On Brane Inflation With Volume Stabilization

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Abstract: The distance between BPS branes in string theory corresponds to a flat direction in the effective potential. Small deviations from supersymmetry may lead to a small uplifting of this flat direction and to brane inflation. However, this scenario can work only if the BPS properties of the branes and the corresponding flatness of the inflaton potential are preserved in the stable volume compactification. We present an “inflaton trench” mechanism that keeps the inflaton potential flat due to shift symmetry, which is related to near BPS symmetry in our model.
1. Introduction

It is difficult to find a mechanism for inflation in string theory. Brane inflation models [1, 2, 3, 4] provide a significant step towards this goal. However, the validity of brane inflation models in string theory depends on the ability to stabilize the compactification volume. This means that an effective 4d theory has to be found in which there is a potential for the scalar field representing the volume of the internal space. This potential has to have a minimum in which the volume field is trapped. This potential should simultaneously describe the inflaton field corresponding to the distance between branes. The effective theory has to fix the volume modulus while keeping the potential for the distance modulus flat.

Brane inflation models preserving supersymmetry should not have any potential for the modulus representing the distance between BPS branes. When supersymmetry is slightly broken, so that the configuration of branes is a near BPS state, a small attractive potential develops for the distance modulus, which in turn may lead to a slow-roll stage of inflation in the effective 4d theory. Brane inflation models of this type, where small deviations from the BPS configuration trigger inflation were studied in [2, 3, 4].

In the effective field theory, this situation may be characterized by a property of the inflaton field: its mass should be small and definitely much smaller than the Hubble parameter, \( m_{\text{infl}}^2 \ll H^2 \). In supergravity there are known mechanisms to provide this type of relation. Brane inflation models developed without an explicit moduli stabilization mechanism assumed that the stabilization mechanism would not lead to a large inflaton mass. However, the recent study of D3/D3 inflation in a warped background with fluxes and
stable volume modulus [5] has shown that the danger of $m_{\text{infl}}^2 \sim H^2$ may persist, unless a fine-tuning can be performed.

The purpose of this paper is to suggest a mechanism which protects the inflaton field from acquiring a large mass in the context of stable volume brane inflation models. The mechanism can be intuitively understood as exhibiting the fact that BPS branes should be in neutral equilibrium. Compactification with a stable volume should preserve this fact under certain conditions.

We originally found this mechanism by trying various potentials in Mathematica. We were looking for the possibility of having some kind of a flat direction: the inflaton should form a trench which gives an additional dimension to the narrow well trapping the volume modulus in the simplest KKL&T model [6] or in more general models [7]. We have dubbed the phenomena as the “inflaton trench” mechanism. Only later did we realize that the corresponding potentials allow the realization of an exact unbroken supersymmetry in AdS space. One may try to keep this property for near BPS branes in near dS space, where it will be only approximately correct, but will preserve the $m_{\text{infl}}^2 \ll H^2$ relation.

As a specific example, we will work with the D3/D7 brane inflation model which is related to the hybrid model [8] of D-term inflation [9]. We will find that it is compatible with the KKL&T mechanism of the volume stabilization. D-term inflation models in supergravity do not require a fine-tuning with regard to the inflaton mass since the source of inflationary energy is coming from the D-term while the F-term vanishes during inflation. Therefore the dangerous Kähler corrections, which could lead to $m_{\text{infl}}^2 \sim H^2$, vanish. However, the superpotential volume stabilization mechanism suggested in the KKL&T model [6] requires some non-vanishing F-terms. In this paper we will suggest a consistent combination of these two ideas which may provide the basis for stringy brane inflation models with a stable volume.

2. D3 branes, D7 branes and the volume modulus

It has been pointed out in [10] that, in brane cosmologies, it is important to identify properly the volume modulus and the modulus responsible for the distance between branes. For compactifications with only D3 and D3 branes present, there are arguments that the physical volume of the internal space is given by the combination of the real part of the chiral superfield $\rho + \bar{\rho}$ and the Kähler potential $k$ of the Calabi-Yau metric.

$$2r = \rho + \bar{\rho} - k(\phi, \bar{\phi}) .$$

One of the arguments presented in [10] is that the kinetic term for the coordinates of the D3 brane $\phi$ has the form

$$T_{D3} \int d^4x \frac{\partial_\mu \phi^i \partial^\mu \bar{\phi}^k k_{ik}}{\rho + \bar{\rho}} .$$

This in turn means that the Kähler potential of the effective 4d theory must be

$$K(\rho, \phi, \bar{\rho}, \bar{\phi})_{D3} = -3 \log(\rho + \bar{\rho} - k(\phi, \bar{\phi})) .$$
At small $|\phi|^2$, $k(\phi, \bar{\phi}) = \phi \bar{\phi} + \ldots$, and one finds that, generically without fine-tuning, in models of inflation with D3 branes the inflaton mass problem, $m^2 \sim H^2$, persists.

In the case of a compactification with only D7 branes present it was pointed out in [11] that, by dimensional reduction, the real part of the volume modulus appears in the gauge coupling constant

$$\frac{1}{g^2}(F_{\mu\nu})^2 \sim (\rho + \bar{\rho})(F_{\mu\nu})^2.$$  \hspace{1cm} (2.4)

Since the gauge coupling has to be given by the real part of the chiral superfield the dependence on $S$, the coordinate of the D7, should not enter in the definition of the volume modulus in the form $\rho + \bar{\rho} - k(S, \bar{S})$, we should have

$$2r = \rho + \bar{\rho}.$$ \hspace{1cm} (2.5)

An additional argument for identifying the physical volume in this fashion is provided by the calculation of the kinetic term for the D7 coordinates, $S$, which gives

$$T_{D7} \int d^8x \partial_{\mu} S^i \partial^{\mu} \bar{S}^k k_{ik}.$$ \hspace{1cm} (2.6)

This suggests that the Kähler potential must be given by

$$K(\rho, S, \bar{\rho}, \bar{S})_{D7} = -3 \log(\rho + \bar{\rho}) + k(S, \bar{S}).$$ \hspace{1cm} (2.7)

For small $|S|$, or in some specific models, $k(S, \bar{S}) \sim S \bar{S} + \ldots$. In the setting of [3] the type IIB string theory is compactified on $K_3 \times \mathbb{T}_2^2 \simeq \mathbb{Z}_2$. The torus is oriented along $x^4, x^5$. Therefore in these directions, for $S \sim x^4 + ix^5$ we have a canonical metric with $K = S \bar{S}$. The distance between branes should be given by the difference between the coordinates of D3 and D7. Therefore there is a puzzle: what is the correct Kähler potential when both D3 and D7 branes are present, as in the brane construction for the D-term inflation model [3]. How do we resolve the contradiction between these two expressions for $K$?

We propose to resolve this contradiction as follows. If one object is much heavier than the other object, the lighter object should be thought of as a probe which does not significantly modify the background.

The case of the D3/D7 system may be considered as a D7 brane probing the geometry of a heavy stack of D3 branes. If there are many coincident D3 branes, one may treat a D7 brane with fluxes as a probe of the D3 background. This was in fact done in [3] in sec. 2.2 where the D7 brane world-volume action in a D3 background was studied.

In a situation when the D7 is probing the geometry defined by the stack of D3 branes, the relative distance between the D7 and the stack of D3 branes is defined by the coordinates of the D7. Since the heavy D3s do not move, we may use an effective field theory with the Kähler potential of the type defined by the D7 worldvolume in eq. (2.7). We will consider the position of the D3 brane, $\phi$, not as a dynamical variable. However, $S$, being the position of the D7 will be treated as a dynamical variable. Thus there are two different physical systems: when the D7 is considered as a probe, the Kähler potential of the effective field
theory is given by (2.7). When the D3 brane is considered as a probe, the Kähler potential of the effective field theory is given by (2.3). We believe that this distinction allows us to resolve the puzzle\(^1\) of having two different choices of the Kähler potential.

The analysis of the Kahler potentials and the significant difference between the coordinates describing the position of the D7 brane and D3 brane suggested here can also be understood in the framework of the special geometry of the moduli space \(^2\) as derived for type IIB string theory compactifications in \([12]\).

### 3. BPS branes in AdS space with volume stabilization

Consider some combination of branes with fluxes which form a BPS state with some number of unbroken supersymmetries. For example, consider a D7 brane placed in the 0, 1, 2, 3, 6, 7, 8, 9 directions with a self-dual flux on its worldvolume in the compact directions, \(F_{mn} = F^*_{mn} \), \(m, n = 6, 7, 8, 9\). Add to this system many D3 branes extended in the 0, 1, 2, 3 directions at some distance from the D7 in 4, 5 directions. This system is supersymmetric. The distance between the branes, \((x^4)^2 + (x^5)^2\), can take any value without changing the fact that this is a BPS state. There is no force between branes, or, to be more precise, there is a balance of gravitational, dilaton and RR form-field forces.

The condition for unbroken supersymmetry for a D7 brane in the background geometry of a D3 brane with a harmonic function \(H(\vec{y}^2)\) was defined in \([3]\) by solving the equation

\[
(1 - \Gamma) \epsilon = 0. \tag{3.1}
\]

Here \(\Gamma\) is a generator of the local \(\kappa\)-symmetry. This gives the following equation defining the BPS state

\[
e^{-\frac{1}{2} \sigma_3 \otimes H^{1/2} (\vec{y}^2) \Gamma_{67}[(\theta_1 + \theta_2)(1 - \Gamma_{6789}) + (\theta_1 - \theta_2)(1 + \Gamma_{6789})]} \otimes \Gamma_{6789} \epsilon_0 = \epsilon_0. \tag{3.2}
\]

Here the fluxes on D7 are \(\mathcal{F}_{67} = \tan \theta_1\) and \(\mathcal{F}_{89} = \tan \theta_2\). If \(\theta_1 = \theta_2 = \theta\) we find

\[
\exp\{-\sigma_3 \otimes H^{1/2} \theta \Gamma_{67}(1 - \Gamma_{6789})\} \otimes \Gamma_{6789} \epsilon_0 = \epsilon_0, \tag{3.3}
\]

where \(\epsilon_0\) is a constant spinor. Thus for self-dual \(\mathcal{F}\) (\(\mathcal{F}^- = 0\)), an exact non-linear Killing condition can be satisfied subject to two projector equations on the Killing spinors, \(\epsilon = i\sigma_2 \otimes \Gamma_{01236789} \epsilon\) and \(\epsilon = \Gamma_{6789} \epsilon\). However, if \(\theta_1 \neq \theta_2\), equation (3.2) cannot be satisfied for non-vanishing \(\epsilon_0\). Thus any small deviation from the condition \(\mathcal{F}^- = 0\) leads to a breaking of supersymmetry.

We would like to find an effective theory language which will describe the case of unbroken supersymmetry (with \(\mathcal{F}^- = 0\)) when the position of the D7 brane in a background of the D3 is a modulus. Afterwards we will consider a deviation from the supersymmetry

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\(^1\)The existence of the puzzle was revealed in discussions with C. Burgess, J. Maldacena, S. Kachru, S. Trivedi and F. Quevedo.

\(^2\)We are grateful to S. Ferrara and C. Angelantonj for the explanation of this issue.
by looking at the case $\mathcal{F}^- \neq 0$ and try to keep the field describing the distance between branes as a modulus.

We will try to construct an F-term potential which coincides with the (AdS part of) KKLT model when the distance between the branes is exactly zero. We will call a specific field $S$ the “inflaton” in anticipation of it behaving as such once the system has been uplifted to dS space. In the AdS stage this field should just be a flat direction of the potential. This means that we need a condition for unbroken supersymmetry to be satisfied in the AdS space.

We assume that the Kähler potentials for the volume field $\rho$ and the inflaton field $S$ are decoupled. This is different from the corresponding assumption in the case of the D3 brane in [5] where the coordinate of the moving D3 brane was related to the inflaton. Thus we assume that the Kähler potential is the one in eq. (2.7) and that the superpotential, $W(\rho, S)$, is such that, at $S = 0$, it has a minimum in $\rho$, e.g. as in the KKLT model

$$W(\rho, S = 0) = W_{KKLT}(\rho) \equiv w(\rho). \quad (3.4)$$

For example, the simplest $w(\rho) = W_{KKLT}(\rho) = W_0 + A e^{-a\rho}$. Since we would like to describe a BPS state of branes, it is natural to require that a condition of an exact unbroken supersymmetry,

$$D_\rho W = 0, \quad D_S W = 0, \quad (3.5)$$

is satisfied at a particular value of $\rho = \rho_{cr}$ and at all values of the separation $s = \text{Re}S$. An exact solution of equations (3.4) and (3.5) is available under the condition that the supersymmetric extremum takes place at

$$\rho = \bar{\rho}, \quad S = \bar{S}, \quad (3.6)$$

which for $\rho = \sigma + i\alpha$ and $S = s + i\beta$ means that the state with unbroken supersymmetry requires

$$\alpha = 0, \quad \beta = 0. \quad (3.7)$$

The solution for the superpotential is given by the product of the KKLT superpotential times an $S^2$-dependent exponent.

$$W(\rho, S) = W_{KKLT}(\rho) e^{-\frac{S^2}{2}} = w(\rho) e^{-\frac{S^2}{2}}. \quad (3.8)$$

This leads to a simple potential which depends only on $\beta$ and does not depend on $s$.

$$V^F(\rho, \bar{\rho}, S, \bar{S}) = e^{-\frac{1}{4}((S - \bar{S})^2}{(\rho + \bar{\rho})^3} \left[ |D_\rho w(\rho)|^2 \frac{(\rho + \bar{\rho})^2}{3} - 3 + (S - \bar{S})^2 |w(\rho)|^2 \right]. \quad (3.9)$$

Note that the field $S$ has a canonical kinetic term due the presence of the $SS$-term in the Kähler potential. Also, the term $e^K |D_S W|^2$ is given by $e^{-\frac{1}{2}((S - \bar{S})^2}{(\rho + \bar{\rho})^3} [-(S - \bar{S})^2 |w(\rho)|^2]$. At $S = 0$ the potential coincides with the F-term potential of the KKLT model,

$$V(\rho, S = 0, \bar{\rho}, \bar{S} = 0) = V^F_{KKLT} = \frac{1}{(\rho + \bar{\rho})^3} \left[ |D_\rho w(\rho)|^2 \frac{(\rho + \bar{\rho})^2}{3} - 3 |w(\rho)|^2 \right]. \quad (3.10)$$
In fact, the perfect flatness of the potential (3.9) in the direction of the $s$ field was originally discovered by plotting the potential as a function of $(\sigma, s)$ and observing that the $s$-direction forms a trench. When the trench is cut at any constant value of $s$ it gives the profile of the KKLT superpotential for the volume $\sigma$. The plot suggests that it may be possible to bring the potential to a form that is manifestly independent of $s$, as one can see in eq. (3.9).

Another way to see the manifest $s$ independence is to note that the model can be described by a new $\tilde{K}$ and $\tilde{W}$ using Kähler transformations so that

$$\tilde{K} = -3 \ln(\rho + \bar{\rho}) - \frac{(S - \bar{S})^2}{2}, \quad \tilde{W}(\rho, S) = w(\rho). \tag{3.11}$$

The kinetic term for the $S$ field remains canonical. The conditions for unbroken supersymmetry in new variables are

$$\tilde{D}_\rho \tilde{W} = D_\rho w(\rho) = 0, \quad \tilde{D}_S \tilde{W} = -(S - \bar{S}) w(\rho) = 0. \tag{3.12}$$

These conditions are satisfied at $\alpha = \beta = 0$, i.e. at $\rho - \bar{\rho} = S - \bar{S} = 0$.

Thus we have presented here an effective model which may be interpreted as describing some BPS branes in a compact space with stabilized volume. The fact that the corresponding potential has a flat direction $s$ follows from the requirement of the existence in the theory of a supersymmetric critical point. The potential has a flat direction describing the distance between branes while keeping the volume fixed.

In the language of the $\tilde{K}$ and $\tilde{W}$, the existence of the supersymmetric flat direction is a consequence of maintaining a shift symmetry of the model under compactification:

$$S \Rightarrow S + \delta, \quad \tilde{K}(S - \bar{S}) \Rightarrow \tilde{K}(S - \bar{S}), \quad \tilde{W}(\rho) \Rightarrow \tilde{W}(\rho), \tag{3.13}$$

where $\delta \in \mathbb{R}$. Clearly, the effective potential derived from $\tilde{K}$ and $\tilde{W}$ (3.11) does not have any dependence on $S + \bar{S}$, i.e. on the inflaton field $s = (S + \bar{S})/2$. In fact, we could...
simply start with the effective $N = 1$ supergravity (3.11) and use shift symmetry to prove the flatness of the potential in the inflaton direction. This would be reminiscent of the proposal in [14], where the requirement of an analogous shift symmetry was used to rescue chaotic inflation in supergravity.

An important difference in our setting is that we derived the condition of shift symmetry of the effective $N = 1$ supergravity as a consequence of our requirement to describe the BPS branes using effective field theory [3, 13]. That is why we needed to check the consistency condition (3.12) in order to verify that the supersymmetric state corresponding to BPS branes does exist in our theory. However, once this property of the theory is established, we can use shift symmetry to prove that the potential remains flat in the inflaton direction, $s$, not only in the supersymmetric state in AdS, but for all values of the fields $s, \alpha$ and $\beta$.

It is instructive to present the potential as an explicit function of all 3 remaining fields. Using the example from KKLT where $w(\rho) = W_0 + A e^{-\rho}$ we find for $V(\sigma, \alpha, \beta)$ the following expression:

$$V(\sigma, \alpha, \beta) = e^{\beta^2 - 2a\sigma} \left[ A^2 [a\sigma(3 + a\sigma) + 3\beta^2] + 3A e^{a\sigma} (a\sigma + 2\beta^2) \cos(a\alpha) W_0 + 3e^{2a\sigma} \beta^2 W_0^2 \right].$$  (3.14)

The dependence on $\alpha$ is very simple:

$$V(\sigma, \alpha = 0, \beta) = f(\sigma, \beta) + g(\sigma, \beta) \cos(a\alpha),$$  (3.15)

where $g(\sigma, \beta)$ is negative for all values of $(\sigma, \beta)$ because $W_0$ is negative in this model. This shows that, at all values of $\sigma$ and $\beta$, the potential has a minimum at $a\alpha = 2n\pi$ and a maximum at $a\alpha = (2n + 1)\pi$. Thus $\alpha = 0$ is in fact a supersymmetric minimum. From now on we will simplify our problem by taking $\alpha = 0$.

$$V(\sigma, \alpha = 0, \beta) = e^{\beta^2 - 2a\sigma} \left[ A^2 [a\sigma(3 + a\sigma) + 3\beta^2] + 3A e^{a\sigma} (a\sigma + 2\beta^2) W_0 + 3e^{2a\sigma} \beta^2 W_0^2 \right].$$  (3.16)

At the critical point for $\sigma$ where $W_0 = -A e^{-\sigma_{cr}} (1 + \frac{2}{3} a \sigma_{cr})$ the dependence on $\beta$ simplifies and we find

$$V(\sigma_{cr}, \alpha = 0, \beta) = \frac{a^2 A^2 e^{2\beta^2 - 2a\sigma_{cr}}}{18\sigma_{cr}} (-3 + 4\beta^2).$$  (3.17)

The potential has, in addition to the supersymmetric critical point at $\beta = 0$, another non-supersymmetric minimum at $\beta^2 = 1/4$ and $D_\rho W(\rho, S) = 0, D_S W(\rho, S) \neq 0$. One finds that

$$\frac{\partial V(\sigma_{cr}, \beta)}{\partial \beta} = 0$$  (3.18)

when

$$D_\rho w(\rho_{cr}) = 0, \quad \beta(-1 + 4\beta^2) = 0.$$  (3.19)

The point $\beta = 0, \alpha = 0, \sigma = \sigma_{cr}$ is a supersymmetric saddle point, which is a maximum in the $\beta$ direction and a minimum in the $\alpha$ and $\sigma$ directions. The point $\beta = \pm 1/2$ and $\alpha = 0, \sigma = \sigma_{cr}$ is a non-supersymmetric minimum of the potential.
In our model the field $S$ is proportional to $x^4 + ix^5$, where the distance between branes is $d^2 = (x^4)^2 + (x^5)^2$. We see that the exact supersymmetric state is realized when the position of the D7 coincides with the position of the stack of D3's in $x^5$ direction. In such a case the distance in the direction of $x^4$, and therefore the total distance $d^2$, is completely arbitrary. It does not cost any energy to move the D7 in $x^4$ direction, which plays the role of the inflaton field in our model.

As in the KKLMT model, the AdS part is just a technical step towards the cosmological theory. The next step is to add a D-term potential related to a D7 brane with the non-self-dual flux on its worldvolume to lift the system to dS.

4. dS uplifting

The procedure of uplifting the AdS minimum to a dS one using an anti-D3 brane was explained in [6]. An alternative method based on non-self-dual fluxes on D7 branes was suggested in [11]. For our purpose of generalizing D-term inflation to the case of stabilized volume compactification models, it is appropriate to use this alternative method. In particular, this will allow us to study the exit from inflation towards Minkowski or dS with a very small cosmological constant. This will involve a few new steps: considering D3-D7 strings stretched between D3 and D7 branes ($\Phi_{\pm}$ fields) and their logarithmic effect on the flat direction of the classical potential. We will also be able to consider the waterfall stage of the hybrid inflation model where $S$ is close to zero and one of the charged fields, e.g. $\Phi_{+}$, will change significantly.

When we are interested only in the inflationary stage of D-term inflation with volume stabilization, the charged fields $\Phi_{\pm}$ are near the dS minimum with $\Phi_{+} = \Phi_{-} = 0$. At $\Phi_{+} = \Phi_{-} = 0$ the D-term is given by [11]

$$\frac{g^2}{2} D^2 = \frac{C}{\sigma^3}. \quad (4.1)$$

Here $g^2$ for the vector fields on the D7 is proportional to $\frac{1}{\rho + \bar{\rho}}$ and the $\rho$-dependent $D$-term is calculated via fluxes in the 6, 7, 8, 9 space of the D7 and is proportional to $\frac{1}{\rho + \bar{\rho}}$.

The total classical potential describing the inflationary stage of inflation with vanishing $\Phi_{\pm}$-fields consists of both $V^F$ and $V^D$ and is given by

$$V(\rho, S, \bar{\rho}, \bar{S}) = V^F(\sigma, \alpha, \beta) + V^D(\sigma) = \frac{e^{2\beta^2 - 2a\sigma}}{6\sigma^3} A^2 [a\sigma(3 + a\sigma) + 3\beta^2] + \frac{e^{2\beta^2 - 2a\sigma}}{2\sigma^3} \left[ Ae^{a\sigma/(2\sigma^2)} (a\sigma + 2\beta^2) \cos(a\alpha) W_0 + e^{2a\sigma} \beta^2 W_0^2 \right] + \frac{C}{\sigma^3}. \quad (4.2)$$

As before, the $s$ direction is a flat one. We will consider a simple case with $\alpha = 0$ and $\beta = 1/2$, i.e. the nonsupersymmetric minimum of $V$. The dependence of $V$ on the volume, $\sigma$, and the flat direction, $s$, is shown in Figure 2.

Thus the volume is stabilized and the inflaton field (the distance between the branes) can change without affecting the stabilization mechanism. We are led at a picture which
Figure 2: dS trench at the minimum of $\alpha = 0$ and $\beta = 1/2$ in which the volume $\sigma$ is stabilized close to the critical value of the AdS stage. There is no dependence of $V$ on the inflaton $s$.

is not exactly D-term inflation since we have a non-vanishing $V^{F}$-term in the potential. However, the main features of the D-term inflationary model and the corresponding D3/D7 brane construction are preserved: the D3-D7 strings stretched between branes (which correspond to the fields $\Phi_{\pm}$) have a mass splitting due to a non-self-dual flux. This leads to the shift symmetry breaking by quantum effects and to the logarithmic corrections to the potential $\sim \log s$, which will drive hybrid inflation.

Here again, the effective field theory tells us that if the distance between D7 from the stack of D3’s in direction $x^5$ is defined by $x^5 \sim \beta = 1/2$, it will stay there as it is a minimum of the potential. The position of D7 in the $x^4$ direction is almost flat: the classical potential is exactly flat$^{3}$ but the first loop corrections $\sim \ln s$ trigger slow-roll inflation as in D-term inflation.

After passing the critical distance, the D7 brane will form a supersymmetric bound state with the D3 and one of the charged fields $\Phi_{+}$ or $\Phi_{-}$ will acquire a vev. Preliminary investigation shows that the volume of internal space will remain stabilized during the rapid waterfall stage at the end of inflation. We plan to perform a detailed investigation of this model in future.

5. Discussion

In this note we proposed an “inflaton trench” mechanism which may lead to a viable model of inflation with volume stabilization in string theory. The shift symmetry of the classical theory protects the small mass of the inflaton in the framework of a string compactification with a stabilized volume modulus. This symmetry is shown to follow from the description of BPS branes in a supersymmetric equilibrium. A small deviation from supersymmetry leads

$^{3}$Note that until one considers the interaction of the inflaton field $s$ with the fields $\Phi_{\pm}$, the flatness of the effective potential is preserved even at the quantum level, despite supersymmetry breaking, because the masses of all other fields do not depend on $s$. We are grateful to A. Linde for the discussion of this issue.
to a small violation of the shift symmetry which allows a slow-roll regime of inflation. A similar mechanism may be responsible for the present stage of acceleration of the universe.

A picture of a “stringy landscape” that has been developed in [15] takes into account, in particular, a profile of the KKL T model. Now it may be enriched with the trenches of the kind shown in Fig. 2 where inflation or late-time acceleration of the universe may take place.

We are planning to develop these considerations in the context of the D3/D7 brane construction, with a hope that it may provide us with a realistic model of brane inflation with stable compactification. One can also expect that the “inflaton trench” mechanism may be useful for other models of brane inflation. We hope to return to the discussion of these possibilities in a separate publication.

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