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Fine-Tuning in Brane-antibrane Inflation

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I give a brief overview of brane-antibrane inflation, with emphasis on the problems of tuning to get a flat potential in the KKLMMT framework, and recent work on the nature of superpotential corrections in that model.

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Brane-antibrane inflation is one of the most important ideas for inflation from string theory. I have reviewed it previously in [1], to which the reader is directed for more complete references to the literature. Here I will recapitulate some of the historical developments that led to the KKLMMT [2] model, then discuss its tuning problems, and the challenges for finding superpotential corrections within string theory which have the right properties for producing a sufficiently flat potential.

1. Inflation from brane annhilation

The interaction energy between a parallel D3-brane and its corresponding antibrane can give rise to inflation in the early universe [3]. The subsequent brane-antibrane annihilation ends inflation and can reheat the observed universe [4], presumably located on some other brane which may or may not be coincident with the inflationary branes. This is illustrated in figure 1. One might wonder whether the branes being parallel requires an extra fine-tuning beyond those which will be discussed below. However for a D3 brane, any nonzero angle would require the brane to wrap some of the compact dimensions, similar to a helix on the surface of a drinking straw. The energy density of the wrapped brane would be greater than that of a zero-angle brane due to the greater volume required by wrapping relative to remaining straight. Thus the zero-angle configuration is energetically preferred.

In this picture, the brane-antibrane separation r plays the role of the inflaton, and the lightest mode of the stretched string between branes becomes tachyonic at a critical separation or order $1/M_s$ (the inverse string mass scale), ending inflation. In this respect, brane-antibrane inflation is quite similar to hybrid inflation.

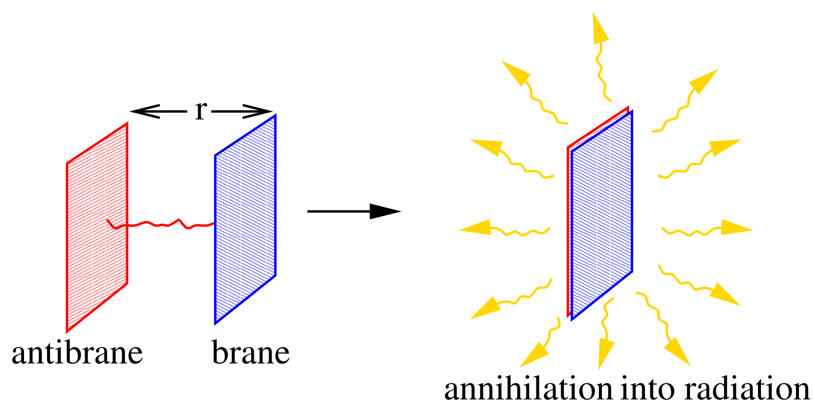

Figure 1: Brane-antibrane inflation and reheating.

1.1 Brane-antibrane action

To understand the inflationary potential, one should first note that parallel BPS (supersymmetric) D3 branes exert no force on each other. The two component forces are

$$V_{\rm grav} = -\kappa_{10}^2 \frac{\tau_3^2}{r^4}, \quad \text{gravitational attraction}$$

$$V_{\rm gauge} = +\kappa_{10}^2 \frac{\tau_3^2}{r^4}, \quad \text{RR gauge field repulsion}$$

where κ_{10}^2 is the 10D gravitational constant, τ_3 the D3-brane tension, and r the separation in the compact dimensions. Notice that due to the BPS condition, these exactly cancel each other. On the other hand, for antiparallel D3 branes, the orientation and charge of one brane is reversed, turning it into an antibrane, as illustrated in figure 2. The gravitational attraction is no longer canceled by RR-gauge repulsion, resulting in the attractive total potential

$$V_{\rm tot} = -2\kappa_{10}^2 \frac{\tau_3^2}{r^4} \tag{1.1}$$

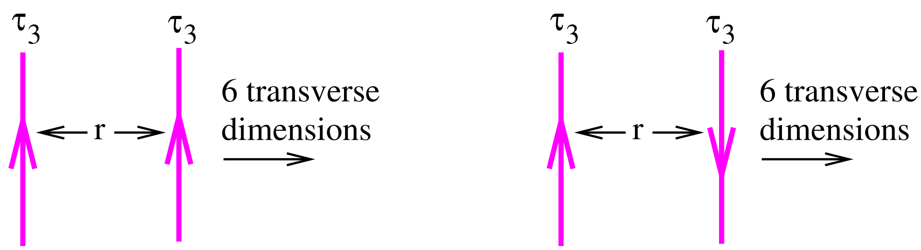

Figure 2: Brane-brane (left) versus brane-antibrane (right) configuration; orientation hence charge of antibrane is reversed.

So far we have treated the brane-antibrane separation r as if it were a single degree of freedom, but branes are not rigid objects; they fluctuate in the transverse directions, so the actual separation is not just a number, but a field $r(x^\mu)$ which depends on the position x^μ in the noncompact directions, as shown in figure 3.

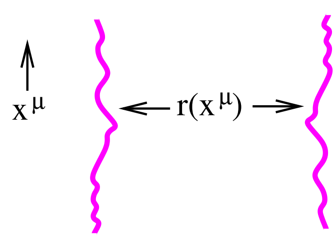

Figure 3: The inflaton field $r(x^\mu)$.

To find the kinetic term for the inflaton, we start with the Dirac-Born-Infeld (DBI) action for a single D3 or $\overline{\rm D3}$ brane,

$$S = -\tau_3 \int d^4x \sqrt{-G} \tag{1.2}$$

where $G_{\mu\nu}$ is the induced metric on the brane,

$$G_{\mu\nu} = g_{AB} \frac{\partial X^A}{\partial x^\mu} \frac{\partial X^B}{\partial x^\nu} = \eta^{\mu\nu} + \frac{\partial \phi^I}{\partial x^\mu} \frac{\partial \phi^I}{\partial x^\nu} \tag{1.3}$$

Here ϕ^I are the transverse oscillations to the brane, and expanding to leading order in them gives

$$\det G = -1 + \left(\frac{\partial\phi}{\partial x}\right)^2 + \dots \tag{1.4}$$

Hence the DBI action takes the approximate form

$$S = -\tau_3 \int d^4x \left(1 - \frac{1}{2}\left(\frac{\partial\phi}{\partial x}\right)^2 + \dots\right) \tag{1.5}$$

To find the action for the canonically normalized inflaton we let $r^I = \phi^I - \bar{\phi}^I$, where ϕ^I, $\bar{\phi}^I$ are the respective fluctuations of the brane and antibrane. The Lagrangian then splits into an uninteresting contribution for the center-of-mass, and the relevant one for the separation,

$$\mathscr{L} = -\tfrac{1}{2}\tau_3(\partial r)^2 - V(r) \tag{1.6}$$

The canonically normalized inflaton is therefore

$$\varphi = \sqrt{\tau_3}\, r = \sqrt{\tau_3}\left(\sum_I (r^I)^2\right)^{1/2} \tag{1.7}$$

and its potential is

$$V = 2\left(\tau_3 - \frac{c}{\varphi^4}\right), \qquad c = \kappa_{10}^2 \tau_3^4 \tag{1.8}$$

It will be important below that the 10D gravitional coupling is

$$\kappa_{10}^2 = M_{10}^{-8} = M_p^{-2} L^6 \tag{1.9}$$

in terms of the compactification volume L^6.

1.2 The flatness problem

To get enough inflation, we need the slow-roll parameters to be small. One finds that the η parameter provides the most stringent constraint,

$$\eta \;\equiv\; M_p^2 \frac{V''}{V} \;\sim\; -\left(\frac{L}{r}\right)^6 \tag{1.10}$$

From this formula, it appears that the only way to make η small is to demand that the brane-antibrane separation satisifes $r \gg L$. However it is impossible to separate them by more than the size of the extra dimensions, so this does not work [5]. In fact the approximation (1.10) is only valid when $r \ll L$; when $r \sim L$ compactification effects become important and the potential no longer behaves like $1/r^4$ as it does in flat space. Nevertheless, the setup is still problematic because of the assumption that the compactification volume is stabilized. Realistically L is a modulus with dynamics that can influence the inflaton. It is not obvious that the introduction of a dynamical stabilization mechanism for L will leave $V(\varphi)$ flat, even if that can be achieved for fixed L. Thus it is important to have a complete picture in which the dynamics of compactification is understood.

2. Flux Compactifications and the KKLMMT model

An important step toward more complete and realistic string-inflationary model building was the realization that background fluxes can stabilize many of the moduli of string theory. In particular, Giddings, Kachru and Polchinski (GKP) [6] showed that fluxes in warped compactifications, using a Klebanov-Strassler (KS) throat [7], generically stabilize the dilaton and complex structure moduli of type IIB string theory compactified on a 6D Calabi-Yau manifold. The situation is illustrated in figure 4. Besides the advantages of moduli stabilization, this has further appealing features: the throat generates a hierarchy through warping like in the Randall-Sundrum (RS) model [8]; a large hierarchy can be generated from natural values of the fluxes, which are quantized.

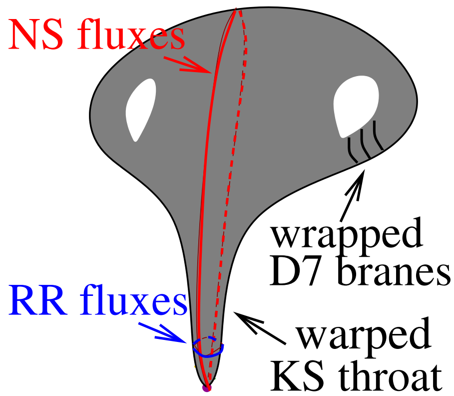

Figure 4: Klebanov-Strassler throat attached to a Calabi-Yau manifold, with fluxes of H_3 and F_3 wrapping dual 3-cycles.

2.1 Klebanov-Strassler Throat

Let us consider the KS warped throat in more detail. It can be thought of as a generalization of the RS model to 10D. The geometry of the throat is approximately $\mathrm{AdS}_5 \times T_{1,1}$, where $T_{1,1}$ is a compact space described by five angular coordinates:

$$ds^2 = a^2(r)(-dt^2 + dx^2) + a^{-2}(r)(dr^2 + r^2 ds^2_{T_{1,1}}) \tag{2.1}$$

and the warp factor takes the form

$$a(r) \cong \frac{r}{R}, \quad R = \text{AdS curvature scale} \tag{2.2}$$

The throat is a generalization of the singular conifold geometry pictured in figure 5. It is similar to a cone, but the base $T_{1,1}$ has the topology of $S_2 \times S_3$ instead of a circle. At the tip of the cone, where $r = 0$, the S_3 shrinks to zero size. One can also consider a deformed conifold in which the manifold closes off smoothly at some nonzero value $r = r_0$. These manifolds, which are complex, can be described in terms of four complex coordinates w_i restricted by one complex condition,

$$\sum_{i=1}^{4} w_i^2 = z \tag{2.3}$$

The case $z = 0$ corresponds to the singular conifold, while $z \neq 0$ describes the deformed conifold. z is a dynamical field, the complex structure modulus, which is a flat direction in the absence of fluxes, but which acquires a potential when fluxes are turned on for $H_{(3)}$, the Kalb-Ramond field strength, and for $F_{(3)}$, the the field strength of the Ramond-Ramond (RR) 2-form $C_{(2)}$. The flux quanta are specified by integers M and K,

$$\left(\frac{M_s}{2\pi}\right)^2 \int_A F_3 = M, \qquad \left(\frac{M_s}{2\pi}\right)^2 \int_B H_3 = -K \tag{2.4}$$

where A and B denote dual 3-cycles of the Calabi-Yau, portrayed as circles in figure 4. The stress-energy of the fluxes fixes the value of z to be

$$z = e^{-2\pi K/g_s M} = a_0^3 \tag{2.5}$$

In language familiar from the RS model, a_0 is the warp factor at bottom of throat, which plays the role of the infrared brane.

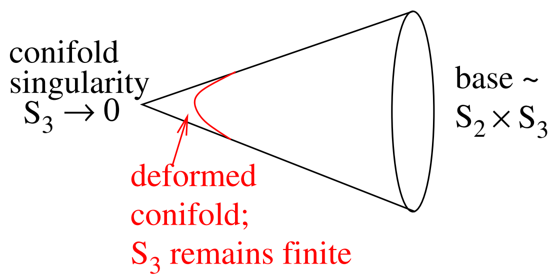

Figure 5: The singular and deformed conifold geometries.

2.2 Getting Inflation: KKLMMT

We have now introduced (almost!) all of the basic ingredients required for building a semi-rigorous inflationary model from string theory. KKLMMT [2] added a D3 and $\overline{\rm D3}$ into the throat, as shown in figure 6. In this configuration, the $\overline{\rm D3}$ sinks quickly to bottom of the throat, while the D3 is almost neutrally buoyant. This comes about because of the background fluxes, which induce a RR 5-form field strength background through its equation of motion,

$$dF_{(5)} \sim H_{(3)} \wedge F_{(3)} \tag{2.6}$$

The corresponding gauge potential is the 4-form, whose solution is

$$C_{(4)} = a^4(r) \tag{2.7}$$

The 4-form couples to D3 and $\overline{\rm D3}$ through the Chern-Simons (CS) action, *i.e.*, the second term in

$$S = -\tau_3 \int d^4x \left(a^4(r)\sqrt{1 + a^{-4}(r)(\partial\phi^I)^2} \mp C_{(4)} \right) \tag{2.8}$$

$$\cong \tfrac{1}{2}\tau_3(\partial\phi^I)^2 + \begin{cases} 0, & \rm D3 \\ -2\tau_3\, a^4(r) \int d^4x, & \overline{\rm D3} \end{cases} \tag{2.9}$$

The first term in eq. (2.8) is the DBI action including the warp factor in the background geometry. Eq. (2.9) is leading term in the slowly-rolling limit. The constant parts of the DBI and CS terms cancel for D3 but add for $\overline{\rm D3}$, explaining why one floats while the other sinks.

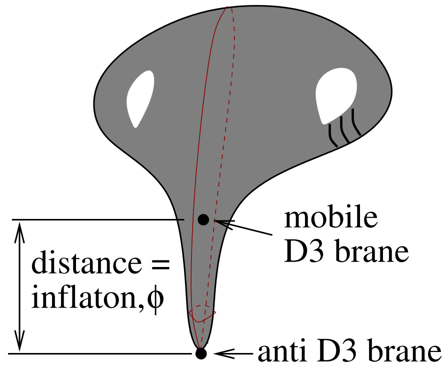

Figure 6: D3 and $\overline{\text{D3}}$ in KS throat in KKLMMT setup

However, we have ignored the D3-$\overline{\text{D3}}$ interaction in the approximation (2.9). To derive it, one can consider the action for a static $\overline{\text{D3}}$ at position $r = r_0$ in the throat:

$$S = -\tau_3 \int d^4x \sqrt{g_4(r_0)} - \tau_3 \int d^4x\, C_{(4)}(r_0) \tag{2.10}$$

If there is no additional brane in the throat, the background fields have solution $\sqrt{g_4} = C_{(4)} = a_0^4$ and the potential for the $\overline{\text{D3}}$ is $V = -2a_0^4\tau_3$ as in (2.9). Now imagine adding a D3 at position r; it perturbs the geometry

$$g^{(6)}_{\mu\nu} \to g^{(6)}_{\mu\nu} + \delta g^{(6)}_{\mu\nu} \tag{2.11}$$

The perturbation satisfies the Poisson equation in the 6 extra dimensions,

$$\nabla^2 \delta g^{(6)}_{\mu\nu} = C\,\eta_{\mu\nu}\,\delta^{(6)}(\vec{r}) \quad \Rightarrow \quad \delta g^{(6)}_{\mu\nu} \sim C\,\eta_{\mu\nu}\,(r-r_0)^{-4} \tag{2.12}$$

Substituting the perturbed background $g_4 \sim 1/g_6 \sim C^2_{(4)}$ back into the action (2.8), one obtains the potential

$$V = \frac{2a_0^4\tau_3}{1 + a_0^4(r-r_0)^{-4}} \tag{2.13}$$

If eq. (2.13) was the final result, it would be an ideal potential for getting slow-roll inflation, because of the new parameter $\varepsilon \equiv a_0^4$ which can be made small without any fine tuning by appropriate choices of the fluxes in (2.5). Notice the potential can be approximated as

$$V \cong 2\varepsilon\tau\left(1 - \frac{\varepsilon}{r^4}\right) \tag{2.14}$$

By simply taking $\varepsilon \ll 1$, one can make V as flat as desired. The η slow-roll parameter is

$$\eta = \frac{V''}{V} \cong -20\varepsilon \tag{2.15}$$

which can easily be made small enough to get 60 e-foldings of inflation and a nearly scale-invariant spectral index.

2.3 η strikes back

Unfortunately, the nice potential (2.14) is not the final answer, because we have ignored the dynamics of the overall volume (Kähler) modulus T. This is the one modulus which is not stabilized by the fluxes. We will now show that the interaction of T with the inflaton φ induces a large mass for φ, which can be expressed as an additional term in the inflaton potential of the form

$$\delta V = \tfrac{1}{2}m^2\varphi^2, \qquad m^2 \sim V_0 \sim H^2 \tag{2.16}$$

Since $m \sim H$, inflation is spoiled:

$$\eta \;=\; \frac{V''}{V} \;\rightarrow\; \frac{2}{3}$$

The inflaton never rolls slowly!

To understand how the problem arises, we must consider how the 10D metric depends on T,

$$ds^2 = e^{-6u}\,a^4\,dx^2 + e^{2u}\,a^{-4}\,\tilde{g}^{(6)}_{ab}\,dy^a\,dy^b \tag{2.17}$$

where u and T are related to the compactification length L through

$$e^{4u} = T + \overline{T} = L^4 \tag{2.18}$$

When (2.17) is used to compute the induced metric that goes into the DBI action, the kinetic term of the inflaton gets modified to

$$(\partial\varphi)^2 \rightarrow \frac{(\partial\varphi)^2}{T+\overline{T}} \tag{2.19}$$

On the other hand, the low-energy effective action for the brane position can also be written in the language of supergravity (SUGRA). Consistency between the DBI and SUGRA approaches implies that the Kähler potential for T gets modified in the presence of the D3 brane to

$$K = -3\ln(T+\overline{T}-|\varphi|^2) \equiv -2\ln 2\sigma \tag{2.20}$$

In SUGRA, the F-term potential then also gets modified, since $V \sim e^K$. This implies that (2.14) is corrected to [9]

$$V \rightarrow \frac{V}{(2\sigma)^2} = \frac{V}{(T+\overline{T}-|\varphi|^2)^2} \tag{2.21}$$

For small values of φ we can expand the new contribution to obtain the Lagrangian

$$\mathscr{L} \sim -\frac{(\partial\varphi)^2}{T+\overline{T}} - \frac{V}{(T+\overline{T})^2}\left(1+\frac{2|\varphi|^2}{T+\overline{T}}\right) \tag{2.22}$$

Because of the new factor $T+\overline{T}$, we must rescale φ to get a canonically normalized kinetic term. Doing so gives the inflaton mass (in units $M_p = 1$)

$$m_\varphi^2 \sim \frac{V}{2\sigma} \sim H^2 \tag{2.23}$$

Thus the warp factor no longer helps to make η small.

3. Tuning with Superpotential Corrections?

The solution which was advocated in ref. [2] to overcome the η problem was to cancel the unwanted positive contribution to m_φ^2 by appropriately modifying the superpotential W. In order to stabilize the Kähler modulus, it was assumed that a nonperturbative contribution Ae^{-aT} was present [10],

$$W = W_0 + Ae^{-aT} \tag{3.1}$$

which generates a potential for T with a nontrivial minimum. Generically one expects this superpotential to also have some φ dependence, which was parametrized in ref. [2] as a correction of the form

$$W \to W_0 + Ae^{-aT}(1+\delta\varphi^2) \tag{3.2}$$

By tuning δ at the level of 1 part in 100, the inflaton mass can be made sufficiently small for inflation.

In an interesting new development, ref. [11] noted that it is not necessary to merely parametrize these corrections; rather, they can be explicitly computed from string theory. One can thus check whether the desired tuning can actually be realized. To make the computation tractable, it is necessary to ignore the Calabi-Yau in the unwarped region and assume that the geometry is well-approximated by the KS throat by itself. The superpotential corrections arise due to the stack of D7 branes wrapping a 4-cycle of the throat, which were a necessary ingredient of the GKP construction. This is illustrated in figure 7.

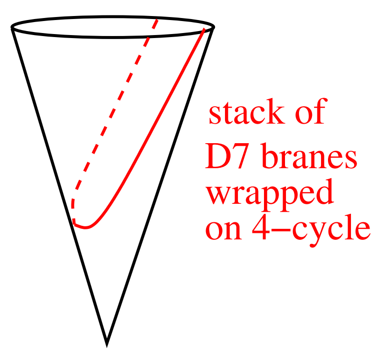

Figure 7: D7 branes wrapped on a 4-cycle of the KS throat

The superpotential corrections are determined by the 4-cycle on which the D7 branes wrap the throat, for which there are infinitely many choices. A simple class of 4-cycles which preserve SUSY is given by [12]

$$\prod_{i=1}^{4} w_i^{p_i} = \mu^P \tag{3.3}$$

where p_i are integers, $P = \sum p_i$ and the parameter μ determines how close to the bottom of the throat the 4-cycle extends. (Notice that the constraint (3.3), together with the original conifold restriction (2.3), indeed reduces the 8D complex manifold $\{w_i\}$ to a 4D subspace.) Within this class, ref. [11] shows that the superpotential corrections take the form

$$W = W_0 + Ae^{-aT}\left(1 - \frac{\prod_i w_i^{p_i}}{\mu^P}\right)^{1/N_{D7}} \tag{3.4}$$

where N_{D7} is the number of D7 branes in the stack.

The string-derived correction to W was used in ref. [13] to find the corresponding correction to the F-term potential:

$$\begin{aligned} V_F &= \frac{\kappa_4^2}{12\sigma^2}\Big[\,(T+\bar{T})|W_{,T}|^2 - 3(\overline{W}W_{,T} + \text{c.c.}) \qquad (3.5)\\ &\qquad + \frac{3}{2}\left(\overline{W}_{,\bar{T}}w^j W_{,j} + \text{c.c.}\right) + \frac{1}{c}k^{\bar{i}j}\overline{W}_{,\bar{i}}W_{,j}\Big] \\ &= \frac{\kappa_4^2}{12\sigma^2}\Big[\left[(T+\bar{T})a^2 + 6a\right]|A|^2 e^{-2a(T+\bar{T})} + 3aW_0(Ae^{-aT} + \bar{A}e^{-a\bar{T}}) \qquad (3.6)\\ &\qquad -\frac{3}{2}ae^{-a(T+\bar{T})}\left(\bar{A}w^j A_{,j} + \text{c.c.}\right) + \frac{1}{c}k^{\bar{i}j}\overline{A}_{,\bar{i}}A_{,j}e^{-a(T+\bar{T})}\Big], \end{aligned}$$

The new terms are those in the last line of (3.7). This can be explicitly evaluated in terms of the angular coordinates on the $T_{1,1}$ manifold, using

$$\begin{aligned} w_1 &= r^{3/2}\, e^{\frac{i}{2}(\psi-\phi_1-\phi_2)} \sin\frac{\theta_1}{2}\sin\frac{\theta_2}{2}, \\ w_2 &= r^{3/2}\, e^{\frac{i}{2}(\psi+\phi_1+\phi_2)} \cos\frac{\theta_1}{2}\cos\frac{\theta_2}{2}, \\ w_3 &= r^{3/2}\, e^{\frac{i}{2}(\psi+\phi_1-\phi_2)} \cos\frac{\theta_1}{2}\sin\frac{\theta_2}{2}, \\ w_4 &= r^{3/2}\, e^{\frac{i}{2}(\psi-\phi_1+\phi_2)} \sin\frac{\theta_1}{2}\cos\frac{\theta_2}{2}, \end{aligned} \tag{3.7}$$

We find that the new contribution δV_F to V_F due to the superpotential corrections cannot help with tuning the inflaton potential, because it gets minimized at the value $\delta V_F = 0$ when $\theta_1 = \theta_2 = 0$. For small θ_i, δV_F takes the form

$$\delta V_F = M_{11}(\theta_1^2 + \theta_2^2) + M_{12}\cos\left(\tfrac{1}{2}\tilde{\psi}\right)\theta_1\theta_2 + \dots \tag{3.8}$$

where $\tilde{\psi} = \psi - \phi_1 - \phi_2$ and $M_{11}^2 > \frac{1}{4}M_{12}^2\cos^2(\frac{1}{2}\tilde{\psi})$ for physically reasonable values of the parameters. The energetically preferred brane position is thus at $\theta_i = 0$, for which δV_F has no effect.

However, there is another correction which, when combined with δV_F, leads to a nonvanishing correction to the potential. Ouyang [12] showed that the D7-branes cause the dilaton to acquire a dependence on position in the compact dimensions:

$$e^{-\Phi} = \frac{1}{g_s} - \frac{N_{D7}}{2\pi}\log\left(\frac{r^{3/2}}{\mu}\sin\frac{\theta_1}{2}\sin\frac{\theta_2}{2}\right). \tag{3.9}$$

SUSY is not broken by this effect, and so by itself it does not contribute to the D3 brane potential. However, if one also introduces nonprimitive G_3 fluxes, which break SUSY spontaneously,

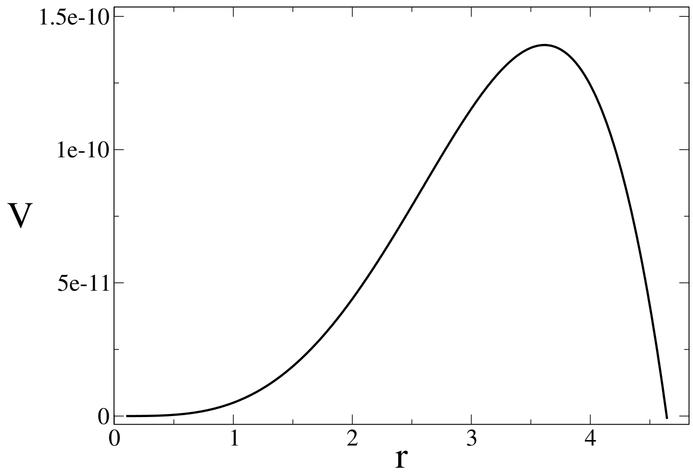

Figure 8: The correction to the inflaton potential due to superpotential and dilaton corrections, as a function of r.

(3.9) gets modified in such a way that the spatially-varying dilaton background leads to an extra contribution to the D3 potential [13, 14]

$$\delta V_O = -\frac{\delta N(\varepsilon)}{2\pi}\frac{T_3\xi_0^4}{R^2}\left(\frac{r}{r_0}\right)^4 \log\left(\frac{r^{3/2}}{\mu}\sin\frac{\theta_1}{2}\sin\frac{\theta_2}{2}\right) + \mathscr{O}(\varepsilon^2) \tag{3.10}$$

Combining this with δV_F, one sees that $\theta_i = 0$ is no longer a minimum of the full potential since δV_O diverges as $\theta_i \to 0$. There is a competition between the two terms which leads to nontrivial values of θ_i, at which the full $\delta V_{\rm tot}$ no longer vanishes.

We can then ask the question: is it possible to tune $\delta V_{\rm tot}$ against the $m^2\varphi^2$ of KKLMMT to get a flat potential for inflation? We find that for the class of embeddings (3.3), the answer is no: the curvature of $\delta V_{\rm tot}$ has the wrong sign, and only exacerbates the η problem coming from the $m^2\varphi^2$ term. Evaluated at the energetically preferred angles and Kähler modulus, $\delta V_{\rm tot}$ as a function of r has the form shown in figure 8. It has a maximum at a value of $r = r_{\rm max}$ which is close to the radius of closest approach of the D7 brane to the bottom of the throat. Near this maximum, the curvature of the potential is negative, but it is much too large to support inflation. The contribution to the η parameter at this point is

$$\eta \sim \frac{(T+\overline{T})M_p^2}{\tau_3\, r_{\rm max}^2} \sim (T+\overline{T})g_s(2\pi)^3\left(\frac{M_p}{M_s}\right)^4 \gg 1 \tag{3.11}$$

Although we do not obtain inflation from this construction, it is interesting to note that it does give us *uplifting*; that is, δVtot gives a positive contribution to V, which is necessary for offsetting the negative value of V_F at its minimum, which comes from the superpotential (3.1). In ref. [10] this problem was overcome by the addition of the $\overline{\rm D3}$, which explicitly breaks supersymmetry, and is thus at odds with the SUGRA formalism used to compute the rest of the potential. An advantage of our uplifting contribution is that it does not explicitly break SUSY, and can thus be derived from a superpotential.

Qualitatively, the uplifting works rather similarly to that from $\overline{\text{D3}}$ branes, as can be seen by comparing the σ-dependence of the two potentials:

$$\delta V_{\overline{D3}} = \frac{c}{\sigma^2}, \qquad \delta V_{\text{tot}} = \frac{c}{\sigma^2}\ln(f(\sigma)) \tag{3.12}$$

The shape of the uplifted potential for σ is illustrated in figure 9.

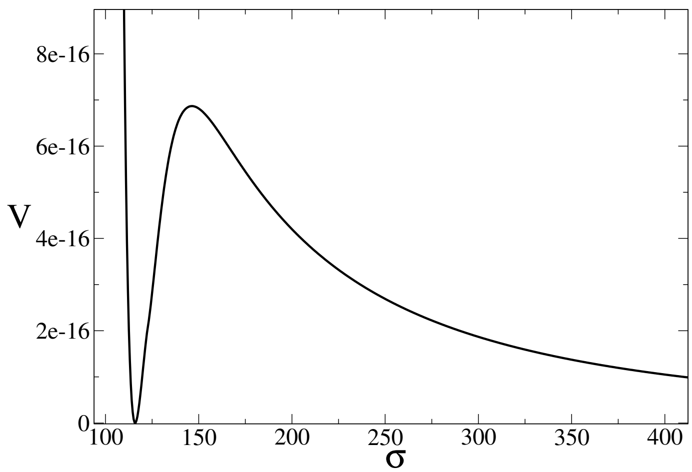

Figure 9: Potential for Kähler modulus, uplifted to Minkowski vacuum by superpotential and dilaton corrections.

3.1 Inflation using symmetric throats

Lest we give the impression that no working models exist, it is worthwhile to note an exception [15], which builds a flat region into the potential by assuming there are two nearby throats on the Calabi-Yau. Clearly a brane at the midpoint between them will be at an unstable maximum—see figure 10. The potential has the form

$$V(r) = V_0 - \frac{\tau_3^2 a^8}{\pi^3 M_{10}^8}\left(|\vec{r}-\vec{r}_1|^{-4} + |\vec{r}+\vec{r}_1|^{-4}\right) \tag{3.13}$$

if the two $\overline{\text{D3}}$'s at the bottoms of the throats are located at $\pm\vec{r}_1$, respectively. Ref. [15] shows that the negative curvature of this potential can be tuned against the unwanted positive contribution from the Kähler modulus to get $\eta \ll 1$ if

$$r_1 \sim a_0^{2/3} L \tag{3.14}$$

which can be naturally achieved. This therefore looks like a good candidate theory for brane-antibrane inflation.

4. Conclusions

Brane-antibrane inflation, which at a qualitative level seems like an intuitively appealing new way of getting inflation from string theory, is much harder to successfully implement than one

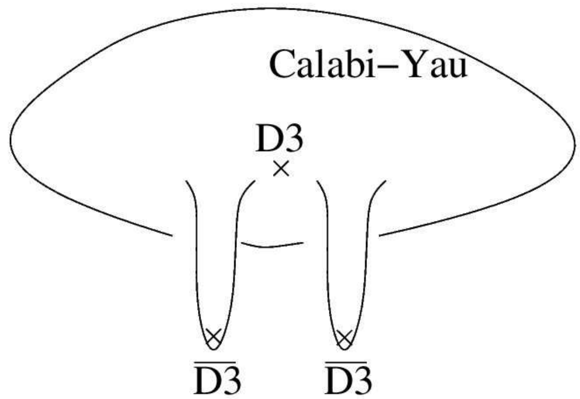

Figure 10: Mobile D3 brane between two throats.

might have guessed. Even if one is willing to fine-tune the potential, it is not obvious that string theory provides the latitude to do so, although the special case of symmetrically-placed throats seems to provide a working example. It may also be possible to achieve the desired tuning by more intricate choices of D7-brane embeddings in the single-throat scenario [16].

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