses on the angular potentials for D-branes at the $S^3$ tip of warped deformed conifolds, for certain moduli stabilizing scenarios. It is necessary and interesting to explore the throat tip potential and also a wide class of compactification scenarios for concrete realizations of the curvaton. We note again that in most part of the paper we simply assumed the SM particles to be realized on the curvaton $\overline{D3}$-brane(s). Inclusion of a more realistic SM sector is necessary for a fully consistent model, and is crucial for detailed studies of the thermal history of the universe.

One of the general lessons of our work is that in string theory, one finds it quite natural to consider fields other than the inflaton for generating the primordial perturbations. Such light fields can be realized in string theory, thus giving rise to mechanisms which may have seemed too intricate from the phenomenological point of view. It is fair to say that top-down approaches to inflationary cosmology can provide us with rich ideas beyond the standard slow-roll inflation pictures.

Acknowledgements

We would like to thank Robert Brandenberger, Damien Easson, Toshiya Imoto, Etsuko Kawakami, Shunichiro Kinoshita, Kazunori Nakayama, Brian Powell, Misao Sasaki, Toyokazu Sekiguchi, Gary Shiu, Shigeki Sugimoto, Fuminobu Takahashi, Atsushi Taruya, Henry Tye, Yi Wang, and Jun’ichi Yokoyama for very helpful discussions. T.K. is also grateful to Masahiro Kawasaki and Katsuhiko Sato for their continuous support. The work of S.M. was supported in part by MEXT through a Grant-in-Aid for Young Scientists (B) No. 17740134, by JSPS through a Grant-in-Aid for Creative Scientific Research No. 19GS0219, and by the Mitsubishi Foundation. This work was supported by World Premier International Research Center Initiative (WPI Initiative), MEXT, Japan.
Figure 2 (left): The consistency conditions on the $\Delta$-$\lambda$ plane for the case where the curvaton contributes mainly to non-Gaussianity. The parameters are set to $g_s = 0.1$, $M = 200$, $L/\alpha'^{1/2} = 10$, $h_0 = 10^{-4}$, $H_{\text{inf}}/M_{\text{Pl}} = 10^{-9}$ (hence $N \sim 20000$, $K \sim 90$, $M_{\text{Pl}}\alpha'^{1/2} \sim 10000$). The lines denote where each condition is saturated, blue: masslessness (2.2), orange: BBN (2.6), green: curvature perturbation (2.10), red: $f_{\text{NL}}$ (2.11), purple: curvaton oscillation energy (3.32). The stringy correction (3.30) condition is satisfied in the entire displayed area. The yellow region satisfies all six conditions. Also, the dashed line denotes where the curvaton starts to dominate the universe before decay. The produced non-Gaussianities are shown as contour lines for $f_{\text{NL}}$.

Figure 3 (right): The consistency conditions on the $\Delta$-$\lambda$ plane for the case where the curvaton is stable until now. The parameters are set to $g_s = 0.1$, $M = 100$, $L/\alpha'^{1/2} = 10$, $h_0 = 4 \times 10^{-9}$, $H_{\text{inf}}/M_{\text{Pl}} = 10^{-13}$ (hence $N \sim 9000$, $K \sim 90$, $M_{\text{Pl}}\alpha'^{1/2} \sim 10000$). The lines denote where each condition is saturated, orange: stability (4.6), green: isocurvature perturbation (4.9), red: $f_{\text{NL}}$ (4.10). The masslessness (2.2), stringy correction (3.30), curvaton oscillation energy (3.32), and DM (4.7) conditions are satisfied in the entire displayed area. The yellow region satisfies all seven conditions. The produced non-Gaussianities are shown as contour lines for $f_{\text{NL}}$. 
References

[1] A. A. Starobinsky, “A new type of isotropic cosmological models without singularity,” Phys. Lett. B 91, 99 (1980).

[2] K. Sato, “First Order Phase Transition Of A Vacuum And Expansion Of The Universe,” Mon. Not. Roy. Astron. Soc. 195, 467 (1981).

[3] A. H. Guth, “The Inflationary Universe: A Possible Solution To The Horizon And Flatness Problems,” Phys. Rev. D 23, 347 (1981).

[4] J. M. Cline, “String cosmology,” arXiv:hep-th/0612129.

[5] R. Kallosh, “On Inflation in String Theory,” Lect. Notes Phys. 738, 119 (2008) arXiv:hep-th/0702059.

[6] C. P. Burgess, “Lectures on Cosmic Inflation and its Potential Stringy Realizations,” PoS P2GC, 008 (2006) [Class. Quant. Grav. 24, S795 (2007 POSCI,CARGSE2007,003.2007)] arXiv:0708.2865 [hep-th].

[7] L. McAllister and E. Silverstein, “String Cosmology: A Review,” Gen. Rel. Grav. 40, 565 (2008) arXiv:0710.2951 [hep-th].

[8] D. Baumann and L. McAllister, “Advances in Inflation in String Theory,” arXiv:0901.0265 [hep-th].

[9] E. Komatsu et al. [WMAP Collaboration], “Five-Year Wilkinson Microwave Anisotropy Probe (WMAP) Observations:Cosmological Interpretation,” Astrophys. J. Suppl. 180, 330 (2009) arXiv:0803.0547 [astro-ph].

[10] M. Badziak and M. Olechowski, “Volume modulus inflation and a low scale of SUSY breaking,” JCAP 0807, 021 (2008) arXiv:0802.1014 [hep-th].

[11] H. Abe, T. Higaki, T. Kobayashi and O. Seto, “Non-perturbative moduli superpotential with positive exponents,” Phys. Rev. D 78, 025007 (2008) arXiv:0804.3229 [hep-th].

[12] J. P. Conlon, R. Kallosh, A. Linde and F. Quevedo, “Volume Modulus Inflation and the Gravitino Mass Problem,” JCAP 0809, 011 (2008) arXiv:0806.0809 [hep-th].

[13] M. Badziak and M. Olechowski, “Volume modulus inflection point inflation and the gravitino mass problem,” JCAP 0902, 010 (2009) arXiv:0810.4251 [hep-th].

[14] H. Y. Chen, L. Y. Hung and G. Shiu, “Inflation on an Open Racetrack,” JHEP 0903, 083 (2009) arXiv:0901.0267 [hep-th].

[15] S. Kachru, R. Kallosh, A. Linde, J. M. Maldacena, L. P. McAllister and S. P. Trivedi, “Towards inflation in string theory,” JCAP 0310, 013 (2003) arXiv:hep-th/0308055.
[16] D. Baumann, A. Dymarsky, I. R. Klebanov, J. M. Maldacena, L. P. McAllister and A. Murugan, “On D3-brane potentials in compactifications with fluxes and wrapped D-branes,” JHEP 0611, 031 (2006) [arXiv:hep-th/0607050].

[17] D. Baumann, A. Dymarsky, I. R. Klebanov, L. McAllister and P. J. Steinhardt, “A Delicate Universe,” Phys. Rev. Lett. 99, 141601 (2007) [arXiv:0705.3837 [hep-th]].

[18] A. Krause and E. Pajer, “Chasing Brane Inflation in String-Theory,” JCAP 0807, 023 (2008) [arXiv:0705.4682 [hep-th]].

[19] D. Baumann, A. Dymarsky, I. R. Klebanov and L. McAllister, “Towards an Explicit Model of D-brane Inflation,” JCAP 0801, 024 (2008) [arXiv:0706.0360 [hep-th]].

[20] A. Linde, “Fast-Roll Inflation,” JHEP 0111, 052 (2001) [arXiv:hep-th/0110195].

[21] L. Kofman and S. Mukohyama, “Rapid roll Inflation with Conformal Coupling,” Phys. Rev. D 77, 043519 (2008) [arXiv:0709.1952 [hep-th]].

[22] T. Kobayashi and S. Mukohyama, “Conformal Inflation, Modulated Reheating, and WMAP5,” Phys. Rev. D 79, 083501 (2009) [arXiv:0810.0810 [hep-th]].

[23] T. Kobayashi, S. Mukohyama and B. A. Powell, “Cosmological Constraints on Rapid Roll Inflation,” arXiv:0905.1752 [astro-ph.CO].

[24] A. D. Linde and V. F. Mukhanov, “Nongaussian isocurvature perturbations from inflation,” Phys. Rev. D 56, 535 (1997) [arXiv:astro-ph/9610219].

[25] K. Enqvist and M. S. Sloth, “Adiabatic CMB perturbations in pre big bang string cosmology,” Nucl. Phys. B 626, 395 (2002) [arXiv:hep-ph/0109214].

[26] D. H. Lyth and D. Wands, “Generating the curvature perturbation without an inflaton,” Phys. Lett. B 524, 5 (2002) [arXiv:hep-ph/0110002].

[27] T. Moroi and T. Takahashi, “Effects of cosmological moduli fields on cosmic microwave background,” Phys. Lett. B 522, 215 (2001) [Erratum-ibid. B 539, 303 (2002)] [arXiv:hep-ph/0110096].

[28] K. Dimopoulos and D. H. Lyth, “Models of inflation liberated by the curvaton hypothesis,” Phys. Rev. D 69, 123509 (2004) [arXiv:hep-ph/0209180].

[29] L. Pilo, A. Riotto and A. Zaffaroni, “Old inflation in string theory,” JHEP 0407, 052 (2004) [arXiv:hep-th/0401004].

[30] G. Dvali, A. Gruzinov and M. Zaldarriaga, “A new mechanism for generating density perturbations from inflation,” Phys. Rev. D 69, 023505 (2004) [arXiv:astro-ph/0303591].

[31] L. Kofman, “Probing string theory with modulated cosmological fluctuations,” arXiv:astro-ph/0303614.
[32] D. H. Lyth, “Generating the curvature perturbation at the end of inflation,” JCAP 0511, 006 (2005) [arXiv:astro-ph/0510443].

[33] L. Alabidi and D. Lyth, “Curvature perturbation from symmetry breaking the end of inflation,” JCAP 0608, 006 (2006) [arXiv:astro-ph/0604569].

[34] D. H. Lyth and A. Riotto, “Generating the curvature perturbation at the end of inflation in string Phys. Rev. Lett. 97, 121301 (2006) [arXiv:astro-ph/0607326].

[35] L. Leblond and S. Shandera, “Cosmology of the tachyon in brane inflation,” JCAP 0701, 009 (2007) [arXiv:hep-th/0610321].

[36] H. Y. Chen, J. O. Gong and G. Shiu, “Systematics of multi-field effects at the end of warped brane inflation,” JHEP 0809, 011 (2008) [arXiv:0807.1927 [hep-th]].

[37] I. R. Klebanov and M. J. Strassler, “Supergravity and a confining gauge theory: Duality cascades and chiSB-resolution of naked singularities,” JHEP 0008, 052 (2000) [arXiv:hep-th/0007191].

[38] S. B. Giddings, S. Kachru and J. Polchinski, “Hierarchies from fluxes in string compactifications,” Phys. Rev. D 66, 106006 (2002) [arXiv:hep-th/0105097].

[39] S. Kachru, R. Kallosh, A. Linde and S. P. Trivedi, “De Sitter vacua in string theory,” Phys. Rev. D 68, 046005 (2003) [arXiv:hep-th/0301240].

[40] J. F. G. Cascales, M. P. Garcia del Moral, F. Quevedo and A. M. Uranga, “Realistic D-brane models on warped throats: Fluxes, hierarchies and moduli stabilization,” JHEP 0402, 031 (2004) [arXiv:hep-th/0312051].

[41] J. H. Brodie and D. A. Easson, “Brane inflation and reheating,” JCAP 0312, 004 (2003) [arXiv:hep-th/0301138].

[42] S. Mukohyama, “Reheating a multi-throat universe by brane motion,” arXiv:0706.3214 [hep-th].

[43] K. Dimopoulos, G. Lazarides, D. Lyth and R. Ruiz de Austri, “Curvaton dynamics,” Phys. Rev. D 68, 123515 (2003) [arXiv:hep-ph/0308015].

[44] E. Komatsu and D. N. Spergel, “Acoustic signatures in the primary microwave background bispectrum,” Phys. Rev. D 63, 063002 (2001) [arXiv:astro-ph/0005036].

[45] P. Candelas and X. C. de la Ossa, “Comments on Conifolds,” Nucl. Phys. B 342, 246 (1990).

[46] A. Klemm, B. Lian, S. S. Roan and S. T. Yau, “Calabi-Yau fourfolds for M- and F-theory compactifications,” Nucl. Phys. B 518, 515 (1998) [arXiv:hep-th/9701023].
[47] O. DeWolfe, L. McAllister, G. Shiu and B. Underwood, “D3-brane Vacua in Stabilized Compactifications,” JHEP 0709, 121 (2007) arXiv:hep-th/0703088.

[48] S. Kachru, J. Pearson and H. L. Verlinde, “Brane/Flux Annihilation and the String Dual of a Non-Supersymmetric Field Theory,” JHEP 0206, 021 (2002) arXiv:hep-th/0112197.

[49] D. Marolf, L. Martucci and P. J. Silva, “Fermions, T-duality and effective actions for D-branes in bosonic backgrounds,” JHEP 0304, 051 (2003) arXiv:hep-th/0303209.

[50] D. Marolf, L. Martucci and P. J. Silva, “Actions and fermionic symmetries for D-branes in bosonic backgrounds,” JHEP 0307, 019 (2003) arXiv:hep-th/0306066.

[51] L. Kofman, A. D. Linde and A. A. Starobinsky, “Reheating after inflation,” Phys. Rev. Lett. 73, 3195 (1994) arXiv:hep-th/9405187.

[52] L. Kofman, A. D. Linde and A. A. Starobinsky, “Towards the theory of reheating after inflation,” Phys. Rev. D 56, 3258 (1997) arXiv:hep-ph/9704452.

[53] O. Aharony, Y. E. Antebi and M. Berkooz, “Open string moduli in KKLT compactifications,” Phys. Rev. D 72, 106009 (2005) arXiv:hep-th/0508080.

[54] A. Ceresole, G. Dall’Agata, R. D’Auria and S. Ferrara, “Spectrum of type IIB supergravity on AdS(5) x T(11): Predictions on N = 1 SCFT’s,” Phys. Rev. D 61, 066001 (2000) arXiv:hep-th/9905226.

[55] A. Ceresole, G. Dall’Agata and R. D’Auria, “KK spectroscopy of type IIB supergravity on AdS(5) x T(11),” JHEP 9911, 009 (1999) arXiv:hep-th/9907216.

[56] O. DeWolfe, S. Kachru and H. L. Verlinde, “The giant inflaton,” JHEP 0405, 017 (2004) arXiv:hep-th/0403123.

[57] D. Baumann, A. Dymarsky, S. Kachru, I. R. Klebanov and L. McAllister, “Holographic Systematics of D-brane Inflation,” JHEP 0903, 093 (2009) arXiv:0808.2811 [hep-th].

[58] S. Kuperstein, “Meson spectroscopy from holomorphic probes on the warped deformed conifold,” JHEP 0503, 014 (2005) arXiv:hep-th/0411097.

[59] A. R. Frey, A. Mazumdar and R. C. Myers, “Stringy effects during inflation and reheating,” Phys. Rev. D 73, 026003 (2006) arXiv:hep-th/0508139.

[60] X. Chen and S. H. Tye, “Heating in brane inflation and hidden dark matter,” JCAP 0606, 011 (2006) arXiv:hep-th/0602136.

[61] J. F. Dufaux, L. Kofman and M. Peloso, “Dangerous Angular KK/Glueball Relics in String Theory Cosmology,” Phys. Rev. D 78, 023520 (2008) arXiv:0802.2958 [hep-th].
[62] R. Kallosh and A. Linde, “Landscape, the scale of SUSY breaking, and inflation,” JHEP 0412, 004 (2004) [arXiv:hep-th/0411011].

[63] S. Dimopoulos, S. Kachru, N. Kaloper, A. E. Lawrence and E. Silverstein, “Small numbers from tunneling between brane throats,” Phys. Rev. D 64, 121702 (2001) [arXiv:hep-th/0104239].

[64] S. Dimopoulos, S. Kachru, N. Kaloper, A. E. Lawrence and E. Silverstein, “Generating small numbers by tunneling in multi-throat compactifications,” Int. J. Mod. Phys. A 19, 2657 (2004) [arXiv:hep-th/0106128].

[65] N. Barnaby, C. P. Burgess and J. M. Cline, “Warped reheating in brane-antibrane inflation,” JCAP 0504, 007 (2005) [arXiv:hep-th/0412040].

[66] L. Kofman and P. Yi, “Reheating the universe after string theory inflation,” Phys. Rev. D 72, 106001 (2005) [arXiv:hep-th/0507257].

[67] D. Chialva, G. Shiu and B. Underwood, “Warped reheating in multi-throat brane inflation,” JHEP 0601, 014 (2006) [arXiv:hep-th/0508229].

[68] P. Langfelder, “On tunnelling in two-throat warped reheating,” JHEP 0606, 063 (2006) [arXiv:hep-th/0602296].

[69] B. v. Harling, A. Hebecker and T. Noguchi, “Energy Transfer between Throats from a 10d Perspective,” JHEP 0711, 042 (2007) [arXiv:0705.3648 [hep-th]].

[70] A. Linde and V. Mukhanov, “The curvaton web,” JCAP 0604, 009 (2006) [arXiv:astro-ph/0511736].

[71] D. H. Lyth, C. Ungarelli and D. Wands, “The primordial density perturbation in the curvaton scenario,” Phys. Rev. D 67, 023503 (2003) [arXiv:astro-ph/0208055].

[72] D. H. Lyth and D. Wands, “The CDM isocurvature perturbation in the curvaton scenario,” Phys. Rev. D 68, 103516 (2003) [arXiv:astro-ph/0306500].

[73] M. Kawasaki, K. Nakayama, T. Sekiguchi, T. Suyama and F. Takahashi, “Non-Gaussianity from isocurvature perturbations,” JCAP 0811, 019 (2008) [arXiv:0808.0009 [astro-ph]].

[74] D. Langlois, F. Vernizzi and D. Wands, “Non-linear isocurvature perturbations and non-Gaussianities,” JCAP 0812, 004 (2008) [arXiv:0809.4646 [astro-ph]].

[75] T. Moroi and T. Takahashi, “Non-Gaussianity and Baryonic Isocurvature Fluctuations in the Curvaton Scenario,” Phys. Lett. B 671, 339 (2009) [arXiv:0810.0189 [hep-ph]].

[76] M. Kawasaki, K. Nakayama, T. Sekiguchi, T. Suyama and F. Takahashi, “A General Analysis of Non-Gaussianity from Isocurvature Perturbations,” JCAP 0901, 042 (2009) [arXiv:0810.0208 [astro-ph]].

24
[77] C. Hikage, K. Koyama, T. Matsubara, T. Takahashi and M. Yamaguchi, “Limits on Isocurvature Perturbations from Non-Gaussianity in WMAP Temperature Anisotropy,” arXiv:0812.3500 [astro-ph].