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Inflationary constraints on type IIA string theory

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Abstract: We prove that inflation is forbidden in the most well understood class of semi-realistic type IIA string compactifications: Calabi-Yau compactifications with only standard NS-NS 3-form flux, R-R fluxes, D6-branes and O6-planes at large volume and small string coupling. With these ingredients, the first slow-roll parameter satisfies $\epsilon \geq \frac{27}{13}$ whenever $V > 0$, ruling out both inflation (including brane/anti-brane inflation) and de Sitter vacua in this limit. Our proof is based on the dependence of the 4-dimensional potential on the volume and dilaton moduli in the presence of fluxes and branes. We also describe broader classes of IIA models which may include cosmologies with inflation and/or de Sitter vacua. The inclusion of extra ingredients, such as NS 5-branes and geometric or non-geometric NS-NS fluxes, evades the assumptions used in deriving the no-go theorem. We focus on NS 5-branes and outline how such ingredients may prove fruitful for cosmology, but we do not provide an explicit model. We contrast the results of our IIA analysis with the rather different situation in IIB.

Keywords: Supergravity Models, Cosmology of Theories beyond the SM, dS vacua in string theory, Flux compactifications.

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1. Introduction

Our desire to understand the large-scale properties of our Universe is one of the motivations for studying fundamental microphysical theories such as string theory. Indeed, observing the early Universe may be our most promising path toward confronting string theory with data. The leading paradigm for explaining the large-scale isotropy, homogeneity, and flatness of the Universe, as well as its $\mathcal{O}(10^{-5})$ seed fluctuations, is cosmological inflation [1 – 4]. Specifically, by assuming that there exist one or more scalar fields that undergo slow rolling in the early Universe in a potential energy function of just the right shape, one can explain these large-scale properties and predict the numerical values of as many as eight cosmological parameters, many of which have now been accurately measured [5 – 7]. The leading candidate for a fundamental microphysical theory is string theory, and so we would like to know how generically string theory can accommodate such potential energy functions.

It is rather well known that the conditions for inflation do not arise easily in string theory, or in other words, that a *generic* point in field space may not be expected to satisfy the slow-roll conditions [8 – 10]. In part this is because of issues like the η problem

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(essentially, that the potential varies too quickly), which also complicate attempts to build inflationary models in quantum field theory and supergravity theories. There are, however, three reasons to suspect a priori that string theory can accommodate inflation: firstly, the potential energy is typically a function of hundreds of fields, which means that there is a large field space to explore. Secondly, there are an exponentially large number of infinite families of potential energy functions, parameterized by typically hundreds of discrete fluxes in the compact space [11 – 15]. Thirdly, there are at least many millions of different topologies, such as Calabi-Yau manifolds, that in general give rise to qualitatively different physical theories in 4 dimensions. It is reasonable to suspect that occasionally in this vast space of possibilities, the conditions for inflation are satisfied.

For this reason, the past few years have seen intense investigation into the possibility of inflation driven by closed string moduli [16 – 20], axions [21 – 25], or brane positions [26 – 33] in the extra dimensions. In the most intensely studied case of IIB string compactifications on Calabi-Yau orientifolds, the conclusion at this point is that one can probably build working models, at the cost of fine-tuning the relevant potentials. Some of these models could even have interesting observable signatures [34, 35]. Reviews of this general subject appear in [36 – 40].

One feature of the existing constructions is that they are implicit, relying at some point on either non-compact models of regions of the compactification space, or on the ability to perform tunes which (though seemingly possible based on detailed theoretical considerations) are not performed explicitly. One would ideally like to build simpler models, where all of the calculations are performed in a completely explicit and reliable way. Commonly, compactifications suffer from unstabilized moduli in the low energy description, or the calculations stabilizing moduli apply in a regime that, while apparently numerically controlled, is not under parametric control.[1]

In some limits of string theory, however, we now have explicit examples of stabilized models with parametric control of the moduli potential. The best understood case occurs in massive IIA string theory; namely IIA string theory with R-R 0-form flux, compactified on a Calabi-Yau orientifold. The 10-dimensional massive IIA supergravity action was suggested in [41]. The compactification of this theory on a Calabi-Yau orientifold was performed in a 4-dimensional supergravity formalism in [42] and the stabilization was obtained in [43].[2] The 10-dimensional description of these compactifications was further studied in [45]. Since these models carry at most $\mathcal{N} = 1$ supersymmetry in 4 dimensions, and gauge groups and chiral matter can be incorporated in this context, we consider them to be at least semi-realistic.

An investigation of cosmology in these IIA compactifications was initiated in ref. [46] by considering some specific simple examples and showing that inflation could not occur in these examples. In the present work we extend that study. We use a simple scaling

[1] This in particular applies in any class of stabilized compactifications where the number of choices of fluxes, branes, etc., while perhaps very large, is *finite*. In such models, there is perforce a limit on how small g_s can be, namely the smallest g_s obtained in the finite list. Of course if the finite number is sufficiently large, the smallest attainable coupling may be quite small, so this may not be a serious limitation.

[2] With additional ingredients (geometric flux) stabilization was achieved in [44].

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analysis of various terms which appear in the low-energy 4-dimensional potential to rule out inflation and de Sitter vacua at large volume and weak coupling in all IIA Calabi-Yau models with conventional fluxes, D-branes, and O-planes. This means that inflation imposes the constraint that our Universe is not in this portion of the landscape. We emphasize that the derivation of this no-go result is only valid in the large volume limit and can be evaded by various other structures including NS 5-branes as well as geometric and non-geometric NS-NS fluxes, indicating that IIA compactifications containing these ingredients may be a good place to look for string models with inflation and/or de Sitter vacua. Indeed, as we were completing this work we received a copy of [47], in which more explicit IIA de Sitter models are constructed by using geometric NS-NS fluxes, 5-branes, and various other ingredients.

The structure of this article is as follows: in section 2 we summarize the IIA supergravity theory in 10 dimensions and outline the dimensional reduction to 4 dimensions. We explain the key step in the analysis of the paper, which involves considering 2-dimensional slices in the full moduli space parameterized by the volume and dilaton moduli of the compactification. The behavior of the four (space-time) dimensional potential on these 2-dimensional slices of moduli space allows us to place a lower bound on the slow-roll parameter ϵ. In section 3 we compute the scaling of the various terms appearing in the 4-dimensional potential energy function V in terms of the two model-independent moduli. We prove that both inflation and de Sitter vacua are forbidden at large volume and weak string coupling when standard fluxes, D6-branes, and O6-planes are included; the slow-roll parameter is bounded below in this case by $\epsilon \geq \frac{27}{13}$ whenever $V > 0$. In section 4 we describe some additional ingredients such as NS 5-branes, geometric fluxes and non-geometric NS-NS fluxes which can be included in type IIA and which lead to terms in the 4-dimensional potential with scaling properties allowing us to evade the no-go theorem. In section 5 we discuss the type IIB theory. We show how the structure of the IIB theory differs from the IIA theory from the point of view taken in this paper and discuss the connection of our IIA results with previous work on inflation in IIB models. We discuss our results in section 6. More details regarding the kinetic energy and potential energy are provided in appendix A.

2. Type IIA compactifications

We investigate large volume and small string coupling compactifications where it is valid to perform computations using supergravity. We study the 10-dimensional type IIA supergravity theory, where we include conventional NS-NS and R-R field strengths, as well as D6-branes and O6-planes:

$$S = \frac{1}{2\kappa_{10}^2}\int d^{10}x\sqrt{-g}\,e^{-2\phi}\left(R + 4(\partial_\mu\phi)^2 - \frac{1}{2}|H_3|^2 - e^{2\phi}\sum_p |F_p|^2\right)$$
$$-\mu_6\int_{\text{D6}} d^7\xi\sqrt{-g}\,e^{-\phi} + 2\mu_6\int_{\text{O6}} d^7\xi\sqrt{-g}\,e^{-\phi} \tag{2.1}$$

where R is the 10-dimensional Ricci scalar, ϕ is the scalar dilaton field, H_3 is the NS-NS 3-form field strength that is sourced by strings, F_p are the R-R p-form field strengths

($p = 0, 2, 4, 6$) that are sourced by branes, $\kappa_{10}^2 = 8\pi G_{10}$ is the gravitational strength in 10-dimensions, and μ_6 ($-2\mu_6$) is the D6-brane (O6-plane) charge and tension. We have set all fermions to zero as we are interested in solutions with maximal space-time symmetry. There are also Chern-Simons contributions to the action. These are essentially topological, and are independent of the dilaton as well as the overall scale of the metric (in string frame). We expect that as in [43] the contribution to the action from the Chern-Simons terms will vanish on-shell,[3] so that we need not consider it further here. Although there are some subtle questions regarding the definition of orientifolds in these massive IIA backgrounds and whether these backgrounds can be described in a weak coupling string expansion [48, 49], these compactifications seem to be described adequately in the 4-dimensional supergravity formalism of [42], corresponding to the 10-dimensional massive supergravity analysis when the sources are uniformly distributed in the compactification space.

2.1 Compactification

We now perform a Kaluza-Klein compactification of this theory from 10 dimensions to 4 dimensions. Let us first focus on the gravity sector. Assuming that we can neglect any dependence of the Ricci scalar on the compact space coordinates, we can integrate over the compact space, giving

$$\int d^{10}x\sqrt{-g_{10}}\, e^{-2\phi}R = \int d^4x\sqrt{-g_4}\,\mathrm{Vol}\, e^{-2\phi}R \tag{2.2}$$

where Vol is the 6-dimensional volume of the compact space. The volume, dilaton, and all the other fields that describe the size and shape of the compact space are scalar fields in the 4-dimensional description, known as moduli. In addition to kinetic energy terms, the remaining terms in the supergravity action, when reduced to 4 dimensions, describe a potential function V which depends on the moduli and fluxes in any given model.

The key observation of this paper is that by studying the dependence of the potential energy function V on only two of the moduli, we can learn a great deal about the structure of the potential relevant for the possibility of inflation. We define the volume modulus of the compact space ρ and the dilaton modulus τ by[4]

$$\rho \equiv (\mathrm{Vol})^{\frac{1}{3}}, \quad \tau \equiv e^{-\phi}\sqrt{\mathrm{Vol}}. \tag{2.3}$$

While V depends on all moduli, we can explore the behavior of this function on the whole moduli space by considering 2-dimensional slices of the moduli space where all moduli other than τ and ρ are fixed. By showing that V has a large gradient in the τ-ρ plane on every slice wherever V is positive, we will be able to rule out inflation on the *entire* moduli space, regardless of which fields we would like to identify as the inflaton.

[3] More precisely, integrating out the 4-dimensional non-dynamical field dC_3 as a Lagrange multiplier gives an equation which must be satisfied by the axions, but the Chern-Simons terms do not otherwise affect the 4-dimensional potential.

[4] Note that we have defined the volume modulus in the string frame, since this definition relates to the Kähler moduli in the IIA theory. This differs from the conventional definition of Kähler moduli in analogous IIB orientifolds, where the Einstein frame metric is used in defining the chiral multiplets.

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In order to bring the gravity sector into canonical form, we perform a conformal transformation on the metric to the so-called Einstein frame:

$$g^E_{\mu\nu} \equiv \frac{\tau^2}{\bar{m}_{\rm P}^2\,\kappa_{10}^2} g^4_{\mu\nu}, \tag{2.4}$$

where $\bar{m}_{\rm P} = 1/\sqrt{8\pi G} \approx 2 \times 10^{18}\,\text{GeV}$ is the (reduced) Planck mass and G is the 4-dimensional Newton constant. By re-expressing R in terms of a 4-dimensional Ricci scalar and performing the conformal transformation, one finds that the gravity sector is canonical and that the fields ρ and τ carry kinetic energy that is diagonal. Although ρ and τ do not have canonical kinetic energies, they are related to fields which do:

$$\hat{\rho} \equiv \sqrt{\frac{3}{2}}\,\bar{m}_{\rm P} \ln \rho, \quad \hat{\tau} \equiv \sqrt{2}\,\bar{m}_{\rm P} \ln \tau. \tag{2.5}$$

Altogether we obtain the effective Lagrangian in 4 dimensions in the Einstein frame

$$\mathcal{L} = \frac{1}{16\pi G} R_E - \left[\frac{1}{2}(\partial_\mu \hat{\rho})^2 + \frac{1}{2}(\partial_\mu \hat{\tau})^2 + \ldots\right] - V \tag{2.6}$$

where the Ricci scalar R_E and all derivatives are defined with respect to the Einstein metric. The dots indicate further kinetic energy terms from all the other fields of the theory associated with the compactification (ϕ_i): the so-called Kähler moduli, complex structure moduli, and axions.[5] The important point is that their contributions will always be *positive*. More details are provided in appendix A.1. All contributions from the field strengths and D-branes & O-planes from eq. (2.1) are described by some potential energy function V.

2.2 The slow-roll condition

From this action we can derive the slow-roll conditions on the potential V for inflation. In order to write down the conditions in detail we would need to know the precise form of the kinetic energy with respect to all the moduli. This can be done cleanly in the 4-dimensional supergravity formalism, as mentioned in appendix A.1. The important point here is that the first slow-roll parameter ϵ involves partial derivatives of the potential with respect to each direction in field space, and that the contribution from $\hat{\rho}$, $\hat{\tau}$, ϕ_i is non-negative. In fact it is roughly the square of the gradient of $\ln V$. The contributions from $\hat{\rho}$ and $\hat{\tau}$ thus give the following lower bound:

$$\epsilon \geq \frac{\bar{m}_{\rm P}^2}{2}\left[\left(\frac{\partial \ln V}{\partial \hat{\rho}}\right)^2 + \left(\frac{\partial \ln V}{\partial \hat{\tau}}\right)^2\right]. \tag{2.7}$$

A *necessary* condition for inflation is $\epsilon \ll 1$ with $V > 0$. We will now prove that this condition is impossible to satisfy in this IIA framework, at large volume and weak coupling where our calculations apply.

[5] The axions arise from zero-modes of the various gauge fields.

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3. No-go theorem

Having discussed the gravity and kinetic energy sector of the theory, let us now describe the form of the potential energy function V. By using the bound on ϵ we will then prove that inflation is forbidden for any Calabi-Yau compactification of type IIA string theory in the large volume and small string coupling limit, when only conventional NS-NS and R-R field strengths, D6-branes and O6-planes are included.

3.1 Potential energy

The potential energy arises from the dimensional reduction of the terms in (2.1) associated with the various field strengths H_3 & F_p ($p = 0, 2, 4, 6$) and the D6-branes & O6-planes. Let us focus on some field strength F_p. Such a field can have a nonvanishing integral over any closed p-dimensional internal manifold (homology cycle) of the compact space, and must satisfy a generalized Dirac charge quantization condition:

$$\int_\Sigma F_p \propto f_\Sigma, \tag{3.1}$$

where f_Σ is an integer associated with number of flux quanta of F_p through each p-dimensional cycle Σ. By choosing different values for f_Σ over a basis set of p-cycles, one obtains a *landscape* of possible allowed potential energy functions V.

The energy arising from a p-form flux F_p comes from a term in (2.1) proportional to $|F_p|^2$; since the p-form transforms as a covariant p-tensor (i.e., has p lower indices), we contract with p factors of $g_6^{\mu\nu}$, so that

$$|F_p|^2 \propto \rho^{-p} \tag{3.2}$$

in the string frame. By including the appropriate factors of the volume and dilaton from the compactification and performing the conformal transformation to the Einstein frame, we have the following contributions to V:

$$\begin{aligned} V_3 &\propto \rho^{-3}\tau^{-2} \quad \text{for } H_3, \\ V_p &\propto \rho^{3-p}\tau^{-4} \quad \text{for } F_p. \end{aligned} \tag{3.3}$$

We also need the contribution from D6-branes and O6-planes. In the Einstein frame they scale as

$$\begin{aligned} V_{\text{D6}} &\propto \tau^{-3} \quad \text{for D6-branes}, \\ V_{\text{O6}} &\propto -\tau^{-3} \quad \text{for O6-planes}, \end{aligned} \tag{3.4}$$

where we have indicated that O6-planes provide a negative contribution, while all others are positive.

Altogether, we have the following expression for the scalar potential in 4-dimensions

$$\begin{aligned} V &= V_3 + \sum_p V_p + V_{\text{D6}} + V_{\text{O6}} \\ &= \frac{A_3(\phi_i)}{\rho^3\tau^2} + \sum_p \frac{A_p(\phi_i)}{\rho^{p-3}\tau^4} + \frac{A_{\text{D6}}(\phi_i)}{\tau^3} - \frac{A_{\text{O6}}(\phi_i)}{\tau^3}. \end{aligned} \tag{3.5}$$

Here we have written the various coefficients as A_j (≥ 0), which in general are complicated functions of all the other fields of the theory ϕ_i, namely the remaining set of Kähler moduli, complex structure moduli, and axions. The coefficients A_j also depend on the choice of flux integers f_Σ. This means that V is in general a function of hundreds of fields, for each of the exponentially large number of infinite families of possible flux combinations on each of the many available Calabi-Yau manifolds. We have simply described its dependence on two of the fields: ρ and τ. An alternative proof of eq. (3.5) from the perspective of the 4-dimensional supergravity formalism is given in appendix A.2. As discussed in ref. [43] this potential ensures that there exist special points in field space which stabilize all the geometric moduli and many axions. We are interested, however, in exploring the full moduli space in search of inflation.

3.2 Proof of no-go theorem

We find that for any potential of the class we have constructed so far, inflation is impossible anywhere in the moduli space. The proof of this result is quite simple. The key is to observe that the potential in eq. (3.5) satisfies

$$\begin{aligned} -\rho\frac{\partial V}{\partial \rho} - 3\tau\frac{\partial V}{\partial \tau} &= 9V + \sum_p p\, V_p \\ &\geq 9V, \end{aligned} \tag{3.6}$$

where the inequality comes from the fact that all $V_p \geq 0$ ($A_p \geq 0$) and $p \geq 0$. Now, assuming we are in a region where $V > 0$, which is necessary for inflation, we can divide both sides by V and rewrite this inequality in terms of $\hat{\rho}$ and $\hat{\tau}$ as

$$\bar{m}_{\rm P}\left|\sqrt{\frac{3}{2}}\left(\frac{\partial \ln V}{\partial \hat{\rho}}\right) + 3\sqrt{2}\left(\frac{\partial \ln V}{\partial \hat{\tau}}\right)\right| \geq 9. \tag{3.7}$$

This implies that it is impossible for both terms in eq. (2.7) to be simultaneously small. Specifically, comparing this inequality with eq. (2.7), we see that $\sqrt{2\epsilon}/\bar{m}_{\rm P}$ is the distance to the origin on the plane spanned by $(\partial \ln V/\partial\hat{\rho}, \partial \ln V/\partial\hat{\tau})$, while a sloped band around the origin is forbidden, implying the existence of a lower bound on ϵ. By minimizing ϵ subject to the constraint (3.7), we find the bound on the slow-roll parameter ϵ to be

$$\epsilon \geq \frac{27}{13} \quad \text{whenever } V > 0. \tag{3.8}$$

Hence both inflation and de Sitter vacua are forbidden everywhere in field space.

Indeed in any vacuum $\partial V/\partial\rho = \partial V/\partial\tau = 0$, so eq. (3.6) implies $V = -(\sum pV_p)/9$. By assuming $V_p > 0$ for at least one of $p = 2$, 4, or 6, then Minkowski vacua are forbidden also.[6] This type of relation was used in [43] to show that vacua in a specific IIA compactification must be anti-de Sitter, and in [50] to rule out a simple F-term uplift of this model; here we

[6] We almost certainly need $V_p > 0$ for at least one of $p = 2$, 4, or 6 in order to be in the large volume and small string coupling limit, since the 3-form, 0-form, and number of D6/O6 planes are tightly constrained by a tadpole condition.

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have shown that this relation holds very generally in IIA compactifications, and furthermore rules out inflation anywhere on moduli space for potentials containing only terms of the form (3.5). For stabilized compactifications, the implication is that the field vector always undergoes fast rolling from a region with $V > 0$ towards an anti-de Sitter vacuum.

The no-go theorem we have derived here can be interpreted as defining a necessary condition for inflation in IIA models: *in order to inflate, a IIA compactification must contain some additional structure beyond that considered so far which gives a term in the potential V whose scaling leads to a term on the r.h.s. of (3.6) with a coefficient less than 9 if positive or greater than 9 if negative.* In the next section we turn to a discussion of specific types of structure which can realize this necessary condition for inflation.

4. Evading the no-go theorem

In the analysis which led to the preceding no-go theorem for inflation in IIA compactifications, we allowed a specific set of ingredients in the IIA models considered. Following [43], we included NS-NS 3-form flux, R-R fluxes, D6-branes and O6-planes, which are sufficient to stabilize all geometric moduli. Because D6-branes and anti-D6-branes give terms to V which scale in the same way, this no-go theorem rules out brane-antibrane inflation as well as inflationary models using other moduli as the inflaton. Corrections to 10-dimensional supergravity which arise at small compactification volume may evade our no-go result, so one approach to finding a IIA compactification with inflation is to include finite volume corrections to the Kähler potential. There are also other structures we can include in a compactification besides those already mentioned which can evade the no-go result. In this section we consider other possibilities which give rise to terms in the potential V which have scaling coefficients violating (3.6). This may give some guidance for where to look to find IIA compactifications with inflation and/or de Sitter vacua. Note that many of the structures suggested go beyond the range of compactifications which are so far well understood in string theory.

4.1 Various ingredients

One obvious possibility is to include Dp-branes and Op-planes of dimensionality other than $p = 6$. The scaling of the resulting contributions to V are as follows

$$\begin{aligned} V_{\mathrm{Dp}} &\propto \rho^{\frac{p-6}{2}}\tau^{-3} \qquad \text{for D}p\text{-branes,} \\ V_{\mathrm{Op}} &\propto -\rho^{\frac{p-6}{2}}\tau^{-3} \qquad \text{for O}p\text{-planes.} \end{aligned} \tag{4.1}$$

In these cases, the right hand of eq. (3.6) is $(12 - p/2)V_{\mathrm{Dp/Op}}$, so the no-go theorem applies to Dp-branes with $p \leq 6$ and Op-planes with $p \geq 6$, but is evaded otherwise. Since the branes must extend in all non-compact directions of space-time[7] in IIA we can only consider $p = 4, 8$. Wrapping such a brane, however, would either require a compactification with non-trivial first homology class H_1 or finite π_1, or perhaps require the use of so-called coisotropic 8-branes [51, 52]. These branes carry charge and so would generate additional

[7] Otherwise they would describe localized excitations in an asymptotic vacuum not including these branes.

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tadpoles which would need to be cancelled, as well as potentially breaking supersymmetry. Compactifications with branes of this type have not been studied extensively in the string theory literature, but it would be interesting to investigate this range of possibilities further.

Another possibility is to include more general NS-NS fluxes, such as geometric fluxes *a la* Scherk and Schwarz [53] or non-geometric fluxes [54]. Geometric fluxes parameterize a "twisting" away from Calabi-Yau topology, generalizing the notion of a twisted torus [53, 55]. These fluxes are associated with a metric with curvature on the compact space. Geometric fluxes arise under T-duality or mirror symmetry when an NS-NS 3-form H has a single index in a dualized direction [56]. Further T-dualities generate non-geometric fluxes, described in [57] through the sequence $T : H_{abc} \to f^a_{bc} \to Q^{ab}_c \to R^{abc}$, where f parameterize geometric fluxes, Q are locally geometric but globally non-geometric fluxes which can in some cases be realized in the language of "T-folds" [58, 59], and R parameterize fluxes associated with compactifications which are apparently not even locally geometric. These general NS-NS fluxes and the associated compactifications are still poorly understood. From the T-duality picture, however, it is straightforward to determine the scalings of these 3 types of fluxes. Each T-duality inverts the size of a dimension of the compactification, replacing a factor of ρ^{-1} in the scaling with ρ, so we have

$$\begin{aligned} V_{\rm f} &\propto \pm\rho^{-1}\tau^{-2} \quad && \text{for geometric (f) flux,} \\ V_{\rm Q} &\propto \pm\rho\,\tau^{-2} \quad && \text{for Q flux,} \\ V_{\rm R} &\propto \pm\rho^{3}\tau^{-2} \quad && \text{for R flux.} \end{aligned} \tag{4.2}$$

Applying the linear operator of eq. (3.6) to each of these terms gives a right hand side of $7V_{\rm f}$, $5V_{\rm Q}$, and $3V_{\rm R}$, respectively. So although our no-go theorem applies when such terms are negative, it is evaded if any of these contributions are positive. Among these fluxes, the best understood are geometric fluxes, which are realized in many simple compactifications such as twisted tori [44, 53, 55]. Compactification on spaces with these fluxes (and other ingredients) is studied in the forthcoming paper by Silverstein [47], where it is shown that de Sitter vacua can indeed be realized in such backgrounds. This is a promising place to look for string inflation models. Note, however, that general NS-NS fluxes cannot in general be taken to the large volume limit. For example, fluxes of the Q type involve a T-duality inverting the radius of a circle in a fiber when a circle in the base is traversed. Thus, somewhere the size of the fiber must be sub-string scale. This makes solutions of the naive 4-dimensional supergravity theory associated with flux compactifications such as those found in [60] subject to corrections from winding modes and also to uncontrolled string theoretic corrections if curvatures become large.

Another possible ingredient which can be added to the IIA compactification models are NS 5-branes; these are the magnetic duals of the string. Such objects are non-perturbative and carry tension in 10 dimensions that scales as $g_s^{-2} = e^{-2\phi}$. While they backreact more significantly than D-branes and so are not as simple to describe at the supergravity level, their presence can be captured by adding a term of the form

$$-\mu_5 \int_{\rm NS5} d^6\xi \sqrt{-g}\, e^{-2\phi} \tag{4.3}$$

to the action of eq. (2.1). The compactification and conformal transformation yields a new term that scales as

$$V_{\rm NS5} \propto \rho^{-2}\tau^{-2}, \tag{4.4}$$

yielding a right hand side of eq. (3.6) of $8V_{\rm NS5}$, hence evading the no-go theorem. To satisfy tadpole cancellation, one would probably wish to add metastable pairs of separated NS 5-branes and anti-NS 5-branes, wrapping distinct isolated curves in the same homology class. Such configurations have been a focus of study in the dual IIB theory in recent works, starting with [61].

4.2 An illustration

In this section we illustrate how the above ingredients may be useful from the point of view of building de Sitter vacua and inflation. We focus the discussion on NS 5-branes, which appear particularly promising. We will not attempt an explicit construction, since that would take us beyond the scope of this article. Our goal is only to show that simple, available ingredients in the IIA theory have energy densities which scale with the volume and dilaton moduli in a way which suffices to overcome our no-go theorem, which was based purely on scalings of energy densities. This should act as a guide to model building, but should be taken in the heuristic spirit it is offered.

Using the potential of eq. (3.5) and adding to it a (necessarily) positive term from NS 5-branes wrapping 2-cycles, we obtain[8]

$$V = \frac{A_3(\phi_i)}{\rho^3\tau^2} + \sum_p \frac{A_p(\phi_i)}{\rho^{p-3}\tau^4} - \frac{A_{\rm O6}(\phi_i)}{\tau^3} + \frac{A_{\rm NS5}(\phi_i)}{\rho^2\tau^2}. \tag{4.5}$$

Let us now streamline V and focus on the most important features of this setup; we set $A_2 = A_6 = 0$, and expect the remaining coefficients to scale with fluxes and numbers of planes and branes as

$$A_3 \sim h_3^2, \;\; A_0 \sim f_0^2, \;\; A_4 \sim f_4^2, \;\; A_{\rm O6} \sim N_{\rm O6}, \;\; A_{\rm NS5} \sim g(\omega) N_{\rm NS5}, \tag{4.6}$$

where we have introduced a function g of the modulus ω associated with the 2-cycle wrapped by each NS 5-brane.

There is a tadpole constraint that the charge on the O6-plane must be balanced by the fluxes from H_3 and F_0, i.e., $h_3 f_0 \sim -N_{\rm O6}$. We use this to eliminate h_3. Now since it has no effect on the kinetic energy, let us rescale our fields as: $\rho \to \rho\sqrt{|f_4/f_0|}$ and $\tau \to \tau\sqrt{|f_0 f_4^3|}/N_{\rm O6}$. Then we have

$$V = V_{\rm flux}\left[\frac{B_3(\phi_i)}{\rho^3\tau^2} + \sum_p \frac{B_p(\phi_i)}{\rho^{p-3}\tau^4} - \frac{B_{\rm O6}(\phi_i)}{\tau^3} + \frac{B_{\rm NS5}(\phi_i)}{\rho^2\tau^2}\right] \tag{4.7}$$

with $V_{\rm flux} \equiv N_{\rm O6}^4/\sqrt{|f_0^3 f_4^9|}$, and

$$B_3 \sim 1, \;\; B_0 \sim 1, \;\; B_4 \sim 1, \;\; B_{\rm O6} \sim 1, \;\; B_{\rm NS5} \sim c(\omega) \equiv g(\omega) N_{\rm NS5}\sqrt{|f_0^3 f_4|}/N_{\rm O6}^2. \tag{4.8}$$

[8] We have set $A_{\rm D6} = 0$ as it adds little to the analysis.

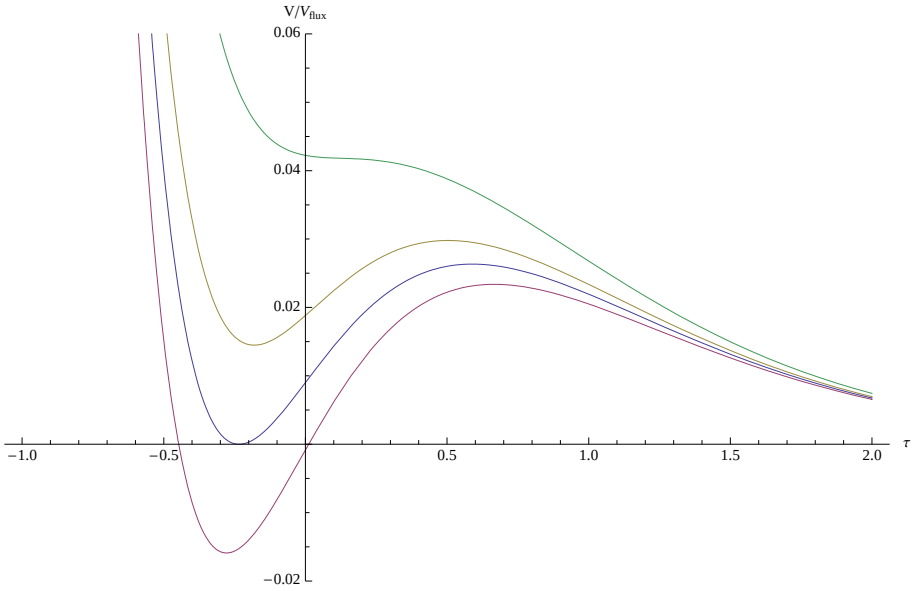

Figure 1: The potential $V(\hat{\tau})/V_{\rm flux}$ with ($\bar{m}_{\rm P}=1$) $B_3=1/4$, $B_0=1/4$, $B_4=3/8$, $B_{\rm O6}=2$, and ρ satisfying $\partial V/\partial\rho=0$. From bottom to top, the curves correspond to the following choices of c: 2.183 (anti-de Sitter), 2.205 (Minkowski), 2.227 (de Sitter), and 2.280 (inflection), respectively.

So the shape of the potential is essentially controlled by the parameter c.

Let us make some comments on the value of c which determines the contribution of the NS 5-brane. In the large f_4 flux limit of [43], c is parametrically larger than all other contributions, including that from the 3-form flux to which we compare: the NS 5-brane contribution has the same scaling with τ, but scales more slowly to zero as $\rho\to\infty$. Therefore, one would expect that one needs *small* $N_{\rm NS5}$ to push V up without causing a runaway to infinite volume. As we will see, we need c to be fined tuned and $\mathcal{O}(1)$. An analogous situation occurs in IIB string theory with compactifications involving anti-D3 branes and non-perturbative volume stabilization. There the presence of strong warping allows one to construct states where the anti-D3 energy density is exponentially suppressed, naturally providing a small coefficient to the perturbation of the energy density [62]. This plays an important role in de Sitter constructions in that context [63], and one might expect an analogous mechanism (involving large warping or a very small cycle) could similarly dynamically explain a small $g(\omega)$ to compensate large f_4 in the IIA context. We could also imagine a compactification where $N_{\rm O6}$ is very large to achieve the same result.

In any case, by treating c as a continuous parameter and ignoring the dynamics of all other moduli, we can obtain a meta-stable de Sitter vacuum. We set ($\bar{m}_{\rm P}=1$) $B_3=1/4$, $B_0=1/4$, $B_4=3/8$, $B_{\rm O6}=2$, and stabilize ρ by satisfying $\partial V/\partial\rho=0$. In figure 1 we plot $V=V(\hat{\tau})$ of eq. (4.7) for different choices of $B_{\rm NS5}=c$. We find that there are critical values of c: for a Minkowski vacuum $c_M\approx 2.205$ and for a point of inflection $c_I\approx 2.280$. For $c<c_M$ an anti-de Sitter vacuum exists, for $c_M<c<c_I$ a de Sitter vacuum exists, and

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for $c > c_I$ no vacuum exists. For c close to c_M^+ we expect the de Sitter vacuum lifetime to be long, as in the KKLT meta-stable vacuum of type IIB [63].

If one could realize such a construction explicitly in a controlled regime of the IIA theory, then we anticipate that the opportunities to realize inflation will be greatly enhanced. For example, the local maximum in figure 1 may be useful since $\epsilon \to 0$ there. It will often suffer, however, from the so-called η problem; the second slow-roll parameter (which measures the second derivative of the potential) will typically be large and negative. It is possible that, for example, by moving in another transverse direction at the hill-top one could build some form of hybrid-inflation. We do note that by choosing $c = c_I$ we would immediately solve the η problem and have inflection-point-inflation as advocated in [64]. In this situation, however, inflation would finish with two difficulties; runaway moduli and little to no reheating. A different approach is to simply fix the volume and the dilaton at a de Sitter (or Minkowski) minimum with high mass and strictly use other lighter moduli to drive inflation in transverse directions. Scenarios such as N-flation may be possible here [22].

5. Type IIB compactifications

In this section we discuss the relationship between the results we have derived here for IIA string theory and previous work on inflation in type IIB string theory. Although the type IIA and IIB string theories are related through T-duality, this duality acts in a complicated way on many of the ingredients used in constructing flux compactifications. The basic IIB flux compactifications of [65] (see also [66 – 68]) involve H-flux, F_3-flux, D3-branes and O3-planes. T-duality/mirror symmetry on such compactifications transforms the NS-NS H-flux into a complicated combination of H-flux, geometric flux and non-geometric flux which generically violates the restrictions needed in the no-go theorem we have proven here. Conversely, the backgrounds which we have proven here cannot include inflation are T-dual/mirror to complicated IIB backgrounds with geometric and non-geometric fluxes. Furthermore, the volume modulus used in our analysis is dual to a complex structure modulus in the IIB theory which is difficult to disentangle from the other moduli, so that proving the analogous no-go theorem in IIB, on the exotic class of backgrounds where it is relevant, would be quite difficult without recourse to duality.

Despite these complications, we can easily explain why standard IIB flux vacua behave so differently with respect to potential constructions of de Sitter space and inflation as seen through the methods of this paper. For IIB flux vacua, the basic ingredients of H-flux, F_3-flux, D3-branes and O3-planes give contributions to the 4-dimensional potential which scale as

$$\frac{e^{2\phi}}{\rho^6}, \quad \frac{e^{4\phi}}{\rho^6}, \quad \frac{e^{3\phi}}{\rho^6}, \quad -\frac{e^{3\phi}}{\rho^6} \tag{5.1}$$

respectively.[9] Thus, the scaling equation (3.6) is not the appropriate equation for gaining

[9] This ρ scales as $\mathrm{Vol}^{1/3}$ defined in the string frame, as in section II, and should not be confused with the ρ modulus of e.g. [65], which scales as $\mathrm{Vol}^{2/3}$ defined in the Einstein frame.

useful information about the cosmological structure. Instead, we simply have

$$-\rho\frac{\partial V(\phi,\rho)}{\partial\rho} = 6V(\phi,\rho)\,. \tag{5.2}$$

This shows immediately that any classical vacuum must have $V = 0$, as is well known from the tree-level no-scale structure of this class of models [65]. When the dilaton and complex structure moduli are chosen to fix $V = 0$ then there is a classical flat direction. When such moduli cannot be chosen then the (positive) potential causes a runaway to large volume. In such a simple setting, one can show that $\epsilon \geq 3$ whenever $V > 0$, but this *should not* be viewed as a serious obstacle to realizing inflation or de Sitter vacua. This is because the classical flat directions along which $V = \partial V/\partial\rho = 0$ can be lifted by including quantum contributions (or other fluxes etc) to the potential, and then the naive bound on ϵ is irrelevant. Typically, non-perturbative corrections to the superpotential are included to stabilize the Kähler moduli. This should be contrasted with the classical stabilization in IIA.[10]

In this general IIB setting, starting with the no-scale vacuum and including various corrections to achieve de Sitter, many inflationary models have been proposed. The basic strategy, starting with [29], has usually been to stabilize the dilaton and volume moduli at a high scale, and inflate at a lower energy scale. Then, the dilaton and volume contributions to ϵ (which were the focus of our no-go theorem in the IIA context) are simply absent. The most explicit models to date appear in [31], though one should consult the reviews [36 – 40] for a much more extensive list of approaches and references.

6. Discussion

In this paper we have demonstrated that a large class of flux compactifications of type IIA string theory cannot give rise to inflation in the regime of moduli space where we have parametric control of the potential. This result applies to large-volume, weak coupling compactifications on arbitrary Calabi-Yau spaces with NS-NS 3-form flux, general R-R fluxes, D6-branes and O6-planes. These ingredients are arguably the most well understood in IIA compactifications. The no-go theorem of section 3 applies in particular to the $T^6/\mathbb{Z}_3^2$ orientifold model of ref. [43] and the $T^6/\mathbb{Z}_4$ orientifold model of ref. [69], and explains the numerical results of ref. [46] which suggested that inflation is impossible in these models. The no-go theorem we have derived here, however, applies to all other Calabi-Yau compactifications of this general type as well. So the simplest part of the IIA flux compactification landscape does not inflate. This implies the following constraint: the portions of the landscape that are possibly relevant to phenomenology will necessarily involve interplay of more diverse ingredients, as has also been found in the IIB theory.

We emphasize that while our derivation has only involved two moduli (the volume ρ and the dilaton τ) we are not assuming that either of those moduli necessarily play the role of the inflaton. Instead, inflation by *any* modulus (or brane/anti-brane) is always spoiled due to the fast-roll of ρ and/or τ. This follows because a necessary condition for slow-roll

[10] Note that in the absence of fluxes, branes and orientifolds, mirror symmetry relates supersymmetric IIA and IIB compactifications. In this case $V = 0$ exactly.

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inflation is that the potential be flat in *every* direction in field space, as quantified by ϵ. In fact because the first slow-roll parameter is so large $\epsilon \geq \frac{27}{13}$, there can never be many e-foldings, even ruling out so-called *fast-roll inflation* [70]. We point out that this result is slightly non-trivial because it requires analyzing *both* ρ and τ, as in eq. (3.6), and cannot be proven by focussing on only one of them.

A simple corollary of our result is that no parametrically controlled de Sitter vacua exist in such models. We emphasize, though, that proving the non-existence of inflation is a much stronger statement than proving the non-existence of de Sitter vacua. In particular we can imagine a priori a scenario where, although the vacuum is anti-de Sitter (or Minkowski) or there is no vacuum at all, inflation is still realized somewhere in a region where $V > 0$. We find, however, that this does not occur. It is intriguing that the proof of the non-existence of inflation is so closely related to that of the non-existence of de Sitter. This *suggests* that there may be a close connection between building de Sitter vacua and realizing inflation.

Our result can be interpreted as giving a necessary condition for inflation in IIA models. To have inflation, some additional structure must be added which gives rise to a potential term in 4 dimensions with scaling such that $-\rho\partial V/\partial\rho - 3\tau\partial V/\partial\tau = \alpha V$ with a coefficient $\alpha < 9$ for a positive contribution and $\alpha > 9$ for a negative contribution. We described various ingredients which give rise to such terms; compactifications including these ingredients may be promising places to look for inflationary models. Some of these ingredients take us outside the range of string compactifications which are understood from a perturbative/supergravity point of view. Among the possibilities which evade the assumptions made in deriving the no-go theorem are other NS-NS fluxes, such as geometric and non-geometric fluxes. It is currently difficult to construct models with such generic fluxes in a regime that is under control, but progress in this direction has been made in [47], where IIA de Sitter vacua are found using a specific set of geometrical fluxes and other ingredients. Another promising direction which we have indicated here (also incorporated in [47]) is to include NS 5-branes and anti-NS 5-branes on 2-cycles in the Calabi-Yau. More work is needed to find explicit models where these branes are stabilized in a regime allowing de Sitter vacua and inflation, but this does not seem to be impossible or ruled out by any obvious considerations. In addition to the mechanisms we have discussed, there are probably other structures (e.g., D-terms [71, 72]) which violate the conditions of the no-go theorem.

Including any of these ingredients does not guarantee that inflation will be realized. It may be the case that a slightly more general no-go theorem for inflation exists with certain combinations of additional ingredients. This may follow from studying other moduli, since any field can ruin inflation by fast-rolling. It may also be that with more work an elegant realization of inflation and de Sitter vacua can be found in type IIA string theory. This deserves further investigation.

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A. 4-dimensional $\mathcal{N} = 1$ supergravity

It was shown in [42] that with the ingredients used in section 3, the dimensional reduction of massive IIA supergravity can be described in the language of a 4-dimensional $\mathcal{N} = 1$ supergravity theory in terms of a Kähler potential and superpotential.

A.1 Kinetic energy from Kähler potential

The authors of ref. [42] showed that in the large volume limit the Kähler potential is given by $K = K^K + K^Q$ where

$$K^K = -\bar{m}_{\rm P}^2 \ln\left(\frac{4}{3}\kappa_{abc}v^a v^b v^c\right), \tag{A.1}$$

$$K^Q = -2\bar{m}_{\rm P}^2 \ln\Big(2\,{\rm Im}(CZ_\lambda){\rm Re}(Cg_\lambda) - 2\,{\rm Re}(CZ_k){\rm Im}(Cg_k)\Big) \tag{A.2}$$

Here v^a are the Kähler moduli, κ_{abc} are the triple-intersection form constants, the set of Z and g are the co-ordinates in some basis of a holomorphic 3-form that describes the complex structure moduli, and C is the "compensator" which incorporates the dilaton. If the set of complex moduli is denoted ψ^i, then the kinetic energy is given by

$$T = -K_{i\bar{j}}\partial_\mu\psi^i\partial^\mu\psi^{\bar{j}} \tag{A.3}$$

with corresponding first slow-roll parameter

$$\epsilon = \bar{m}_{\rm P}^2\frac{K^{i\bar{j}}V_iV_{\bar{j}}}{V^2}. \tag{A.4}$$

Let us focus on the Kähler contribution. We write the Kähler moduli as $\psi^a = a^a + i\,v^a$, so the kinetic energy is given by

$$T^K = -\frac{1}{4}\frac{\partial^2 K^K}{\partial v^a\partial v^b}\left(\partial_\mu v^a\partial^\mu v^b + \partial_\mu a^a\partial^\mu a^b\right). \tag{A.5}$$

Now we change coordinates from v^a to $\{\rho, \gamma^a\}$ as follows:

$$v^a = \rho\,\gamma^a, \;\text{ with }\; \kappa_{abc}\gamma^a\gamma^b\gamma^c = 6, \tag{A.6}$$

so $\mathrm{Vol} = \rho^3$. Then using $\partial_\mu(\kappa_{abc}\gamma^a\gamma^b\gamma^c) = 0$, we obtain:

$$T^K = -\,\bar{m}_{\rm P}^2\left[\frac{3(\partial_\mu\rho)^2}{4\,\rho^2} - \frac{1}{4}\kappa_{abc}\gamma^c\partial_\mu\gamma^a\partial^\mu\gamma^b + \frac{\kappa_{acd}\gamma^c\gamma^d\kappa_{bef}\gamma^e\gamma^f - 4\kappa_{abc}\gamma^c}{16\,\rho^2}\partial_\mu a^a\partial^\mu a^b\right]. \quad \text{(A.7)}$$

By switching from ρ to $\hat{\rho}$, we see that the first term is precisely the kinetic energy for $\hat{\rho}$. The remaining kinetic energy terms for γ^a and a^a are block diagonal (there are no cross terms involving $\partial_\mu\rho\,\partial^\mu\gamma^a$ etc), and this has an important consequence: We know that in the physical region the total kinetic energy must be positive, so *each* of the above 3 terms must be positive. Hence, $T^K = -(\partial_\mu\hat{\rho})^2/2 + \text{positive}$.

For the complex structure/dilaton sector the procedure is similar, although more subtle. In (A.2) the expression for K^Q is not a completely explicit function of the moduli; although $\mathrm{Re}(CZ_k)$ and $\mathrm{Re}(Cg_k)$ are explicitly half of the complex structure moduli, $\mathrm{Im}(CZ_k)$ and $\mathrm{Im}(Cg_k)$ are only *functions* of the remaining complex structure moduli. Nevertheless, the kinetic term is again block diagonal. To see this, note that the compensator C, and hence all moduli in this sector, are proportional to τ. Furthermore, the Z and g are constrained to the surface: $K^Q = -2\bar{m}_{\rm P}^2\ln(\tau^2)$. This is analogous to the Kähler sector. Without going through the details here, we find $T^Q = -(\partial\hat{\tau})^2/2 + \text{positive}$. In fact we know this *must* be true from the 10-dimensional point of view; the dilaton modulus is inherited directly from 10 dimensions, and so cannot possibly give rise to mixed kinetic terms with the complex structure moduli in 4 dimensions.

A.2 Potential energy from superpotential

From [42] the superpotential in the IIA theory is given by $W = W^K + W^Q$ where

$$W^K = f_6 + f_{4a}t^a + \frac{1}{2}f_{2a}\kappa_{abc}t^bt^c - \frac{f_0}{6}\kappa_{abc}t^at^bt^c, \quad \text{(A.8)}$$

$$W^Q = \left(\bar{h}_\lambda\xi^\lambda - h_k\xi^k\right) + 2i\left(\bar{h}_\lambda\mathrm{Re}(Cg_\lambda) - h_k\mathrm{Re}(CZ_k)\right). \quad \text{(A.9)}$$

We will not explain all the details of this here; the interested reader is pointed to refs. [42, 43]. For our present purposes it suffices to note that $\mathrm{Im}(t^a) = v^a \propto \rho$ and $\mathrm{Im}(W^Q) \propto \tau$. Hence, the superpotential is cubic in ρ and linear in τ.

From the supergravity formula for the Einstein frame potential

$$V = e^{K/\bar{m}_{\rm P}^2}\left(D_iWK^{i\bar{j}}\overline{D_jW} - 3\frac{|W|^2}{\bar{m}_{\rm P}^2}\right), \quad \text{(A.10)}$$

we easily infer the dependence on ρ and τ. Firstly, since the constraints on γ^a and complex structure imply that $K = -\bar{m}_{\rm P}^2\ln(8\rho^3\tau^4)$, the pre-factor scales as

$$e^{K/\bar{m}_{\rm P}^2} \propto \rho^{-3}\tau^{-4}. \quad \text{(A.11)}$$

Also, the scaling contributions from the parenthesis in eq. (A.10), which is roughly $|W|^2$, can be easily determined. By analyticity, only terms of the form $\rho^p\tau^q$ can appear where $p+q$ is even. This leaves only the following 7 possible scalings: $\tau^2, \rho^6, \rho^4, \rho^2, 1, \rho^3\tau, \rho\tau$. By multiplying by the pre-factor, we see that the first 6 terms are precisely those that arise

from H_3, F_0, F_2, F_4, F_6, and D6/O6, respectively. The 7th term ($\rho\tau$) is new, but cancels between the two terms inside the parenthesis of eq. (A.10). Hence we obtain the form of the potential given in eq. (3.5).

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