Brane inflation and cosmic string tension in superstring theory

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Abstract. In a simple reanalysis of the KKLMMT scenario, we argue that the slow roll condition in the $D3$–$\bar{D}3$-brane inflationary scenario in superstring theory requires no more than a moderate tuning. The cosmic string tension is very sensitive to the conformal coupling: with less fine-tuning, the cosmic string tension (as well as the ratio of tensor to scalar perturbation mode) increases rapidly and can easily saturate the present observational bound. In a multi-throat brane inflationary scenario, this feature substantially improves the chance of detecting and measuring the properties of the cosmic strings as a window to the superstring theory and our pre-inflationary universe.

Keywords: string theory and cosmology, inflation

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1. Introduction

Cosmic strings were originally proposed as an alternative to inflation, providing the seeds for structure formation in our universe [1]. The cosmic microwave background data from WMAP and earlier experiments [2, 3] strongly support the inflationary universe scenario [4] as being the origin of the early universe, and rules out the cosmic string scenario for structure formation. However, the origin of the inflaton and its corresponding inflaton potential are not well understood.

Branes play a key role in superstring theory [5], and the brane world is a natural realization of our universe in string theory. In this scenario, standard model particles are open string modes confined to the branes (which have three large spatial dimensions) while the gravitons are closed string modes in the bulk. Brane inflation [6], where non-BPS branes move towards each other in the brane world, is a generic inflationary scenario in superstring theory. In this scenario, inflation ends when the branes collide and heat the universe, initiating the hot big bang. Cosmic strings (but not domain walls or monopoles) are copiously produced during the brane collision [7, 8]. Properties of superstring theory cosmic strings are well studied [9]–[15]. Their contribution to density perturbations is small compared to that coming from the inflaton. Their properties are compatible with all observations today, but are likely to be tested in the near future. The properties of the cosmic strings depend on the type of brane involved in the scenario. So detecting and measuring the properties of the cosmic strings provides a window to the superstring theory and our pre-inflationary universe.

The simplest brane inflationary scenario is realized by considering the motion of a D3-brane moving towards an $\overline{D}3$-brane in a compact manifold [16]. This scenario can be realized by introducing an extra $D3-\overline{D}3$-brane pair in a KKLT-like vacuum [17], where all moduli are dynamically stabilized. The $\overline{D}3$-brane is sitting at the bottom of a throat (the A throat) as the mobile D3-brane is moving towards it. This is the KKLMMT scenario [11]. For an inflationary scenario to work, the production of defects other than cosmic strings must be suppressed by many orders of magnitude. The $D3-\overline{D}3$ scenario achieves this property automatically. A simple generalization of the above model, namely,
the multi-throat scenario [18]–[20], can be easily achieved by putting the standard model branes in another throat (the $S$ throat), where the gauge hierarchy problem may be solved with enough warping [21], while the inflationary constraints require a more modest warping for the $A$ throat. Recent arguments [22] suggest that heating up the standard model branes to start the hot big bang may not be an issue. We believe that this multi-throat scenario is quite generic and realistic. In addition, it also allows one to ensure the (initial) stability of cosmic strings [12, 13] so that they can evolve into a scaling network [23] that may be detected. Here we shall consider this scenario in some detail.

Consider the following generic potential in the $A$ throat in the multi-throat scenario:

$$
V = V_K + V_A + V_{DD} = \frac{1}{2} \beta H^2 \phi^2 + 2T_3h_A^4 \left(1 - \frac{1}{N_A} \frac{\phi^4}{\phi^4} \right) + \cdots
$$

where the first term $V_K$ receives contributions from the Kähler potential and various interactions in the superpotential [11] as well as possible $D$-terms [24]. $H$ is the Hubble constant and the inflaton $\phi$ is the position of the $D3$-brane, so this interaction term behaves like a conformal coupling. In general, $V_K$ depends on where the $D3$-brane is sitting so $\beta$ is actually a function of $\phi$. However, we expect $\beta$ to stay more or less constant in each throat, so we use that approximation here. $\beta$ receives many contributions and is non-trivial to determine [25]–[27]. Generically $\beta \sim 1$; for slow roll, $|\beta|$ must be small; this is a fine-tuning. The second term $V_A$ is the effective cosmological constant coming from the presence of the $\overline{D3}$-brane (with tension $T_3$) sitting at the bottom of the $A$ throat. $h_A \ll 1$ is the corresponding warp factor. This term drives inflation. The last term $V_{DD}$ is the Coulomb-like attractive potential between the $D3$-brane and the $\overline{D3}$-brane.

Generically we expect $\beta \sim 1$ and the slow roll condition for enough e-folds of inflation is not satisfied. Unless the $D3$-brane sits in a trench with a shift symmetry [28, 25], the model requires a fine-tuning on $\beta$. In the KKLMMT scenario, $|\beta| \lesssim 1/100$ so that $V_{DD}$ dominates over $V_K$. This is the two-orders-of-magnitude fine-tuning alluded to in [11]. Furthermore, in [18] it was demonstrated that a fine-tuning of 1/1000 was required on the parameters in the model to achieve enough inflation, although the inflationary potential did not contain the $V_K$ term. Here we would like to ask exactly how small $\beta$ must be to satisfy the slow roll condition and other constraints in inflation. We find that $\beta \sim 1/6$ is possible, in which case $V_K$ (actually all three terms in $V$) plays an important role in the inflationary scenario. This means that only a moderate tuning of $\beta$ is needed. Note that $\beta$ depends on where the throat is sitting in the bulk. If there are many throats, it is likely that one of the throats will have a small enough $\beta$, so a throat with $\beta \lesssim 1/6$ may not require any tuning. In this case, we may start with a number of extra pairs of $D3$–$\overline{D3}$-branes. The $D3$-brane around the throat with the smallest $\beta$ is likely to move the slowest and ends up as the brane responsible for inflation.

One may understand the $\beta$ relaxation simply by looking at the slow roll parameter $\eta = \beta/3 - C/\phi^6$, where $C$ is a constant obtainable from equation (1). Suppose that $\eta$ is positive at the beginning of the inflationary epoch. As $\phi$ decreases, the cancellation between the two terms guarantees that $\eta$ will be very small for some range of $\phi$. For relatively large $\beta$, $\eta$ would be positive at $N_e \simeq 55$ e-folds before the end of inflation, so the density perturbation power spectrum index $n_s \simeq 1 + 2\eta > 1$, that is, the scale-invariant spectrum is blue-tilted. In comparison, the spectrum is red-tilted when $\beta = 0$. 

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Figure 1. The cosmic string tension $G\mu$ and the density perturbation power spectrum index $n_s$ are shown as functions of $\beta$ in the KKLMMT-like brane inflationary scenario. Here $g_s = 1$. The horizontal line at $-6.2$ in the $\log(G\mu)$ graph indicates the present observational bound on the cosmic string tension, which corresponds to $\beta \lesssim 1/7$. The observational constraint on $n_s \leq 1.086$ also gives $\beta \lesssim 1/7$.

What may be most interesting in the above $\beta$ relaxation is that the cosmic string tension $\mu$ grows rapidly with increasing $\beta$. For $(p, q)$ strings [12], we find (with $G$ as the Newton’s constant and $g_s$ the string coupling)

$$G\mu_{p,q} = 4 \times 10^{-10} \sqrt{p^2 g_s + q^2 / g_s} f(\beta)$$

where $f(\beta)$ is a function of $\beta$ and $N_e$, where $N_e \simeq 55$ is the number of e-folds before the end of inflation. For $\beta = 0$, $f(\beta = 0) = 1$, so $G\mu$ reduces to that obtained for the KKLMMT scenario [11, 12], where either the F-string ($(p, q) = (1, 0)$) or the D-string ($(p, q) = (0, 1)$) has $G\mu \geq 4 \times 10^{-10}$. For large $\beta N_e$, $f(\beta)$ grows like $(\beta N_e)^{5/4} \exp(\beta N_e/2)$, so $G\mu$ grows rapidly even with moderate values of $\beta$. The cosmic D-string tension and the power spectrum index $n_s$ as functions of $\beta$ are shown in figure 1.

As $\beta$ becomes negative, the $\beta$ term will tend to move the D3-brane out of the throat. To make sure that the attractive $D3\overline{D3}$ force is strong enough to pull the D3-brane towards the $\overline{D3}$-brane at the bottom of the throat (so inflation can end), $\beta$ cannot be too negative. It turns out that this condition requires any negative $\beta$ to be exponentially close to zero.

The WMAP analysis of inflationary parameters [29] gives $n_s \leq 1.28$, which means $\beta \leq 0.4$, which is too weak to be useful. The observational bound on the ratio $r$ of the tensor to the scalar mode perturbation, $r \leq 0.89$, translates to $\beta < 0.22$, which is not stringent at all. Here $r$ as a function of $\beta$ is shown in figure 2. Note that $r$ increases exponentially as $\beta$ increases. The observational bound on the running of $n_s$, namely, $dn_s/d\ln k$, however, is not very conclusive. $dn_s/d\ln k$ as a function of $\beta$ is also shown in figure 2. If one uses only the WMAP data, the bound is $dn_s/d\ln k \leq 0.03$, which gives no bound on $\beta$. If one uses the WMAP + 2dFGRS data, the bound $dn_s/d\ln k \leq 0.01$ gives
the bound $\beta \lesssim 1/5$. If we take seriously\footnote{In contrast to the WMAP data, the cosmological application of the Lyman $\alpha$ data involves some non-trivial astrophysics and may incur unknown uncertainties. It is important to understand better the Lyman $\alpha$ clouds because of its implication to inflation data analysis.} the bound from WMAP + 2dFGRS + Lyman $\alpha$ on $dn_s/d\ln k$, then we find $\beta \lesssim 1/22$. For $\beta \simeq 1/22$, $G\mu \simeq 10^{-8}/\sqrt{g_s}$. However, in the more recent analysis of WMAP data combined with SDSS Lyman $\alpha$ and other data\footnote{The bound on the gravitational wave background from pulsar timing also puts a bound on $G\mu$. However, they are not as stringent [33]. Reference [34] apparently gives a stronger bound than that quoted here. However, that bound assumes a gravitational wave background at one frequency, which is suitable for massive binary black holes but not for cosmic strings, which have a very wide frequency band. We thank E Flanagan for pointing out this issue.}, the inflationary parameters are determined. Depending on how they fit the data, they find $n_s$ goes from $0.98 \pm 0.02$ to $1.02 \pm 0.033$. If the bound on $n_s$ is taken to be $2\sigma$ above the larger central value, we find $n_s \leq 1.086$, or around $\beta \leq 1/7$. In this paper, we shall use this bound on $\beta$. It is clear that more data and additional analysis will be important. If the data/analysis holds up, extensions of this scenario should also be explored thoroughly. Generically, any model extension will in general relax the bounds on the parameters.

We see that, in this $D3-\overline{D3}$-brane inflationary scenario, the cosmic string tension has its value in the range

$$4 \times 10^{-10} \lesssim G\mu \lesssim 6 \times 10^{-7}$$

where a fine-tuning is needed to reach the lower values and the upper value is the present observational bound [10,31] coming from WMAP data\footnote{This observational upper bound gives $\beta \lesssim 1/7$.}. This observational upper bound gives $\beta \lesssim 1/7$.

It is interesting that the constraint on $\beta$ ($\lesssim 1/6$) from model building (enough e-folds of inflation etc) is quite close to that from the data. For practical purposes, we shall use

$$0 \lesssim \beta \lesssim 1/7.$$  

We note that slight variations of the above scenario can relax the bounds further. We summarize the various bounds in table 1.
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Table 1. Various bounds on $\beta$ coming from observational data [29, 10, 31, 30]. The more recent analysis [30] using all available data gives $n_s \leq 1.086$ which implies $\beta \lesssim 1/7$.

| Model building   | $\beta \lesssim 1/6$ |
|------------------|----------------------|
| $n_s$            | $\beta \lesssim 1/7$ |
| $r$              | $\beta \lesssim 1/5$ |
| $G\mu$           | $\beta \lesssim 1/7$ |

With larger $G\mu$, more ways to detect cosmic strings are opened up. The signature of the cosmic strings coming from the superstring theory typically has a non-trivial spectrum in the tension [9, 12]. Larger $G\mu$ drastically improves the chances of detecting and measuring the tension spectrum and so the testing of superstring theory.

Larger $G\mu$ also corresponds to a blue-tilt of the scale-invariant power spectrum of the density perturbation. As $\beta$ increases (say $\beta \geq 0.02$) the power spectrum index $n_s$ becomes blue-tilted. For small $G\mu$, $\beta$ is negligibly small so the power spectrum is red-tilted, as given by the KKLMMT scenario. We note that the range $G\mu$ in this $D3\bar{D}3$-brane inflationary scenario is comparable to that given in [8] for a variety of brane inflationary scenarios. There, the larger $G\mu$ values correspond to the branes-at-a-small-angle scenario [32]. In that scenario, the power spectrum is red-tilted, so a measurement of $n_s$ can distinguish between these two scenarios. Measuring the cosmic string tension spectrum will also be important for distinguishing them. In general, measuring $n_s$, $r$, $dn_s/d\ln k$ together with the $G\mu$ tension spectrum and the corresponding number densities will go a long way to pinpointing the particular brane inflationary scenario. We note that the predictions of the above inflationary quantities and the cosmic string tension are insensitive to the warp factor $h_A$ and the background charge $N_A$. For example $10^{-3} \lesssim h_A \lesssim 10^{-1}$ and $N_A \lesssim 10^4$ are compatible with the inflationary scenario and the above predictions. Note that $N_A \gg 1$ is assumed for the validity of the supergravity approximation used. For our purpose, $N_A \geq 100$.

Fast roll in the multi-throat scenario was proposed as an interesting alternative to slow roll [20, 35]. However, this fast roll scenario requires a fine-tuning (e.g., the background charge $N \sim 10^{14}$), so the above slow roll scenario seems most natural. For larger values of $\beta$ at the $A$ throat, we also consider the possibility that inflation takes place while the $D3$-brane is moving slowly out of another throat. In this case, $\beta$ for this third throat must be moderately small and negative.

2. The set-up

The realistic set-up is a type IIB orientifold (or F theory) compactified on a Calabi–Yau threefold with fluxes [36, 17], where all moduli are stabilized. Inside the bulk of the Calabi–Yau manifold, there are local regions, or throats, with warped geometry [37]. The metric in any throat has the approximate $\text{AdS}_5 \times X_5$ form, where $X_5$ is some orbifold of $S_5$ and the $\text{AdS}_5$ metric in Poincaré coordinates takes the form

$$ds^2 = h(r)^2 (-dt^2 + a(t)^2 d\vec{x}^2) + h(r)^{-2} dr^2$$

(5)
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with the warp factor

\[ h(r) = \frac{r}{R} = e^{-2\pi K/3M_g}, \tag{6} \]

where \( K \) and \( M \) are the background NS–NS and RR fluxes respectively, and \( R \) represents the curvature radius of the AdS throat and is given by \( [37] \)

\[ R^4 = \frac{27}{4\pi g_s N^2}. \tag{7} \]

Here \( N = K M \) is equal to the number of the background \( D3 \) charges or, equivalently, the products of NS–NS and RR flux as constructed in \( [36] \). \( N \) is taken to be relatively large. There can be a number of such throats in the compact region. As long as the tadpole cancellation imposed on the charge conservation is satisfied, there is no restriction on the number of throats. Following the convention of \( [20] \) we consider a scenario with at least two throats: the \( A \) throat, where the \( D3 \)-brane is located, and the \( S \) throat, where the standard model branes are located. As a \( D3 \)-brane moves towards the \( D3 \)-brane at the bottom of the \( A \) throat, inflation takes place. The position of the \( D3 \)-brane is \( r \) in the six-dimensional compact space. Around a throat, we choose the coordinate with respect to the bottom of the throat. This allows us to consider only \( r = |r| \). Note that \( \phi = \sqrt{T_3 r} \) is the inflaton, where \( T_3 \) is the \( D3 \)-brane tension.

In another scenario, the \( D3 \)-brane can simply be a mobile brane in the bulk, or it can be produced in another throat (the \( B \) throat) via the flux–brane annihilation mechanism \( [38] \). For the \( D3 \)-brane to move out of the \( B \) throat, \( \beta \) there must be negative. Each of these throats has its own warping which we denote by \( h_A, h_B \) and \( h_S \) respectively.

In slow roll inflation, the \( D3 \)-brane is moving slowly so that the square root term in DBI action of the \( D \)-brane world volume action can be expanded up to quadratic order. One can easily see that the Chern–Simons and the gravitational part of the DBI action for the \( D3 \)-brane cancel each other and we are left with a canonically normalized kinetic term for \( \phi \).

There are a number of contributions which affect the motion of the \( D3 \)-brane either in the bulk or in a throat. An important one comes from the Kähler modulus stabilization \( [11] \). Other contributions come from a variety of sources \( [25] \). The leading term of these contributions on the effective potential has the form

\[ V_K = \frac{1}{2} \beta H^2 \phi^2, \tag{8} \]

where \( \beta(\rho, \phi) \) is a function of the Kähler modulus \( \rho \) and depends on where the brane is sitting in the compactified manifold. Around any particular throat, the \( \phi \) dependence is small and may be neglected, so we can treat \( \beta \) as a constant.

We assume that there is an \( D3 \)-brane located in the \( A \) throat. Its tension provides the vacuum energy for inflation, given by

\[ V_A = 2h_A^4 T_3 = \frac{2h_A^4}{(2\pi)^3 g_s \alpha'^2}. \tag{9} \]

The factor of 2 is due to equal contributions of the Chern–Simons and the gravitational part of the action for the \( D3 \)-brane. The factor \( h_A^4 \) represents the warping effect. The resulting inflaton potential for this inflaton field comes from equations (8) and (9) and
from the attractive potential \( V_{D^3\bar{D}^3} \) between the \( D^3 \)-and the \( \bar{D}^3 \)-branes,

\[
V = \frac{1}{2} \beta H^2 \phi^2 + 2 h_A^4 T_3 + V_{D^3\bar{D}^3}
\]

\[
= \frac{1}{2} \beta H^2 \phi^2 + \frac{64 \pi^2 \phi^4_A}{27 N} \left( 1 - \frac{\phi^4_A}{N \phi^2} \right)
\]

(10)

where \( \phi_A = \sqrt{T_3 r_0} \) is the location of the \( \bar{D}^3 \)-brane at the bottom of the \( A \) throat. In the KKLMMT scenario, \( \beta \) is fine-tuned to 1 part in 100 so that its effect is negligible. In this case, the doubly warped interaction \( V_{D^3\bar{D}^3} \) is readily compatible with the slow roll conditions.

At the end of inflation, the \( D^3 \)-brane collides with the \( \bar{D}^3 \)-brane at the bottom of the \( A \) throat. Part of the energy produced will be transferred to the \( S \) throat to start the hot big bang, while \( D^1 \)-branes (i.e., \( D \)-strings) and fundamental closed strings are also produced. The quantity of interest is the \( D \)-string tension \( G\mu \), where \( \mu \) is the effective tension measured from the four-dimensional effective action point of view. It is related to the intrinsic tension \( T_1 \) of the \( D^1 \)-brane via

\[
G\mu = GT_1 h_A^2 = \sqrt{\frac{1}{32 \pi g_s}} \left( \frac{T_3}{M_p^2 h_A^4} \right)^{1/2}
\]

(11)

where \( G^{-1} = 8\pi M_p^2 \) and \( g_s \) is the string coupling.

3. The scenario

Here we investigate the range of \( \beta \) allowed by the inflationary constraints, in particular the slow roll condition in the \( A \) throat. Within the allowed range of \( \beta \), we evaluate the cosmic string tension and various inflationary parameters. The KKLMMT scenario is reproduced in the \( \beta \to 0 \) limit.

To satisfy the slow roll condition, the slow roll parameter \( \eta \) must be small enough:

\[
\eta = M_p^2 \frac{V''}{V} = \frac{\beta}{3} - \frac{20}{N_A} \frac{M_p^2 \phi_A^4}{\phi^6}.
\]

(12)

It is easy to check that the other slow roll parameter \( \epsilon \) is always small, so we shall ignore it in the discussion on slow roll. Inflation ends when \( \eta \) ceases to be small. So we can determine the final value of the inflaton field, \( \phi_i \) when \( \eta \sim -1 \), which gives

\[
\phi_i^6 = \frac{1}{(1 + \beta/3)} \left( \frac{20}{N_A} M_p^2 \phi_A^4 \right).
\]

(13)

The number of e-folds is given by

\[
N_e = \frac{M_p^2}{V} \int \frac{V d\phi}{V'}
\]

(14)

so the value of \( \phi \), namely \( \phi_i \) at \( N_e \) before the end of inflation is given by

\[
e^{2\beta N_e} = \frac{\beta N_A \phi_i^6 + 12 M_p^2 \phi_A^4}{\beta N_A \phi_i^8 + 12 M_p^2 \phi_A^4}.
\]

(15)
Combining equations (13) and (15), we find

\[
\phi_i^6 = \frac{24N_c}{N_A} M_P^2 \phi_A^4 \Omega(\beta)
\]

\[
\Omega(\beta) \equiv \frac{(1 + 2\beta) e^{2\beta N_e} - (1 + \beta/3)}{2\beta(N_e + 5/6)(1 + \beta/3)}
\]

\[
\approx \frac{e^{2\beta N_e} - 1}{2\beta N_e}
\]

(16)

where \(\Omega(\beta)\) is normalized such that \(\Omega = 1\) as \(\beta \to 0\). The last formula is a good approximation for small \(\beta\). The density perturbation \(\delta_H = 1.9 \times 10^{-5}\) measured by means of cosmic microwave background radiation [2] is given by

\[
\delta_H \equiv \frac{1}{\sqrt{75\pi M_P^2}} \frac{1}{V^3/2} \Omega(\beta)^{1/2}
\]

\[
= \left(\frac{2^{11}}{3 \times 5^6 \times \pi^4}\right)^{1/6} N_e^{5/6} \left(\frac{T_3}{M_P^4 h_A^4}\right)^{1/3} f(\beta)^{-2/3}
\]

(17)

where \(f(\beta)\) is given by

\[
f(\beta) = \left[\frac{2\beta(N_e + 5/6)}{(1 + 2\beta) e^{2\beta N_e} - (1 + \beta/3)}\right]^{5/4} \frac{(1 + 2\beta)^{3/2}}{(1 + \beta/3)^{1/4} e^{\beta N_e}}
\]

\[
\approx \left[\frac{2\beta N_e}{e^{2\beta N_e} - 1}\right]^{5/4} e^{\beta N_e}
\]

(18)

where the last formula is a good approximation for small \(\beta\). In the limit \(\beta \to 0\), \(f(\beta) \to 1\) and our results reduce to those of KKLMMT.

One can easily show that

\[
\frac{\phi_i}{\phi_R} = \left(\frac{32\pi^2}{27 N_A}\right)^{1/4} \left(\frac{T_3}{M_P^4}\right)^{-1/4} \left(\frac{\phi_i}{M_P}\right)
\]

(19)

where

\[
\frac{\phi_i}{M_P} = \left(\frac{3^{13} \times 5^6}{2^{15}}\right)^{1/12} \Omega(\beta)^{1/6} f(\beta)^{1/3} N_e^{-1/4} \delta_H^{1/2}
\]

(20)

and \(\phi_R \equiv \sqrt{T_3 R}\). Now we can write

\[
\eta = \frac{\beta}{3} \frac{5}{6N_e} \Omega(\beta)
\]

\[
\epsilon = \frac{M_P^2}{2} \left(\frac{V'}{V}\right)^2 = \frac{1}{18} \left(\frac{\phi_i}{M_P}\right)^2 \left(\beta + \frac{1}{2N_e \Omega}\right)^2
\]

\[
\zeta = \frac{M_P^4}{V^2} \left(\frac{V''}{V}\right)^2 = \frac{5}{3N_e \Omega} \left(\beta + \frac{1}{2N_e \Omega}\right)
\]

(21)

where \(\phi_i/M_P\) is given in equation (20). Note that, given COBE normalization \(\delta_H = 1.9 \times 10^{-5}\) at \(N_e \simeq 55\), \(\eta, \epsilon\) and \(\zeta\) are functions of \(\beta\) only. In general, \((\phi_i/M_P)^2 \beta^2 \ll 1\), so \(\epsilon\) may be neglected. Let us focus on this case, where the slow roll condition is guided
by the slow roll parameter \( \eta \). Due to the opposite sign of the two contributions in \( \eta \), we see that the presence of the \( \beta > 0 \) term actually improves the slow roll condition. This property allows a small but not necessarily fine-tuned \( \beta \). The \( \eta \) at \( N_e \) before the end of inflation measures the deviation from the scale-invariant power spectrum. The power spectrum index \( n_s \) at \( N_e \simeq 55 \) is given by

\[
\begin{align*}
    n_s - 1 &\simeq 2\eta - 6\epsilon \simeq \frac{2\beta}{3} - \frac{5}{3N_e \Omega(\beta)} \\
\end{align*}
\]

and is plotted as a function of \( \beta \) in figure 1. For the sake of completeness, both the running of \( n_s \),

\[
\frac{d n_s}{d \ln k} = 16\epsilon \eta - 24\epsilon^2 - 2\zeta ,
\]

and the ratio of the tensor to the scalar mode perturbation, \( r = 12.4\epsilon \), are plotted as functions of \( \beta \) in figure 2. Although the ratio \( r \) is small, its value \( r(\beta \simeq 0.1) \simeq 10^{-4} \) is larger by about five orders of magnitude than the value obtained in KKLMMT limit, where \( r \sim 10^{-9} \). This provides another good test of the scenario.

The cosmic string tension is given by

\[
G_\mu = \left( \frac{3 \times 5^6 \times \pi^2}{2^{21}} \right)^{1/4} g_s^{-1/2} \delta_H^{3/2} N_e^{-5/4} f(\beta) = G_{\mu_0} f(\beta) \]

where \( \mu_0 \) is the cosmic D-string tension at \( \beta = 0 \) (the KKLMMT scenario) and has the value \([11,12]\)

\[
G_{\mu_0} \simeq 4 \sqrt{g_s} \times 10^{-10} .
\]

Both \( G_\mu \) and \( n_s \) as a function of \( \beta \) are shown in figure 1. The range of values for \( G_\mu \) is given in equation (3), where the upper bound comes from WMAP data \([10,31]\). Actually, this bound is obtained for Abelian Higgs model vortices. For \((p,q)\) strings, the analysis is clearly more complicated. It is reasonable to assume that the \((p,q)\) string network also evolves to a scaling solution. It is also reasonable to assume that the high tension (i.e., high \((p,q)\)) string densities are suppressed, so that the range (3) may be approximately applied to \( G_\mu \) for the F- and/or the D-string.

Next, we need to find the bounds on \( \beta \) such that the constraints from the total number of e-folds (which must be larger than \( N_e \), though in general, we expect it to be much larger than \( N_e \)), the COBE normalization and the consistency of the inflationary model building in the A throat are satisfied. It turns out that the last condition puts the most stringent upper bound on \( \beta \). We demand that

\[
\phi_A < \phi_t < \phi_i < \phi_R .
\]

This condition is to make sure that the last 55 e-folds take place while the D3-brane is moving inside the A throat. The condition (19) can easily be satisfied. One may want to impose the more common condition \( \phi_i/M_P < 1 \). Using equation (20), one finds for \( N_e = 55 \), \( \beta \lesssim 1/6 \).

\[3\] After correcting a typo in (C.9) in [11].
As mentioned above, COBE normalization and $n_s$ put weaker bounds on $\beta$. To see that, we note from equation (17) that $\delta_H$ is given in terms of three undetermined parameters: $T_3/M_P^4$, $h_A$ and $\beta$. One can get a crude upper bound for $\beta$ as follows. One expects that $T_3/M_P^4 < 1$ and $h_A < 1$. Consider the extreme case when $h_A^4 T_3/M_P^4 = 1$. For $N_e \simeq 55$ and $\delta_H = 1.9 \times 10^{-5}$, we find $\beta < 0.46$. The constraints of WMAP data on inflationary parameters are presented in [29]. The observational bound on $n_s$ for $\eta > 0$ and $n_s > 1$ is $1.00 \leq n_s \leq 1.28$ [29]. This in turn implies that $\beta < 0.4$. The observational bound on $r$ is $r \leq 0.89$ [29]. This gives the bound $\beta < 0.22$. The observational bound on $dn_s/d\ln k$, however, is more restrictive. If we take the bound from WMAP $+$ 2dFGRS $+$ Lyman $\alpha$ on $dn_s/d\ln k$ seriously, we find that $\beta \lesssim 1/22$ (see figure 2). For $\beta \simeq 1/22$, $G\mu \simeq 10^{-8}/\sqrt{\sigma}$. However, this bound will be relaxed if the number of parameters to be fitted is increased. Also, if one uses only the WMAP data on $dn_s/d\ln k$, the bound is not restrictive at all. Using the WMAP $+$ 2dFGRS data, the bound on $\beta$ is $\beta \lesssim 1/5$, which is still less stringent than the other bounds. In the recent analysis of WMAP data combined with SDSS Lyman $\alpha$ and other data [30], the inflationary parameters as well as other parameters are determined. Depending on how the data are fitted, $n_s$ ranges from $0.98 \pm 0.02$ to $1.020 \pm 0.033 + 0.066 - 0.061$, where the last two numbers are $2\sigma$ uncertainties (note that $2\sigma$ is not twice $1\sigma$). The lower value of $n_s$ imposes no constraint on the model. If the upper bound on $n_s$ is taken to be $2\sigma$ above the larger central value, we find $n_s \leq 1.086$, or around $\beta \leq 1/7$. The bounds on $\beta$ from the bounds on $dn_s/d\ln k$ and $r$ are not as restrictive. Here, we shall use $\beta \leq 1/7$. However, one may be concerned about the astrophysical uncertainties involved in applying the Lyman $\alpha$ data here. This issue should wait for more data as well as analysis.

Having obtained the upper bound on $\beta$, we now determine the lower bound on it. The strongest lower bound comes from the requirement that the $D3$-brane moving in the outskirt of the throat eventually moves toward the $\overline{D3}$-brane so inflation can end. Since the $D3$-brane is moving very slowly this can be achieved by an attractive force, which means $V'_{|\phi=\phi_3} > 0$. We have

$$V'_{|\phi=\phi_3} = \frac{\phi_R V_A}{3M_P^4} \left( \beta + 16\pi \sqrt{\frac{2}{3}} N_A^{-3/2} h_A^4 \left( \frac{T_3}{M_P^4} \right)^{-1/2} \right).$$

(27)

The combination $N_A^{-3/2} h_A^4$, independent of the details of model, is very small. This indicates that an attractive force at the beginning of the throat for a negative $\beta$ is not possible unless $\beta$ is exponentially close to zero. This reduces to the KKLMMT scenario. To conclude, we find that $0 \lesssim \beta \lesssim 1/7$.

Let us summarize the situation with the parameters. With $N_e \simeq 55$, there are at least five parameters, namely, $h_A$, $g_s$, $N_A$, $\beta$ and $T_3/M_P^4$. The physically measurable parameters are $G\mu(\beta, g_s, \delta_H(\beta, h_A^4 T_3/M_P^4), n_s(\beta), r(\beta, \delta_H))$. It is interesting to realize that the measurable cosmological quantities are insensitive to $h_A$ and $N_A$. For example, using equation (17), it is easy to show that $h_A(T_3/M_P^4)^{1/4}$ is approximately equal to $10^{-4}$ and $10^{-3}$ at $\beta = 0$ and 0.1, respectively. This indicates that a range of $10^{-3} \leq h_A \leq 10^{-1}$ is possible with the appropriate value of $T_3/M_P^4$. An upper bound on $N_A$ can be obtained by imposing that inflation ends before the distance between $D3$ and $\overline{D3}$ reaches $\sqrt{\alpha'}$. At this separation, the tachyon appears and inflation must end as in hybrid inflation. However, inflation can end earlier. Using metric (5) to calculate the
and the bottom of the $B$ throat, so in equation (29) we set $A_B=\bar{A}_B$. Also note that the exponential dependence here is weak because $\beta N_{e}/3 \sim 1$.

The COBE normalization and the cosmic string tension, respectively, are

$$
\delta_H = \left( \frac{9^7}{3\pi^2} \right)^{1/4} N_B^{-1/4} |\beta|^{-1} \phi |\beta| N_{e}/3 h_A^2 \left( \frac{T_3}{M_P^4} \right)^{1/4} 
$$

(30)

$$
G_{\mu} = \left( \frac{3}{2^{17}} \right)^{1/4} N_B^{-1/4} |\beta| e^{-|\beta| N_{e}/3} \left( \frac{T_3}{M_P^4} \right)^{1/4} \delta_H. 
$$

(31)

One can see that with sufficient warping the observational value for $\delta_H$ is easily obtained. On the other hand, the degeneracy of the combination $h_A^3 T_3/M_P^4$ is now broken in the $G_{\mu}$ expression in the $B$ throat. So far in all calculations in the $A$ throat, $h_A$ and $T_3/M_P^4$, although independent parameters of the model, appeared jointly. In the $B$ throat inflation, however, the location of the cosmic string and the period of the inflation

$$
\frac{\phi_t}{\phi_A} \geq \exp(\sqrt{\alpha}/R).
$$

(28)

Using equations (13) and (17) in the above expression, we find

$$
N_A \exp((4/27\pi g_s N_A)^{1/4}) \leq \left( \frac{2^{29}}{5^2 \times 3^7} \right)^{1/6} \frac{N_0^{5/6}}{(1 + \beta/3)^{2/3} f(\beta)^{2/3} \delta_H}. 
$$

For example this gives $N_A \lesssim 10^4$ at $\beta = 1/6$ and $\lesssim 10^7$ at $\beta = 0$. Recall that $N_A \gg 1$ in the supergravity approximation that we are using. Typically $N_A \geq 100$ should be sufficient.

4. Slow roll in the $B$ throat

In this section we consider the case where $\beta$ is not small enough in the $A$ throat. We imagine that the slow roll inflation takes place in the $B$ throat which has sufficiently small $\beta$. We begin with the inflaton potential (1). $V_A$ is as given in equation (9), coming from the $D3$-brane locating at the bottom of the $A$ throat. $V_K$ has the same form as in equation (8) but with $\beta \rightarrow \beta_B$. The interaction term, $V_{DD}$, has the combination of warp factors $h_A$ and $h_B$ with the Coulombic form $1/(\phi - d)^4$ where now $\phi$ is measured with respect to the bottom of the $B$-throat and $d$ is the distance between the $D3$-brane and the $\bar{D}3$-brane. The exact form of $V_{DD}$ is a little complicated, but we do not need it in our analysis. This is justified by noting that if the $A$ throat and the $B$ throat are not overlapping, $d > R_A + R_B$, one can assume that $V_K$ dominates over $V_{DD}$ while the $D3$-brane is moving out of the $B$ throat. When it exits the $B$ throat, $\phi \sim \sqrt{T_3 R_B}$, $V_{DD}$ starts to become comparable to $V_K$ and we take that as the end of $B$ throat inflation. For the $D3$-brane to move out of the $B$ throat, it is necessary that $\beta_B < 0$, as in [20].

The slow roll parameter $\eta \simeq \beta/3$, so to satisfy the slow roll condition, one needs $|\beta| \leq 1/20$. The number of e-foldings is given by

$$
\phi_t = \sqrt{T_3 R_B} \exp(-|\beta| N_{e}/3). 
$$

(29)

As explained above, it is assumed that inflation ends when the $D3$-brane reaches the top of the $B$ throat, so in equation (29) we set $\phi_t = \sqrt{T_3 R_B}$. Also note that the exponential dependence here is weak because $\beta N_{e}/3 \sim 1$.

The COBE normalization and the cosmic string tension, respectively, are
are in two different throats which in turn results in the separation of $h_A^4$ from $T_3/M_P^4$ in the $G_\mu$ expression. With equation (31), it is easy to obtain $G_\mu$ ranging from $10^{-10}$ to $10^{-7}$, depending on the value of $T_3/M_P^4$. For example, taking $T_3/M_P^4 \sim 10^{-7}$, $\beta \sim 1/20$, $N_B \sim 200$, one finds $G_\mu \sim 1.8 \times 10^{-8}$. Decreasing $T_3/M_P^4$ to $10^{-9}$ reduces $G_\mu$ to $10^{-9}$.

Due to the attractive NS–NS plus RR forces, $D3$-branes in the bulk tend to move towards the bottom of throats. With enough $D3$-branes falling into the $B$ throat, brane–flux annihilation takes place [38]. Recall that $K$ is the NS–NS background flux and $M$ is the RR background flux of the throat. For $p \lesssim M$, $pD3$-branes can cluster to form a maximal size fuzzy NS-5-brane which then rolls down, leading to the brane–flux annihilation, i.e., the disappearance of the $pD3$-branes. This results in $K \rightarrow K - 1$ and the production of $M - pD3$-branes. So it is quite generic to expect a number of $D3$-branes being produced in the $B$ throat [20]. Inflation takes place as these $D3$-branes slowly move out of the $B$ throat. One may visualize a spray of $D3$-branes leaving the $B$ throat. They may enter different throats in the compact region. Some may enter the $A$ throat while some may enter the $S$ throat. With $k$ numbers of $D3$-branes colliding with $\bar{k}$ numbers of $D3$-branes in the $S$ throat, the resulting annihilation will heat up the remaining $\bar{k} - k$ standard model $D3$-branes. This way, heating of the standard model branes in the $S$ throat happens more or less simultaneously with the annihilation of branes in the $A$ throat. This mechanism may help to solve the reheating problem.

5. Fast roll inflation

Two different kinds of inflationary scenarios in string theory are studied in the literature, namely, the conventional ‘slow roll’ inflation as in the KKLMMT model and the ‘relativistic’ fast roll inflation in [39,35] and [20]. Unlike the slow roll inflation case, one must retain the DBI action for the kinetic term in fast roll inflation. It seems that the fast roll inflation suffers a serious problem regarding the COBE normalization and the non-Gaussianity constraints from the WMAP data. Here we shall comment on Chen’s multi-throat case briefly.

We begin with the DBI action coupled to the gravity

$$S = \frac{1}{2} \int dx^4 M_P^2 \sqrt{-g} (R - 2V) + S_{\text{DBI}},$$

where

$$S_{\text{DBI}} = \int dx^4 \sqrt{-g} f(\phi)^{-1} \left( \sqrt{1 + f(\phi)g^{\mu\nu}\partial_\mu \phi \partial_\nu \phi} - 1 \right)$$

where as in [35]

$$f(\phi) = \frac{\lambda}{\phi^4}.$$  

Here, $\phi = \sqrt{T_3} r$ and $\lambda \equiv T_3 R^4 = 27 N/32 \pi^2$.

$V$ is now given by $V = V_A + V_K$. Here we consider that the relativistic $D3$-brane is moving out of the $B$ throat, which means $\beta \sim -1$ [20]. The positive $\beta$ scenario was studied in [35].
As discussed in appendix B of [35], the condition that $V$ dominates over the kinetic energy requires
\[
\frac{V^{3/2}}{V'M_P} \sqrt{\frac{3f}{g_s}} \gg 1
\] (35)
while the condition for relativistic motion ($\gamma \gg 1$) is
\[
\frac{g_s}{3V} V'^2 f M_P^2 \gg 1.
\] (36)
The number of e-folds is given by
\[
N_e = \int \frac{d\phi}{M_P} \sqrt{\frac{fV}{3g_s}} \sim \sqrt{\frac{2\lambda}{3g_s}} \left( \frac{T_3}{M_P^2 h_A} \right)^{1/2} \left( \frac{M_P}{\phi_i} \right).
\] (37)
The COBE normalization is given by
\[
\delta_H = \frac{1}{15\pi} \sqrt{\frac{fV}{g_s M_P^2}} = \frac{2}{15\pi} \sqrt{\frac{\lambda}{g_s}} \left( \frac{T_3}{M_P^2 h_A} \right) \left( \frac{M_P}{\phi_i} \right)^2.
\] (38)
Equations (37) and (38) can be combined to yield
\[
\frac{\lambda}{g_s} = \frac{N_e^4}{25\pi \delta_H^2}.
\] (39)
Taking $\delta_H^2 \sim 10^{-10}$ and $N_e = 55$, we find $\lambda/g_s \sim 10^{14}$, or the background $D$-brane charge to have an extremely large value:
\[
N \sim 10^{14}
\] (40)
which implies a fine-tuning. A similar result was also obtained for the case in [35]; they suggest that a non-trivial orbifold can reduce this number to a more reasonable value. It would be interesting to see how such a large effective background charge can be achieved in a realistic model. It is also possible that the value of $N$ may be substantially reduced by a more careful analysis of the brane collision in the infrared end of the throat [40]. This interesting possibility needs further investigation. In any case, such scenarios have large non-Gaussianity effects, thus offering a good signature for detection.

The cosmic string tension, using equations (11) and (37), is
\[
G\mu = \frac{N_e}{8} \sqrt{\frac{3}{\lambda \pi}} \left( \frac{\phi_i}{M_P} \right) \sim 7 \left( \frac{\phi_i}{M_P} \right) \times 10^{-7}.
\] (41)
Depending on the value of $\phi_i/M_P$, $G\mu$ as big as $10^{-7}$ is possible, while $G\mu$ as small as in KKLMMT corresponds to $\phi_i/M_P \sim 10^{-3}$, which requires some tuning.
6. Conclusion

One may consider a multi-throat brane inflationary scenario where the brane annihilation takes place in a (the $A$) throat other than the throat where the standard model branes are sitting (the $S$ throat). This multi-throat scenario has the advantage that the warped geometry [21] of the standard model throat allows us to solve the hierarchy problem while the warping of the $A$ throat allows us to fit the inflationary data separately. To summarize, we study a simple generalization of the original KKLMMT scenario by allowing a small but not necessarily very small conformal coupling. More specifically, we consider the following cases:

(1) The original KKLMMT scenario is reviewed, where $|\beta| \ll 1/10$, so that the $\beta$ term in the inflaton potential may be neglected. This requires a fine-tuning. As we relax this fine-tuning to $\beta \lesssim 1/6$, slow roll inflation still works fine, while the cosmic string tension increases rapidly. In this very simple generalization of the KKLMMT scenario, the cosmic string tension can easily saturate the observational bound. As the fine-tuning is relaxed, the density perturbation power spectrum moves from red-shifted to blue-shifted, and $r$ increases exponentially. We also note that the predictions of inflationary quantities and the cosmic string tension are insensitive to the warping $h_A$.

(2) If $\beta$ of the $A$ throat is not small, inflation may take place in another throat where $\beta$ is small and negative. This can be justified by the brane–flux annihilation picture of [38]. In this scenario, the $D3$-brane moves out of that ($B$) throat and then drops into the $A$ throat. Most (if not all) inflation takes place when the brane moves out of the $B$ throat. In a more general scenario, we expect a number of $D3$-branes moving out of the $B$ throat, which may end up in different throats. It is also possible that inflation happens when the $D3$-brane is moving in the bulk towards the $A$ throat.

(3) Suppose that $|\beta| \sim 1$. The positive $\beta$ case is studied in [35] while the negative $\beta$ case is studied in [20]. In both cases, there is no slow roll (in fact, the brane is moving very relativistically) and it is possible to achieve enough inflation. On the other hand, fitting the density perturbation data requires $N \sim 10^{14}$, a fine-tuning, although this fine-tuning may be substantially ameliorated with non-trivial orbifolding.

Since slow roll and fast roll inflations are both possible, a general analysis of an arbitrary roll will be very informative. In the end, we expect the inflationary constraints on $\beta$ and the value of the cosmic string tension to be rather relaxed.

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