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The inflationary brane-antibrane universe

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Abstract: We show how the motion through the extra dimensions of a gas of branes and antibranes can, under certain circumstances, produce an era of inflation as seen by observers trapped on a 3-brane, with the inflaton being the inter-brane separation. Although most of our discussion refers to arbitrary p-branes, when we need to be specific we assume that they are D-branes of type-II or type-I string theory. For realistic brane couplings, such as those arising in string theory, the inter-brane potentials are too steep to inflate the universe for acceptably long times. However, for special regions of the parameter space of brane-antibrane positions the brane motion is slow enough for there to be sufficient inflation. Inflation would be more generic in models where the inter-brane interactions are much weaker. The spectrum of primordial density fluctuations predicted has index n slightly less than 1, and an acceptable amplitude, provided that the extra dimensions have linear size $1/r \sim 10^{12}$ GeV. Reheating occurs as in hybrid inflation, with the tachyonic instability of the brane-antibrane system taking over for small separations. The tachyon field can induce a cascade mechanism within which higher-dimension branes annihilate into lower-dimension ones. We argue that such a cascade naturally stops with the production of 3-branes in 10-dimensional string theory.

Keywords: D-branes, Cosmology of Theories beyond the SM, Physics of the Early Universe.

Dedicated to the memory of Detlef Nolte

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Contents

1. Introduction and summary

Inflationary cosmology [1] provides a compelling explanation for many otherwise unnatural features of the observed early universe. However, no equally compelling microscopic model of the dynamics underlying inflation has yet emerged. All extant models seem to require fine-tuning in order to reproduce all of the features which successful inflation must have: sufficient universal expansion, an acceptable density spectrum and amplitude of primordial density fluctuations, and so on [2].

The relatively recent realization that spacetime may be full of p-branes — or $(p+1)$-dimensional defects — on which most or all of the observed particle species may be trapped [3], provides a new framework within which to search for natural mechanisms for inflation. Indeed, for many choices of parameters (like the dimension of space, and the energy scale of brane physics) inflationary cosmology may provide one of the few constraints we have on the various possible brane constructions. This has inspired several investigations of how inflationary ideas might be rooted in the brane-world framework [4, 5, 6].

In this paper we examine the possibility that inflation can be due to the motion of branes and antibranes[1] within the extra dimensions, with inflation ending when branes and antibranes collide and annihilate. Other scenarios were considered before in terms of only branes, without antibranes. These earlier analyses either made phenomenological assumptions about the nature of brane-brane interactions [5], or

[1] An antibrane is an extended object which has the same tension as the corresponding brane, but with opposite Ramond-Ramond (R-R) charge.

were performed assuming the collisions of BPS branes in M-theory [6], for which the interactions are known to vanish to all orders in perturbation theory and so for which some form of non-perturbative potential must be assumed.

Brane-antibrane collisions are natural candidates for inflation, since branes and antibranes interact with one another in a way which has long been well understood at large distances [7], and since their annihilation might be expected to provide a relatively efficient mechanism for converting the energy of relative brane motion into reheating the universe. Furthermore, much progress has been made of late in understanding their interactions at short distances within string theory [8]. (See [9] for a discussion of some inflationary issues associated with short-distance brane interactions.)

In the calculations to be presented here, we explore inflation due to brane-antibrane collisions for branes of arbitrary dimension, using what is known about their long- and short-distance interactions. In a nutshell, our main findings are these:

- We find that if the brane-antibrane couplings are very weak then their motion towards one another is slow enough for sufficient inflation to occur for generic initial brane positions. Unfortunately, the couplings derived from string theory are too strong for this mechanism to work for generic initial conditions, since the branes and antibranes collide too quickly to provide sufficient inflation. Inflation of this type, if it occurs, gives rise to a spectrum of primordial density fluctuations with a spectral index $n \gtrsim 1$.

- For string theory couplings inflation can nonetheless arise, but only if the initial brane-antibrane configuration starts near special parts of the extra dimensions where their interactions are particularly weak. We identify these special initial conditions explicitly for toroidal compactifications and show that they are sufficiently large for inflation not to be unlikely given a reasonable density of branes and antibranes in the extra dimensions. Primordial density fluctuations are also predicted to have $n \gtrsim 1$, and have the correct amplitude in this scenario if the compactification scale is of order $1/r \sim 10^{12}$ GeV.

- Once the brane and antibrane get sufficiently close to one another they very naturally enter a hybrid inflation regime [10], where the slow change in inter-brane separation becomes replaced with a quick roll down a tachyonic direction. (We have checked that the tachyonic roll is too quick to provide sufficient inflation on its own, without the inter-brane relative motion.) This provides an efficient release of the energy of relative motion into reheating.

- Brane-antibrane collisions do not always annihilate to a no-brane final state, so the energy released by reheating can easily have significant portions trapped on the daughter brane. This is a good feature since we imagine all standard-model particles to be so trapped, and so the reheat energy can plausibly find its way into our observable universe. Typical reheat temperatures are of order 10^{13} GeV.

- Since brane-antibrane annihilation typically produces branes of lower dimension, it would be natural to have a cascade of annihilations where higher dimensional branes successively annihilate into lower ones. Furthermore, the endpoint of such a cascade may furnish a dynamical explanation for the present-day conditions which are widely assumed within brane scenarios. For instance, since larger-dimension branes can find one another more easily and annihilate, the cascade scenario qualitatively suggests a migration of the brane abundance towards lower-dimensional branes as the universe evolves. Furthermore, branes which are $(3+1)$-dimensional (or less) can plausibly fail to find one another throughout the history of the universe in a 10-dimensional theory, potentially providing an explanation within string theory for why the late universe might prefer to be populated by 3-branes. Finally, these final survivors of brane collisions can be argued to prefer to be aligned parallel with the large dimensions of space, as is usually assumed in phenomenological applications.

We now turn to the more detailed explanation of these results.

2. Inflation due to inter-brane motion

In this section we present formulae for the predictions of slow-roll inflation, where the slowly rolling inflaton is the separation between two branes, $y^\mu = x_1^\mu - x_2^\mu$. After first outlining our assumptions concerning the brane and bulk dynamics, and the nature of the interaction between branes, we describe two scenarios in which inflation might arise. Although both can provide viable inflation if the strength of the inter-brane interaction is allowed to be sufficiently weak, only the second is viable (see section 2.5) if the coupling strengths are taken from string theory.

2.1 The effective action

We take as our starting point the following low-energy effective action governing the inter-brane interactions: $S = S_B + S_{b1} + S_{b2}$, where the bulk action is[2]

$$S_B = -\int \mathrm{d}^4x \mathrm{d}^d y \sqrt{-g} \left[\frac{M_s^{2+d}}{2} e^{-2\phi} R + \cdots \right] . \tag{2.1}$$

[2]Our metric has signature $(- + \ldots +)$ and we use Weinberg's curvature conventions.

Here ϕ is the dilaton,[3] R is the full $(4+d)$-dimensional scalar curvature, and eq. (2.1) defines the fundamental (string) scale, M_s. (For string theory applications, $d = 6$.) The ellipses denote all other bulk massless fields which mediate forces between branes.

The dynamical variables of the brane actions are the brane positions, x_i^μ, with $i = 1, 2$ and $\mu = 0, \ldots, 3+d$. The relevant part of the action is:

$$S_{bi} = -\int \mathrm{d}^4x \mathrm{d}^{p-3}y \sqrt{-\gamma}[T_p + \cdots] \,, \tag{2.2}$$

where $\gamma_{ab} = g_{\mu\nu}\, \partial_a x_i^\mu \partial_b x_i^\nu$ is the induced metric on the brane, and $T_p = \alpha M_s^{p+1} e^{-\phi}$ is the brane tension, where α is a dimensionless constant which drops out of most of our results.

For our purposes we will assume the branes to be parallel, and choose coordinates for which the branes lie along coordinate surfaces. If we then separate the part of x_i^m transverse to the brane into the relative motion, $y^m = (x_1 - x_2)^m$, of the branes and the motion of their center-of-mass, $\overline{x}^m = (x_1 + x_2)^m/2$, then expanding in powers of $\partial_a y^m$ gives

$$S_{b1} + S_{b2} = -\int \mathrm{d}^4x \mathrm{d}^{p-3}y \sqrt{-\gamma}T_p \left[2 + \frac{1}{4} g_{mn}\gamma^{ab}\partial_a y^m \partial_b y^n + \cdots\right] . \tag{2.3}$$

2.2 Inter-brane interactions

We assume the dominant interaction between the branes at large distances (compared to $1/M_s$) to be due to the exchange of massless bulk modes, including the metric, dilaton and any other massless bosons which couple to the branes. Any such interaction gives an inter-brane potential energy (per unit brane area) which falls with separation like $1/y^{d_\perp - 2}$ where

$$d_\perp = d - p + 3 \,, \tag{2.4}$$

is the number of spatial dimensions transverse to the branes. In most models the couplings of these modes is gravitational in strength, so we parameterize the long distance potential energy by

$$\frac{E}{A_p} = -\beta \left(\frac{e^{2\phi}}{M_s^{2+d}}\right) \frac{T_p^2}{y^{d_\perp - 2}} \,, \tag{2.5}$$

where β is a model-dependent dimensionless number which characterizes the overall strength of the inter-brane force, and depends on the details of the number and kinds of bulk states which are exchanged.

By far the best motivated candidate for these branes are the p-dimensional BPS Dirichlet branes (Dp-branes) of type-II string theory [7]. In string theory, the potential for the brane-antibrane (Dp-$\overline{\mathrm{D}p}$) system is attractive, since the contributions

[3] For non-string applications $\phi = 0$ can be taken everywhere, although we shall find that ϕ drops out of most formulae.

due to dilaton/graviton exchange are equal to and reinforce the attraction due to the opposite R-R charges carried by these branes. This is to be contrasted with the Dp-Dp case where the dilaton/graviton contributions exactly cancel the repulsion due to the like-sign R-R charges. Direct calculation [7] gives

$$\beta = \pi^{-d_\perp/2}\Gamma\left[\frac{d_\perp - 2}{2}\right]. \tag{2.6}$$

2.3 The effective 4D inflaton

Several assumptions are required in order to analyze the implications of these interactions for inflating the 4 non-compact dimensions. Our most crucial assumption is that the moduli of the compact dimensions are stabilized by some unknown physics. This assumption precludes the dynamics of these moduli — including in particular the space's overall volume — from playing a role in inflation, allowing us to use the volumes $V_\perp \equiv r_\perp^{d-p+3}$ (transverse to the brane) and $V_{||} \equiv r_{||}^{p-3}$ (along the brane) as parameters which we may adjust to our own ends.

The volumes $V_\perp$ and $V_{||}$ may be freely adjusted, subject to two constraints. First, their product is related to the four-dimensional reduced Planck mass, $M_p = (8\pi G_N)^{-1/2} \sim 10^{18}$ GeV, by

$$M_p^2 = e^{-2\phi} M_s^{2+d} V_\perp V_{||} \tag{2.7}$$

as is seen by dimensionally reducing the bulk action, eq. (2.2). We use this expression in what follows to eliminate the volume $V_{||}$ in favor of M_p, leaving only $V_\perp$ as an adjustable parameter.

The second constraint we demand of these volumes is that they be large when measured in units of M_s: $M_s r_\perp \gg 1$, $M_s r_{||} \gg 1$. For instance, this is required within string theory in order to be able to treat the bulk dynamics using only the low-energy effective field theory corresponding to the massless string states.[4]

In what follows we find that the amplitude of primordial density fluctuations depends only on the combination $M_p r_\perp$ (and not separately on e^ϕ or M_s, for example), with agreement requiring $M_p r_\perp \sim 10^6$. Since weak-coupling in string theory requires $M_s < M_p$, we see that M_s may easily be chosen to satisfy the constraints $1/r_\perp \ll M_s < M_p$.

With these assumptions, we dimensionally reduce to recast the dynamics of the inter-brane separation, y, as that of a canonically normalized four-dimensional scalar field, which we denote by ψ. (It is effectively a scalar field because only one component of y is dynamical if the brane and antibrane move toward each other

[4]This is also the condition under which the nonlinear contributions of Einstein's equations are negligible when considering the gravitational field of a single D-brane, say. Therefore our approximate 4D effective field theory treatment is self-consistent.

along a straight line.) Inspection of the brane action, eq. (2.3), shows the canonically normalized field to be related to y by:

$$\psi = y\sqrt{\frac{T_p V_{\|}}{2}} = y\left[\frac{\alpha e^{\phi}}{2(M_s r_{\perp})^{d_{\perp}}}\right]^{1/2} M_s M_p\,. \tag{2.8}$$

The effective scalar potential for y, obtained by combining the tension term of eq. (2.3) with the inter-brane potential energy of eq. (2.5), becomes:

$$V(y) = A - \frac{B}{y^{d_{\perp}-2}}\,, \tag{2.9}$$

with

$$\begin{aligned} A &= 2T_p V_{\|} = \frac{2\alpha e^{\phi}}{(M_s r_{\perp})^{d_{\perp}}} M_s^2 M_p^2\,,\\ B &= \frac{\beta e^{2\phi}}{M_s^{2+d}} T_p^2 V_{\|} = \frac{\alpha^2 \beta e^{\phi} M_p^2}{M_s^{2(d_{\perp}-2)} r_{\perp}^{d_{\perp}}}\,. \end{aligned} \tag{2.10}$$

These expressions for $V(y)$ apply if the separation between the branes is large compared to $1/M_s$, but is not large enough to be comparable to the overall size of the compact transverse dimensions. In our application to string theory, we shall find that in order to obtain sufficient inflation, one requires inter-brane separations which are comparable with $r_{\perp}$. We therefore require the modifications to V which arise in this case.

Although we cannot provide explicit expressions for the modified V in general geometries, we can do so for simple cases like tori. The potential energy of a brane on a torus can be represented by summing the above potential over a lattice of "image" branes. As a brane is moved to an antipodal point from the source brane on a torus, it can interact equally strongly with two or more of the image charges. Inflationary calculations in this instance may be performed using the potential of eq. (2.9), but with the interaction term summed over the contributions of all images.

The potential (2.9) neglects several effects. First, it neglects dependence on the inter-brane speeds and orientations. The velocity dependence is legitimate for a slow roll, where speeds are assumed to be initially very small. Neglect of the dependence on orientation presumes the branes and antibranes to be moving strictly parallel to one another, and has been done for simplicity of analysis. We believe further study of the dependence of inflation on these variables to be warranted.[5]

The inter-brane potential also neglects angular momentum terms, corresponding to the assumption that the branes of interest are moving directly towards one another. Since the angular momentum, L, contributes to the radial potential as $V_L \sim L^2/y^2$, we see that it dominates the large-y limit unless there are three transverse dimensions

[5] See next section for a short discussion of these points.

(corresponding to $d-p=0$, e.g. 6-branes in 10 spacetime dimensions). We believe the neglect of angular momentum to be justified for the branes which initiate inflation, since these branes must move very slowly. Branes which attract and which start initially essentially from rest should approach each other directly.

2.4 The case $M_s^{-1} \ll y \ll r_\perp$

Let us first consider the case of a parallel p-brane and antibrane separated from each other by a distance y. We assume that the branes wrap a $(p-3)$-dimensional space of volume $V_{\|}$, and that the space transverse to the branes is compactified on a space of volume $V_\perp$. We will show that it is only possible to satisfy the slow-roll conditions for inflation in this set-up if the constant β parameterizing the strength of the inter-brane force is chosen sufficiently small. The β predicted for Dp-$\overline{\mathrm{D}p}$ branes in string theory (see eq. (2.6)) turns out to be too large to permit slow-roll inflation, precluding the generation of inflation using this mechanism in string theory.

We use the potential of eq. (2.9) which is valid when the inter-brane separation is large compared to $1/M_s$, but small compared with the scale $r_\perp$ of the compact transverse dimensions. The standard equations of motion of a Friedmann-Robertson-Walker universe are

$$\begin{aligned} \ddot{\psi} + 3H\dot{\psi} &= -\frac{\mathrm{d}V}{\mathrm{d}\psi}\,, \\ H^2 &= \frac{1}{3M_p^2}\left[V + \frac{\dot{\psi}^2}{2}\right], \end{aligned} \tag{2.11}$$

where H is the Hubble parameter. The slow roll conditions require the friction and potential terms dominate, i.e., $|\ddot{\psi}| \ll 3H|\dot{\psi}|$ and $\dot{\psi}^2 \ll V$. Consistency with the field equations requires small values for the two slow-roll parameters

$$\epsilon = \frac{M_p^2}{2}\left(\frac{V'}{V}\right)^2, \qquad \eta = M_p^2\frac{V''}{V}\,, \tag{2.12}$$

where the primes denote derivatives with respect to ψ. For slow-roll inflation, both $|\epsilon|$ and $|\eta|$ should be much less than 1.

In our case, $|\epsilon|$ is always smaller than $|\eta|$ and will be neglected in the following discussion. The η parameter is

$$\eta \approx -\beta(d_\perp - 1)(d_\perp - 2)\left(\frac{r_\perp}{y}\right)^{d_\perp}. \tag{2.13}$$

Here we see the requirement for small β. For the string theoretic value of β as in eq. (2.6), the slow-roll condition $|\eta| \ll 1$ would require a brane separation y much bigger than the size of transverse dimension $r_\perp$, which violates the assumptions under which eq. (2.13) was obtained.[6]

[6] On the last point, we disagree with ref. [11].

One might imagine attempting to satisfy the slow-roll condition using β from string theory in eq. (2.13) by using a compactification manifold which is asymmetric, with some (say, n) transverse dimensions much smaller than the others. This permits $y > r_\perp$ if the branes are separated along one of the long directions. This stratagem turns out to fail, however, although the analysis which shows this requires the inclusion of the corrections to the inter-brane potential which arise when $y \sim r_\perp$, such as we consider in the next section. The potential which results may be shown by a simple scaling argument — or by considerations of T-duality — to be modified in this case to $V \sim 1/a^n y^{d_\perp - n}$ (where a is the "small" radius). Repeating the calculation for this potential gives $y \gg (V_\perp/a^n)^{1/(d_\perp - n)}$, which is no improvement over eq. (2.13).

If we adopt at this point a phenomenological approach, and simply treat β as a model-dependent parameter, we may compute the predictions of the resulting slow-roll inflation.[7] The number of e-foldings occurring after the scales probed by the COBE data leave the horizon can be computed as

$$N = \int_{\psi_*}^{\psi_{\rm end}} \frac{H}{\dot\psi} \mathrm{d}\psi \approx \frac{(y_*/r_\perp)^{d_\perp}}{d_\perp(d_\perp - 2)\beta}\,, \tag{2.14}$$

where y_* is the brane separation when the primordial density perturbation exits the de Sitter horizon during inflation. We assume $y_* \gg y_{\rm end}$ in deriving this result. The number of e-foldings can be made sufficiently large ($N \lesssim 60$) by choosing small enough β.

The amplitude of this density perturbation when it re-enters the horizon, as observed by Cosmic Microwave Background (CMB) experiments, may be written as follows:

$$\delta_H \approx \frac{1}{5\sqrt{3}\pi} \frac{V^{3/2}}{M_p^3 V'} \approx \frac{2}{5\sqrt{3}\pi} \frac{(d_\perp N)^{(d_\perp - 1)/d_\perp}}{[(d_\perp - 2)\beta]^{1/d_\perp}} \frac{1}{M_p r_\perp}\,, \tag{2.15}$$

where the COBE results imply $\delta_H = 1.9 \times 10^{-5}$ [13].

The spectrum of perturbations is similarly calculated. For a (gaussian) perturbation $\delta(\vec{x}, t)$, it is conventional to Fourier transform it into the momentum space, $\delta(\vec{x}, t) = \int \mathrm{d}^3\vec{k}\,\delta(\vec{k}, t) \exp(\mathrm{i}\vec{k}\cdot\vec{x})$, where $\vec{k}$ is the wave vector in the comoving frame and k is its norm. The power spectrum $P_\delta(k)$ is then defined in terms of the two-point correlation functions $\langle \delta(\vec{k})\delta(\vec{k}')\rangle = \delta^3(\vec{k} - \vec{k}')(2\pi^2/k^3)P_\delta(k)$. The function $P_\delta(k)$ has the form $\sim k^{n-1}$ where n is known as the spectral index. In these conventions, inflationary models predict a scale invariant Harrison-Zel'dovich spectrum ($n = 1$) to high precision. In our case the spectral index and its derivative can be readily calculated

$$n - 1 \approx -\frac{2(d_\perp - 1)}{d_\perp N}\,, \qquad \frac{\mathrm{d}n}{\mathrm{d}\ln k} \approx -\frac{2(d_\perp - 1)}{d_\perp N^2}\,. \tag{2.16}$$

[7] Potentials similar to eq. (2.9) were considered in [12] with $B < 0$, inspired by nonperturbative supersymmetric potentials. The fact that $B < 0$ was the reason they found a blue spectrum, in contrast with our result.

2.5 The case $y \sim r_\perp$ and special points

In this subsection we study the case where the brane-antibrane separation is comparable to the size of compactified transverse dimensions. We will show that the difficulty encountered in the previous case is resolved by an interesting cancellation in the brane potential.

For our purposes, we assume the compactified transverse manifold to be a $d_\perp$-dimensional square torus with a uniform circumference $r_\perp$. When the brane-antibrane separation is comparable to $r_\perp$, we have to include contributions to the potential from p-brane images, i.e., we have to study the potential in the covering space of the torus, which is a $d_\perp$-dimensional lattice, $(\mathbb{R}/\mathbb{Z})^{d_\perp}$.

The potential at the position of the antibrane is

$$V(\vec{r}) = A - \sum_i \frac{B}{|\vec{r} - \vec{r}_i|^{d_\perp - 2}} , \tag{2.17}$$

where the parameters A and B are defined in eq. (2.10), $\vec{r}$ and $\vec{r}_i$ are the vectors denoting the positions of the p-branes and antibranes in the $d_\perp$-dimensional coordinate space, and the summation is over all the lattice sites occupied by the brane images, labelled by i. We schematically show our set-up in figure 1.

It is easy to show the first and second derivatives of the above potential vanish when the antibrane is located at the center of a hyper-cubic cell, which we denote as $\vec{r}_0$. To see the vanishing of these quantities, compute

$$\begin{aligned}
\left.\frac{\partial V}{\partial r_a}\right|_{\vec{r}=\vec{r}_0} &= (d_\perp - 2) \sum_i \frac{B(r_0 - r_i)_a}{|\vec{r}_0 - \vec{r}_i|^{d_\perp}} , \\
\left.\frac{\partial^2 V}{\partial r_a \partial r_b}\right|_{\vec{r}=\vec{r}_0} &= (d_\perp - 2) \sum_i \frac{B}{|\vec{r}_0 - \vec{r}_i|^{d_\perp}} \left[\delta_{ab} - d_\perp \frac{(r_0 - r_i)_a (r_0 - r_i)_b}{|\vec{r}_0 - \vec{r}_i|^2}\right] ,
\end{aligned} \tag{2.18}$$

where $a, b \in \{1, 2, \ldots, d_\perp\}$ are coordinate-axis indices. The first derivative obviously vanishes when evaluated at $\vec{r}_0$ because of a reflection symmetry of the lattice, showing that the net force on the antibrane at this point is zero.

It is perhaps more remarkable that the second derivative also vanishes at $\vec{r}_0$. The vanishing of the leading contribution may be verified by explicit summation of the $2^{d_\perp}$ nearest neighbors. More generally, the vanishing of the entire sum may be seen as follows. Consider the two cases, $a = b$ and $a \neq b$:

1. When $a = b$, given any lattice point $\vec{r}_1$, we can generate other lattice sites equidistant from $\vec{r}_0$ (if necessary, we choose a coordinate system such that $\vec{r}_0$ is the origin, with coordinate axes aligned along those of the lattice) by cyclically permuting the coordinates of $\vec{r}_1 - \vec{r}_0$. We can generate at most $d_\perp - 1$ new points in this way. If all these $d_\perp$ points (including $\vec{r}_1$) are distinct, then $\sum_{i=1}^{d_\perp} (r_0 - r_i)_a^2 = |\vec{r}_0 - \vec{r}_1|^2$ for all a, where the summation is over these lattice

sites. If instead only $\ell < d_\perp$ points are distinct (so that they transform under a $\mathbb{Z}_\ell$ subgroup of the $\mathbb{Z}_{d_\perp}$ group of cyclic permutations of the coordinates), then Lagrange's factor theorem for the dimension of a subgroup implies that $d_\perp \sum_{i=1}^{\ell}(r_0 - r_i)_a^2/|\vec{r}_0 - \vec{r}_1|^2 = \ell$ for all a, which cancels the contribution of the δ_{ab} term in eq. (2.18).

2. When $a \neq b$, the second derivative vanishes due to a reflection symmetry along any plane where, say, $(r_0 - r_i)_a$ is a constant.

We therefore conclude the second derivative of the potential is zero along any direction at this special point of the hyper-cubic cell. Of course, this conclusion depends upon our choice of the square lattice.

Consider now antibrane motion when the antibrane is near the center of a hyper-cubic cell. In this case we may expand the potential in terms of power series of this small displacement $\vec{z}$ from the center. From a simple symmetry argument, one can easily see that all the odd powers in $\vec{z}$ vanish. Furthermore, we have shown above that the quadratic terms in $\vec{z}$ also vanish, so that the leading contribution to the potential is the quartic term in $\vec{z}$.

We can therefore model the relative motion of the branes in a quartic potential

$$V(z) = A - \frac{1}{4}Cz^4\,, \tag{2.19}$$

where A is defined as before and

$$C = \gamma M_s^{-(2+d)} e^{2\phi} T_p^2 V_{\|} r_\perp^{-(2+d_\perp)} \tag{2.20}$$

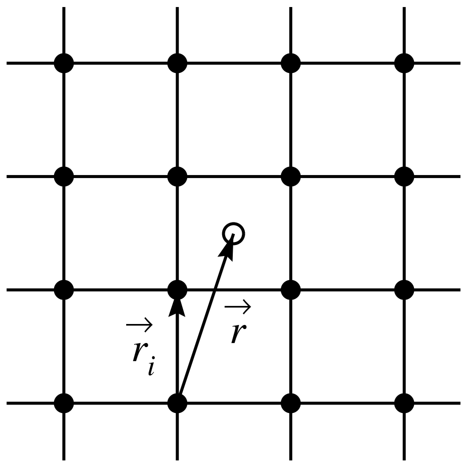

Figure 1: The covering space of a square torus. Solid circles represent images of the Dp-brane, and the open circle represents the $\overline{\mathrm{D}p}$-brane, which is near the center of a hyper-cubic cell during the slow-roll inflationary phase. The images of the antibrane are omitted because they do not contribute to the potential of the antibrane since the force between an antibrane and another antibrane is zero.

with γ a constant of order $\mathcal{O}(1)$. It is straightforward to derive for this potential the slow roll parameter η and the density perturbation δ_H

$$\eta \approx -3\gamma \left(\frac{z}{r_\perp}\right)^2, \qquad \delta_H \approx \frac{2}{5\pi}\sqrt{\frac{\gamma}{3}}\frac{N^{3/2}}{M_p r_\perp}\,, \tag{2.21}$$

where again we have used the standard slow-roll equations. We see now that slow-roll is guaranteed for sufficiently small z.

The value of the antibrane displacement from the center of the hyper-cubic cell at the time when the COBE perturbation leaves the horizon during the inflation is

$$z_* \approx \frac{r_\perp}{\sqrt{2\gamma N}}\,. \tag{2.22}$$

From this the spectral index and its derivative in this model can be easily determined

$$n - 1 \approx -\frac{3}{N}\,, \qquad \frac{\mathrm{d}n}{\mathrm{d}\ln k} \approx -\frac{3}{N^2}\,. \tag{2.23}$$

One can also compute similar quantities for primordial tensor perturbations.

Using $N \approx 60$ we find $z_* \approx 0.1\, r_\perp$ and the value of the displacement at the end of slow roll $z_{\mathrm{end}} \approx 0.5\, r_\perp$ (from $|\eta| \approx 1$). Both of these are fairly natural numbers. Comparing to the COBE normalization [13], $\delta_H \approx 1.9 \times 10^{-5}$, we obtain the size of the compactified transverse dimension

$$M_p\, r_\perp \approx 1.3 \times 10^6 \qquad \text{or} \qquad r_\perp^{-1} \approx 10^{12}\,\mathrm{GeV}\,. \tag{2.24}$$

The deviation of the spectral index n from the scale-invariant form is at a few percent level, and the spectrum is slightly red. $\mathrm{d}n/\mathrm{d}\ln k$ is at the order of $\mathcal{O}(10^{-3})$. These features of the predicted CMB spectrum are quite common for inflationary models, and are consistent with current data [14]. They will be tested more precisely in future satellite experiments such as MAP and Planck.

3. Collisions, reheating and cascade

In the final stages of brane-antibrane annihilation, many more degrees of freedom are relevant than simply the inter-brane separation. This is a good thing, since these provide many channels for reheating which are not constrained by the flatness of the potential responsible for the initial slow roll of the inflationary phase. In this general sense brane collisions are an extreme example of hybrid inflation [10].

The connection to hybrid inflation can be made much more explicit within string theory. The amplitude for brane-antibrane annihilation develops a negative mode when the distance between a brane and antibrane shrinks to M_s^{-1}. This signifies the opening up of new channels for annihilation. In particular, it signals the appearance of tachyonic open strings connecting the branes to the antibranes [8, 15]. The mass of the tachyon is a function of the inter-brane distance and is given by [7] $M_T^2 = y^2/(2\pi\alpha')^2 - 2/\alpha'$. Here, y is the inter-brane separation and $\alpha' \sim M_s^{-2}$. Clearly this mass is tachyonic for small y, changing sign when $M_s y \lesssim 1$.

Although the complete form of the potential as a function of both y and T is not known, the limiting forms for $T = 0$ and $y = 0$ are. For large y, the large positive M_T^2 localizes T at zero, leaving only the relatively flat function of y on which the analysis of the previous sections was based. More recent work has also given the potential for $y = 0$, in which case the two-derivative truncation of the exact tachyon action (in the dimensionally-reduced form) is conjectured to be [16]

$$S_T = -M_s^2 M_p^2 \int \mathrm{d}^4 x e^{-|T|^2} \left[1 + \kappa_1 \left(\frac{2 + \kappa_2 |T|^2}{M_s^2} \right) |\partial T|^2 \right]. \tag{3.1}$$

Here κ_1 and κ_2 are calculable, dimensionless $\mathcal{O}(1)$ numbers, whose values need not concern us further. Only one feature of eq. (3.1) is important for our present purposes: the qualitative shape of the potential well, which is a Mexican hat potential once expressed in terms of the canonically normalized scalar field.

Thus, the tachyon potential behaves as follows: while the branes are still far apart, the potential is parabolic and the tachyon field is localized at the bottom of the potential, $T = 0$. Once the distance shrinks to less than M_s^{-1}, the potential develops a double well shape and the tachyon starts to roll downwards until it reaches one of the degenerate minima of the potential. If the winding number of the tachyon is zero, then the brane and antibrane will annihilate. According to a famous conjecture of Sen [17], the height of the tachyon potential is equal to the brane tension and the end product of brane-antibrane annihilation is the closed string vacuum. On the other hand, if the tachyon has non-trivial winding, then stable lower dimensional branes will be formed after it condenses.

Many types of lower dimensional branes may be formed [18]. For example, on the worldvolume of the Dp-$\overline{\mathrm{D}p}$ pair there are kink-like solutions to the tachyon field equations that correspond to unstable non-BPS D$(p-1)$-branes. These D-branes in turn contain a real tachyon field which induces them to decay into stable BPS D$(p-2)$-branes. Equivalently, these D$(p-2)$-branes correspond to vortex configurations of the complex tachyon that lives on the worldvolume of the Dp-$\overline{\mathrm{D}p}$. By considering more than just a single brane-antibrane pair it is actually possible to describe a system of D$(p-2)$-branes. Furthermore, by looking at anti-kink solitons on the unstable D$(p-1)$-branes one obtains $\overline{\mathrm{D}(p-2)}$-branes. This then allows the formation of $(p-2)$-dimensional brane-antibrane systems that can lead, in turn, to D$(p-4)$-branes and so on. Starting from a system of D9-$\overline{\mathrm{D}9}$-branes we can thus generate all possible D-brane configurations in type-I and type-IIB string theory. A similar story holds for type IIA, with the non-BPS D9-brane as the starting point. This is the basis of the K-theory classification of D-branes [19].

The tachyon potential above may have important applications for cosmology. We have checked that in its simplest setting it does not produce slow-roll inflation by itself, as might be expected given the absence of small parameters. Its role begins once inflation ends, when the inter-brane separation becomes of order $1/M_s$. Once this point is reached, the tachyon mass changes sign and T is free to quickly roll away from zero.

In a field theory example [20], this roll along the tachyonic direction has been shown to be very rapid, with the initial potential energy being quickly converted into gradient energy for the tachyon field, and may also result in topologically non-trivial configurations. We regard the quickness of this transfer of energy into gradients of T as a sign that a significant number of lower-dimensional daughter branes are produced in each brane-antibrane collision. If most of these daughter branes very quickly find nearby antibranes to annihilate, liberating the gradient energy to heat, then we are left with a post-collision world involving some remnant lower-dimensional branes, and with much of the collision energy converted into exciting both bulk modes and the internal degrees of freedom on the daughter branes. A crude estimate of the reheat temperature is given by equating $T_{\rm rh}^4$ with the initial potential energy density, leading to

$$T_{\rm rh} \sim \left[\frac{2\alpha e^{\phi}}{(M_s r_\perp)^{d_\perp}} \right]^{1/4} \sqrt{M_s M_p} = \left[\frac{2\alpha e^{\phi}}{(M_p r_\perp)^{d_\perp}} \left(\frac{M_p}{M_s} \right)^{d_\perp - 2} \right]^{1/4} M_p \,. \tag{3.2}$$

For instance if $d = 6$ and $p = 5$ and we take $M_p r_\perp \sim 10^6$ (as is suggested by the amplitude of primordial density fluctuations) and $\alpha e^{\phi} \sim 10^{-2}$ and $M_p/M_s \sim 10^3$ then we find $T_{\rm rh} \sim 10^{13}\,$GeV.

Although a careful analysis of this collision process is beyond the scope of this paper, we remark that the picture has several attractive features:

- Because the collisions can have daughter branes, the reheating process can be expected to efficiently excite the brane modes as well as those of the bulk, as is required if most of the observable universe now lives on one of these branes.

- Since the collision process produces daughter branes of lower dimension, successive episodes of brane collisions will move the brane population away from the higher-dimensional branes. Furthermore, since lower dimensional branes have more difficulty finding one another than do higher dimensional ones, this annihilation process will eventually be self-limiting. The branes found in the late universe will have a dimension which is small enough to preclude their having found anti-partners to annihilate during the history of the universe.

- There is a simple estimate — following an argument similar to that of ref. [21] — of what dimension a brane must have in order to plausibly fail to find partners with which to annihilate, and so to dominantly survive into the present. The basic observation is that p-branes — whose world-sheets are $(p+1)$-dimensional — generically have difficulty missing one another if they are moving within a spacetime whose dimension is less than or equal to $2(p+1)$. For 10-dimensional spacetime, passing p-branes with $p \geq 4$ will generically collide, while those with $p \leq 3$ will generically miss. We see then that if we start with a 10-dimensional spacetime only branes of dimensions greater or equal than 3 have finite probability of colliding with each other. If we start with a system of D9-$\overline{\rm D9}$-branes they will give rise to pairs of D7-$\overline{\rm D7}$ which can still annihilate each other to produce pairs of D5-$\overline{\text{D5}}$ and then D3-$\overline{\text{D3}}$. However, assuming the 3-branes can resolve the compact dimensions, these 3-branes will have negligible probability of finding each other. This implies that the cascade naturally stops there.

- Furthermore, as is illustrated in figure 2, brane collisions can plausibly act to align the surviving branes parallel to the noncompact dimensions of space. This is because collisions between branes wrapped around the compact dimensions and branes parallel to the noncompact dimensions tend to leave daughter branes which are both parallel to the noncompact dimensions and wrapped around the compact ones. Furthermore, the residual wrapping about the extra dimensions can be removed through the collisions of branes with opposite winding number, such as discussed in ref. [21].

- Within the cascade picture there are many branes colliding, but only some will cause inflation: those which move slowly enough to give sufficient e-foldings. Once inflation happens, the inflated area dominates the volume of the universe. For branes wrapped about a common direction there is an energetic cost to be not straight and parallel with one another, so given enough time branes will tend to align and straighten. Now one of the reasons it is hard to find real string potentials which inflate is that, in general, adjustments of brane energy (like removal of wiggles, or alignment of branes, or their relative rotations) take place on string timescales, making them too fast for inflation. But since inflation preferentially happens for the slowest collisions, it is plausible that these other processes will have had time to already have taken place, tending to make the branes align, be straight, etc. So we believe that the simplifying assumptions we make for our initial conditions are plausible for those relatively rare collisions which are of most interest for inflation. This plausibility needs to be further investigated within our cascade scenario.

It is quite intriguing that a D3-brane universe with branes aligned with the large dimensions of space seems to be selected in such a natural way. Notice that our arguments are not identical to the one given in [21], but complements it in an interesting way. Their argument was based on the fact that winding modes tend to compactify the space dimensions where they are winding around and then prefer only three spatial dimensions to become large (in the original counting, winding strings with $p = 1$ cannot annihilate in dimensions larger than $4 = 2(p + 1)$), therefore splitting naturally the ten dimensional spacetime into four large and six small dimensions. Our argument would complement this by stating that in this spacetime there will

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be stable D3-brane-worlds where we may live.[8] Notice that in some sense the assumptions of our argument are less constrained than that in the discussion of [21]. For their argument to hold they have to assume all the spatial dimensions to be toroidal, in our case the argument applies independent of the properties of the full ten-dimensional spacetime.[9]

We see that a cosmology of cascading brane annihilations may explain our finding ourselves on a 3-brane parallel to the large dimensions, by relating it to the prediction that spacetime is ultimately 10- or 11-dimensional.[10] Although the arguments we propose along these lines are preliminary and qualitative, we believe they have sufficient promise to deserve more careful further investigation.

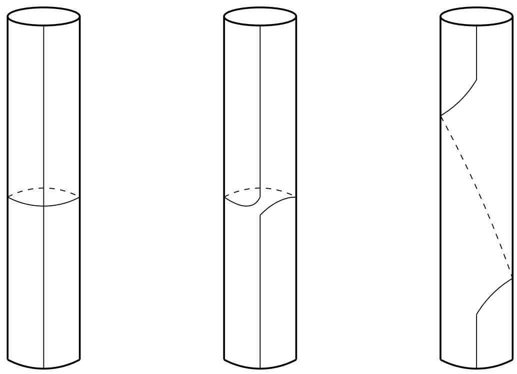

Figure 2: The figure on the left shows two branes along orthogonal directions of a long cylinder. In the middle, the branes reconnect. The final configuration, shown on the right, consists of a single brane winding once around the compact direction.

Let us finish by recapitulating the results of this article. In the first part of our paper we asked the question: Is the brane/antibrane potential leading naturally to inflation? Our answer was no. We then asked the question: Is inflation possible at all in this system? and the answer to this is yes, but only in special configurations when the potential happens to be flat enough. This is the first time that inflation is achieved from a computable potential in string theory. Furthermore the, also computable, stringy tachyon potential

[8]More specifically, one can imagine the following scenario: Branes of all types are initially there, and all dimensions are compact but small. The large dimension branes annihilate out instantly, leading a population of branes which dominantly involves 3-branes or lower. The winding of these various branes keeps any dimensions from growing. Then the Brandenberger-Vafa mechanism kicks in, making 4 dimensions large and 6 small, with no windings about the large spatial directions. Energetic arguments can prevent these small dimensions from growing or collapsing. After this, we have the particular collision which causes inflation of the large 4 dimensions, as studied in the previous section.

[9]Actually, the D3-brane-world is preferred if we start with type IIB or type I strings for which the starting point in the cascade are pairs of D9-$\overline{\text{D9}}$-branes. In type IIA, the starting point would be non-BPS D9-branes which then would give rise to even dimensionality BPS D-branes and the cascade would stop with 4-branes.

[10]On the remaining 3-brane we may have remnant cosmic strings and monopoles which could overclose the Universe. In order to properly pose and solve this question a more detailed knowledge of the final state distribution of branes after annihilation is required, including a good understanding of what the branching fraction for the production of daughter branes (as opposed to no branes) given a brane-antibrane collision. Also an understanding of how these daughters are distributed given the geometry of the initial brane/antibrane collision. These considerations are beyond the scope of our paper (see the second article of reference [21] for a related discussion).

provides the natural way to exit inflation in an elegant realization of hybrid inflation in terms of string theoretical fields. The fact that at the minimum of the tachyon potential supersymmetry is recovered guarantees the vanishing of the vacuum energy at that point without any fine-tunning.

Our proposal of the cascade mechanism, obtained from the tachyon potential, then addresses the issue of how to accomplish inflation naturally if it is so hard to achieve. If only two branes exist, it would be fine-tuned to expect them to collide in the way which inflates. But in the cascade picture the universe starts in a generic state in which many branes are present and colliding, and so the inflationary collisions are rare, but happen. Furthermore we can trace back the origin of the branes gas to overlapping, static, 9-brane/anti 9-brane systems as providing the initial conditions for the multi-brane Universe. We believe our scenario has many interesting features deserving further investigation.

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