

Hybrid Inflation and Brane-Antibrane System

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Abstract

We study a string theory inspired model for hybrid inflation in the context of a brane-antibrane system partially compactified on a compact submanifold of (a caricature of) a Calabi-Yau manifold. The interbrane distance acts as the inflaton, whereas the end of the inflationary epoch is brought about by the rapid rolling of the tachyon. The number of e-foldings is sufficiently large and is controlled by the initial conditions. The slow roll parameters, however, are essentially determined by the geometry and have little parametric dependence. Primordial density fluctuations can be made consistent with current data at the cost of reducing the string scale.

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1 Introduction

The knowledge of the history of the observable Universe before the epoch of nucleosynthesis is rather limited. However, it is widely believed that an early era of cosmological inflation existed [1]. The dynamics of the Universe was dominated by a homogeneous scalar field, called the inflaton field, during this era. The potential for the inflaton field accounted for almost all of the energy density of the Universe. The latter decreased with time as the scalar field rolled slowly down the slope of the potential. Apart from explaining the observed large scale homogeneity and isotropy of the Universe, the inflationary paradigm explains structure formation as having been seeded by tiny primordial density fluctuations. String theory, being a consistent theory of quantum gravity, ought to play a crucial role in providing a viable theory of inflation. An interesting direction that has been followed during the last couple of years is the study of inflation in the context of annihilation of extended objects called Dirichlet branes. There are unstable configurations of these branes in which various scalar fields arise in the low energy effective actions. Among these are the transverse scalars governing the motion of the branes, the ‘tachyon’ field corresponding to the instability of the configuration and moduli fields describing the background geometry. One or more of these could play important roles in inflation. An incomplete list of references is [2, 3, 4, 5, 6, 7, 8, 9, 10, 11, 12, 13]; see also [14] for a recent review.

In this paper we study the annihilation of a brane-antibrane system as a model of hybrid inflation [15, 1] inspired by string theory. The basic idea was first proposed in Ref. [3] (see also [4]). In these models, interbrane separation acts as the inflaton field, and the one that causes exit from the inflationary universe is the tachyon field. The latter is a scalar field corresponding to the lowest level excitation of the open string which connects the brane and the antibrane. It becomes tachyonic when the interbrane

separation reaches the string scale. We consider a pair of Dirichlet six and anti-sixbrane in (a caricature of) a Calabi-Yau space. Three of the worldvolume directions of the branes are wrapped on a three dimensional torus T^3 . The branes fill the spacetime we live in, but are separated in an internal three space, which we take to be a three dimensional sphere S^3 or a real projective three-space RP^3 . The inflaton is the four dimensional scalar field that corresponds to the separation between the branes. There is an attractive potential between the brane and the anti-brane, which, at large distances, is given by the solution of the Poisson equation on the transverse space S^3 (or RP^3). This potential drives inflation. However, when the branes come close enough, an instability develops and the rolling of this tachyonic scalar field causes the universe to exit from the inflationary phase. We find that it is possible to obtain enough expansion, although the density perturbations tend to be slightly higher, with reasonable input parameters (string coupling, string length and sizes of the internal space). We emphasize that ours is *not* a ‘brane world model’. In our model, the universe, at the end of inflation and the subsequent tachyon condensation, settles to a closed string vacuum determined by the Calabi-Yau manifold (possibly with fluxes [16], so that we have minimal supersymmetry).

It has been suggested that this annihilation scenario could be modified by starting with different numbers of branes and antibranes, so that one finally ends up with one (or more) brane(s) on which we live [3]. However, this would require one to start with a net brane charge on the compact transverse space. This is disallowed by the requirement charge neutrality on a compact space.

The plan of the paper is as follows. In the next section, we review some aspects of brane-antibrane geometry and dynamics [17]. In Sec. 3, we couple this system to gravity and derive the equations of motion. These are then analyzed numerically in Sec. 4. In the final section, we end with some comments.

2 Dynamics of brane-antibrane & spacetime geometry

Type II string theory has Dirichlet branes of all dimensions. More specifically, a stack of parallel Dp -branes for even (respectively odd) values of p are supersymmetric in type IIA (IIB) in flat ten dimensional spacetime. These D-branes satisfy BPS condition and break half of the supersymmetries². An antibrane is oriented opposite to a brane, and breaks the other half. Together, a parallel brane-antibrane pair (or a stack of them) give rise to a non-supersymmetric configuration. There is an attractive force between branes and antibranes. This is due to the exchange of massless graviton, dilaton and Ramond-

²Type IIA string has non-supersymmetric, *i.e.*, non-BPS branes for all odd values of p , and conversely in IIB.

Ramond (RR) forms³ (that are responsible for the charge of the branes), as well as the infinite tower of stringy fields; all of which are closed string modes.

Let us now focus our attention on a brane-antibrane pair, which we will simply refer to as branes in the following. When the distance r between the branes is large (in units of the string length scale $\sqrt{2\pi\alpha'}$), this force can be calculated from an effective field theory, and is of Coulomb type $\sim r^{p-8}$, where r is the radial coordinate transverse to the branes. This results in a potential for the open string scalar fields Y^i , ($i = 1, \dots, 9 - p$, $r = 2\pi\alpha'|Y|$) on the $(p + 1)$ -dimensional worldvolume of the branes.

Below a critical value $\pi\sqrt{2\alpha'}$ of interbrane separation, an instability develops. This manifests itself by making the lowest scalar excitation T on the open string (connecting the branes) tachyonic. (At the critical separation, the energy due to the stretching of the string is just enough to balance the negative zero point energy.) This instability is captured by a potential $V(T, Y)$. Since both T and Y^i are excitations of the open string, their interaction is at the tree level (classical) and dominates over the potential term involving the Y s. This is because the latter arises from exchange of gravitons etc. in the loops. Moreover, at separations comparable to the string scale, all the massive modes of the string need to be taken into account, and the simple analysis based on the effective field theory breaks down.

The spacetime geometry we are interested in is not the ten dimensional flat space, but a product of (3+1)-dimensional flat spacetime with a six dimensional internal space \mathcal{M} . In the absence of any brane, type II theory has $\mathcal{N} = 2$ supersymmetry in four dimensions when \mathcal{M} is a manifold of SU(3) holonomy, a Calabi-Yau space. (In [3] the internal manifold is taken to be a torus, which preserves all the supersymmetries. Subsequently, Ref. [5] considered an orbifold of torus which is a singular limit of a Calabi-Yau space. However, in the latter, interbrane separation is not the inflaton.)

There are a large number of moduli fields in any of these compactification schemes. However, recent developments [16] show that almost all these moduli can be frozen by turning on suitable fluxes in \mathcal{M} , and/or taking an ‘orientifold quotient’. What is more, the resulting theory has minimal, *i.e.*, $\mathcal{N} = 1$ supersymmetry. Be that as it may, a detailed model is beyond the scope of the present work. We work with the assumption that some mechanism is at place to freeze the unwanted fields, so that the only closed string fields are the graviton and the RR $(p + 1)$ -form.

We will further assume, following the picture advocated in Ref. [18], that the Calabi-Yau space is a three-torus T^3 fibration over a three dimensional base space \mathcal{B} . Among possible base spaces are those that could topologically be the three-sphere S^3 , or the projective space RP^3 , especially if an orientifold quotient is used. In other words, as far as the topology of \mathcal{M} is concerned, it can locally be written as a product $\mathcal{M} = T^3 \times S^3$ (though at special points of \mathcal{B} , the fibre T^3 may degenerate). A single D6-brane which fills our spacetime, preserves half of the supersymmetries if the compact part of

³For BPS, *i.e.*, supersymmetric brane configurations, the attractive force due to graviton and dilaton is cancelled by the repulsive force due to RR gauge fields.

its worldvolume is along T^3 [19]. An anti-brane, as usual, preserves the other half and together they break all the supersymmetries. This configuration is a non-supersymmetric excitation over the supersymmetric vacuum determined by \mathcal{M} . From now on we will work with this system of D6- $\bar{D}6$ branes. As far as the base \mathcal{B} is concerned, these branes are zero dimensional and sit at specific points. A realistic calculation would require the knowledge of explicit Calabi-Yau metrics, which unfortunately are not available. We will, therefore, consider a crude approximation of a Calabi-Yau manifold by treating \mathcal{M} to be $T^3 \times \mathcal{B}$. Moreover, we will put a flat metric on T^3 and the standard round metrics on S^3 and RP^3 , and neglect the back reaction of the branes on the metric of these spaces.

The potential energy of this system of branes comes from two sources. First, there is a contribution from the tension of the 6-branes wrapping T^3 . This is

$$\mathcal{E}_0 = 2\mathcal{T}_6 V_{\parallel},$$

where \mathcal{T}_6 is the 6-brane tension and V_{\parallel} is the volume of T^3 . Secondly, we have the energy due to the Newtonian and Coulombic interactions. If we assume that the branes are far apart on \mathcal{B} , this can be obtained by solving the Poisson equation. The strength of this interaction is determined by $G_{N(\mathcal{B})} = (2\pi g_s^2)(2\pi\alpha')^4/8V_{\parallel}$, and the interaction energy is

$$\mathcal{E}_{\text{int}} = 2\mathcal{T}_6 V_{\parallel} \frac{g_s \sqrt{2\pi\alpha'}}{8\sqrt{2\pi}} G(r) \equiv 2\mathcal{T}_6 V_{\parallel} \Phi(r), \quad (2.1)$$

where, r is the geodesic distance between the branes and

$$G(r) = \begin{cases} [(r - \pi R_S) \cot\left(\frac{r}{R_S}\right) - R_S] / 4\pi^2 R_S^2, & \text{for } S^3 \\ (2r - \pi R_P) \cot\left(\frac{r}{R_P}\right) / 4\pi^2 R_P^2, & \text{for } RP^3, \end{cases} \quad (2.2)$$

is the solution of Poisson equation obtained by following the general procedure outlined in [20]. As mentioned earlier, we have assumed the standard metric and ignored back-reaction due to the branes.

The effective field theory in (3+1)-dimensional spacetime has three scalar fields corresponding to the transverse motion on \mathcal{B} . The second contribution to the potential energy described above gives rise to a potential term $2\mathcal{T}_6 V_{\parallel} \Phi(Y)$ for the field Y which corresponds to the interbrane separation in \mathcal{B} . To get $\Phi(Y)$, one substitutes $r = 2\pi\alpha'Y$ in eqns.(2.1) and (2.2). Notice that, while the interaction term \mathcal{E}_{int} is due to (6+1)-dimensional gravity, the size of the internal space being small, its effect is to provide a potential term for the field Y in (3+1) dimensions.

Incorporating the kinetic energy of Y , which we recall is of the Born-Infeld form [21], we can write the lagrangian for the field Y and the ‘tachyon’ field T as

$$\mathcal{L}_{\text{matter}} = -2\mathcal{T}_6 V_{\parallel} \left(V(T, Y) \sqrt{1 - (2\pi\alpha')^2 (2f_c \partial_{\mu} T \partial^{\mu} \bar{T} + \partial_{\mu} Y \partial^{\mu} Y)} + \Phi(Y) \right), \quad (2.3)$$

where, $f_c = 2 \ln 2/\pi$ [7], and $V(T, Y)$ is a potential for the tachyon field alluded to earlier. An expression for this potential can be proposed by a straightforward generalization of the result derived in boundary string field theory (BSFT) [22]

$$V(T, Y) = \exp \left(2\pi\alpha' \left(Y^2 - \frac{1}{2\alpha'} \right) |T|^2 \right). \quad (2.4)$$

This form of the potential can be motivated in the following way. When the interbrane separation is zero, this potential matches with that proposed by boundary string field theory. For small values of tachyon, by expanding the exponential it is easy to see that it correctly produces the tachyon mass formula.

The potential (2.4), however, does not give the expected asymptotic behaviour for the tachyon [23, 24]. A better choice follows from the recent proposal of Ref. [13, 25]

$$V(T, Y) = \sec \left(2\pi\alpha' |T| \sqrt{Y^2 - \frac{1}{2\alpha'}} \right). \quad (2.5)$$

Both the functions have the property that for $Y^2 > 1/2\alpha'$, the mass-square of the field T is positive, indeed the potential is extremely steep for large separation of branes. On the other hand, it becomes tachyonic below the critical separation. We find that our conclusions are largely independent of the two possible choices.

3 Coupling to gravity

We will now couple the above field theory of scalars T and Y to (3+1)-dimensional gravity. We will assume a Friedmann-Robertson-Walker metric,

$$ds^2 = -dt^2 + a^2(t) (dx^2 + dy^2 + dz^2),$$

(with zero spatial curvature for simplicity). Since we are primarily interested in the time evolution of the fields, let us further assume $T = T(t)$ and $Y = Y(t)$, *i.e.*, both the tachyon T and the scalar Y are functions only of time.

The lagrangian that describes the dynamics of this universe is

$$\mathcal{L} = \frac{3}{\kappa^2} (a\dot{a}^2 + a^2\ddot{a}) - 2\mathcal{T}_6 V_{\parallel} a^3 \left(V(T, Y) \sqrt{1 - (2\pi\alpha')^2 (2f_c |\dot{T}|^2 + \dot{Y}^2)} + \Phi(Y) \right). \quad (3.1)$$

The four dimensional gravitational constant κ^2 is obtained in terms of the ten dimensional coupling constant $\kappa_{(10)}^2$ through

$$\frac{1}{\kappa^2} = \frac{V_{\parallel} V_{\perp}}{\kappa_{(10)}^2} = \frac{2\pi^2 \hat{R}_T^3 \hat{R}_P^3}{g_s^2 (2\pi\alpha')} = 2\mathcal{T}_6 V_{\parallel} \frac{\pi^{3/2} \hat{R}_P^3}{\sqrt{2} g_s} (2\pi\alpha') \quad \text{for } RP^3, \quad (3.2)$$

where, $\hat{R}_P = R_P/\sqrt{2\pi\alpha'}$ is the radius measured in string units. (The value of κ^2 for S^3 base space is twice that of RP^3 .) The ratio of the string scale to the four dimensional Planck scale that follows from the above is $(m_s/m_P)^2 = g_s^2/8\pi^2\hat{R}_T^3\hat{R}_P^3$.

The equations of motion are easily derived:

$$\frac{2}{3}\dot{H} + H^2 = \frac{g_s\sqrt{2}}{3\pi^{3/2}\hat{R}_P^3(2\pi\alpha')} \left(V\sqrt{\mathcal{B}_I} + \Phi \right), \quad \text{for } RP^3, \quad (3.3)$$

$$(2\pi\alpha')^2\ddot{Y} = -\mathcal{B}_I \left(\frac{1}{V} \frac{dV}{dY} + (2\pi\alpha')^2 3H\dot{Y} + \frac{1 - (2\pi\alpha')^2\dot{Y}^2}{\sqrt{\mathcal{B}_I}V} \frac{d\Phi}{dY} \right), \quad (3.4)$$

$$2f_c\ddot{T} = -\mathcal{B}_I \left(\frac{1}{(2\pi\alpha')^2 V} \frac{dV}{dT} + 6Hf_c\dot{T} - \frac{2f_c\dot{T}\dot{Y}}{\sqrt{\mathcal{B}_I}V} \frac{d\Phi}{dY} \right), \quad (3.5)$$

where, $H = \dot{a}/a$ and $\mathcal{B}_I = 1 - (2\pi\alpha')^2(2f_c|\dot{T}|^2 + \dot{Y}^2)$. In addition, we have

$$H^2 = \frac{1}{3}\kappa^2\rho = \frac{2}{3}\kappa^2\mathcal{T}_6V_{\parallel} \left(V/\sqrt{\mathcal{B}_I} + \Phi \right). \quad (3.6)$$

As a matter of fact, the above equations simplify in practice. At early times, when the separation between the branes is large, the steep potential V (eqn.(2.4) or (2.5)) keeps T pinned at $T = 0$, thereby yielding $V = 1$. Once the tachyon is excited, the potential Φ can be neglected. (Recall that it is due to one loop effect, while the interaction between tachyon and Y is at tree level.) In our numerical computation, we take care of this by freezing Φ at the point where the field T is excited.

The overall potential seen by the field Y at the early stages of evolution, and indeed until the excitation of the scalar T , is

$$\mathcal{V}_{\text{eff}}(Y) = 2\mathcal{T}_6V_{\parallel} (1 + \Phi(Y)) = 2\mathcal{T}_6V_{\parallel} \left(1 + \frac{g_s\sqrt{2\pi\alpha'}}{8\sqrt{2\pi}} G(2\pi\alpha'Y) \right). \quad (3.7)$$

For large brane separations comparable to the radius of \mathcal{B} , this potential goes as $\mathcal{V}_{\text{eff}} \sim 1 - A(Y/\hat{R})^2$. Although this is not flat, we will see that slow roll conditions are satisfied. Note that the behaviour of \mathcal{V}_{eff} is different from that of [3], where it is of the form $a - bx^4$.

4 Time evolution and Y-driven inflation

The equations of motion (3.3)–(3.6) are unfortunately not amenable to analytic solution. We therefore proceed to solve them numerically. It has already been noticed in a similar set up in Ref. [3] that, in order to get sufficient inflation, the branes should initially be as far as apart as possible in the compact space. For the transverse space S^3 or RP^3 , this means that they must be close to a pair of antipodal points to begin with. To be

concrete, we will present most of our numerical results for RP^3 and comment on the (mostly quantitative) differences for the S^3 at the very end.

To begin with, we also assume that the branes start with zero initial velocity, so once they start moving, they approach each other head on. These are the initial conditions that we use for Y . The field T , at this point, is in an extremely steep potential, and therefore it is a very massive one. Fluctuations of T in this case would cost a lot of energy and within our framework it makes sense to set T to zero⁴.

As the branes approach each other, T starts becoming lighter and when the interbrane separation nears the critical value $Y \sim 1/\sqrt{2\alpha'}$, the dynamics of T becomes important. Close to this transition point, we mimic the fluctuation of T by transferring a small part of the kinetic energy of Y to T at

$$Y_c = c_Y/\sqrt{2\alpha'}, \quad (1.0 \lesssim c_Y \lesssim 1.5). \quad (4.1)$$

The fluctuation sets off an oscillation of T . Once the separation goes below the critical value, T becomes tachyonic and starts to roll down the potential $V(T, Y)$. We have already seen that the interaction between T and Y is at tree level. Therefore, once T is excited, it is a valid approximation to neglect the potential $\Phi(Y)$. In our solution, we take care of this by freezing the value of Φ to $\Phi(Y_c)$. Strictly speaking, for $Y \sim Y_c$, we are in a stringy regime, and the low energy effective field theory description really breaks down. However, once the tachyon starts rolling, its dynamics dominates. An effective field theory is once again a good description [26].

It might be argued that the parameter c_Y introduces an unnecessary element of arbitrariness, one that could, in principle, have been eliminated had we engaged in a full string theoretic calculation. This exercise being intractable, we find it instructive to examine, at this stage, the dependence of the time evolution of the fields T and Y on c_Y . For this purpose, and even for the rest of the numerical analysis, it is convenient to consider the normalized field θ_Y :

$$\theta_Y = \frac{\sqrt{2\pi\alpha'}Y}{\hat{R}_P} \equiv \frac{Y}{\hat{R}_P}, \quad 0 \leq \theta_Y \leq \frac{\pi}{2}. \quad (4.2)$$

As an examination of Figs. 1(a,b) shows, the subsequent evolution does not suffer any qualitative change as c_Y is varied over a reasonably large range. Quantitatively, the onset of the rapid roll down towards $\theta_Y = 0$ is delayed or brought forward. This, in turn, determines the instant at which T becomes tachyonic and rolls down the potential. As we shall see below, most of the inflation would have taken place before this transition point is reached. Therefore, the number of e-foldings is largely independent⁵ of c_Y . In the remainder of this analysis we shall use a representative value of c_Y , namely $c_Y = 1.2$.

⁴As we did for massive string states. In spite of this, it may be instructive to look at the fluctuations of T at the initial epoch, and we will return to this point later in this section.

⁵Of course, if c_Y were to be so large as to necessitate the freezing of $\Phi(Y)$ relatively close to the antipodal points, the above arguments clearly would not hold.

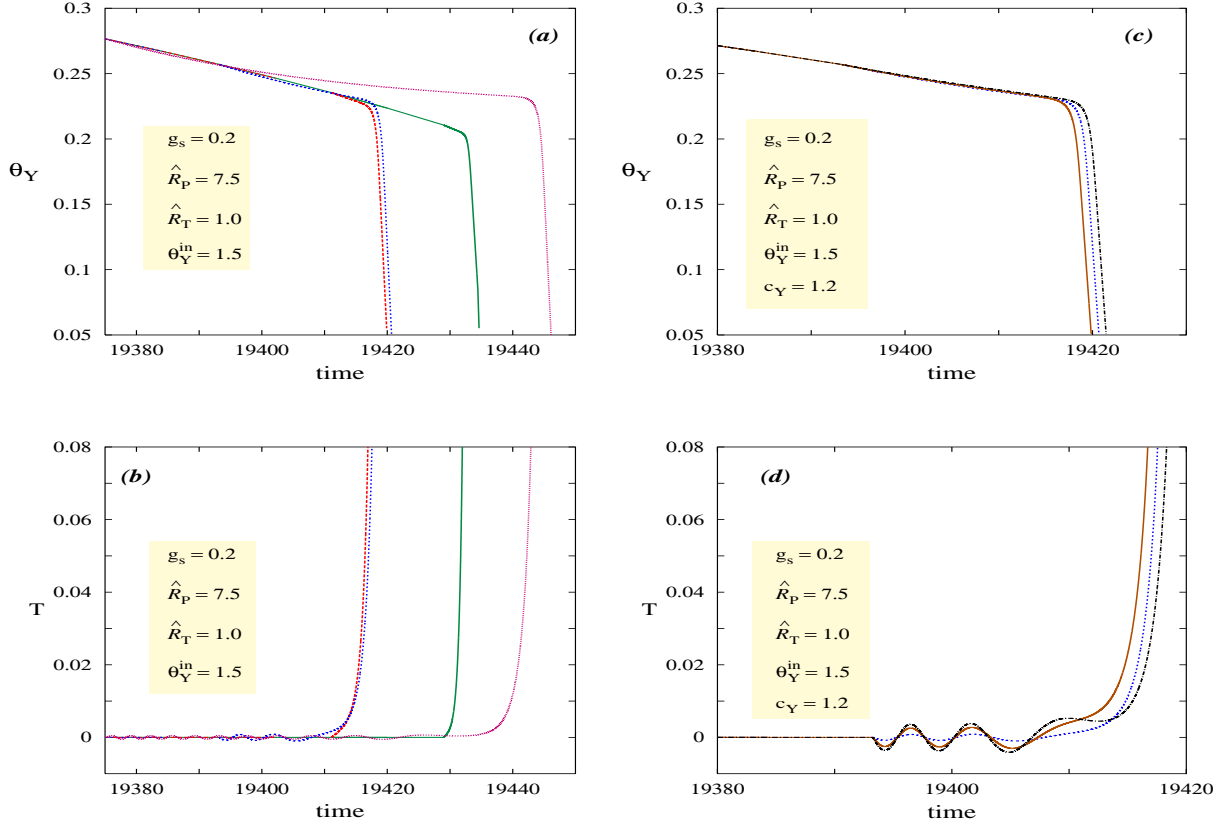


Figure 1: *The left panels depict the time evolution of the scalar fields θ_Y and T subsequent to the transfer of kinetic energy from the former to the latter, for different transfer locations. The curves range over $c_Y = 1$ to 1.5 (from right to left). In each case, 1% of the energy is transferred. The panels on the right concentrate on a single location ($c_Y = 1.2$) but for different amounts of energy transfer (1% to 20%).*

Another aspect that needs looking at is the dependence on the amount of kinetic energy that is transferred from Y to T . While the heuristic picture presented above suggests that this fraction should be a small one, in our investigations we have allowed for a fairly large variation. The results are depicted in Figs. 1(c,d). Once again, it is obvious that no qualitative change is brought about in the dynamics by transferring a large amount of energy to the would-be tachyon. Hence, for the remainder of this analysis we shall consider the case in which only 1% of the kinetic energy of Y is transferred to T , at the point $Y = c_Y/\sqrt{2\alpha'}$.

Let us now study, in detail, the complete time evolution of the system for a representative set of parameter values. As has already been pointed out, when the interbrane separation is large compared to the string length, the field T should not develop at all. In

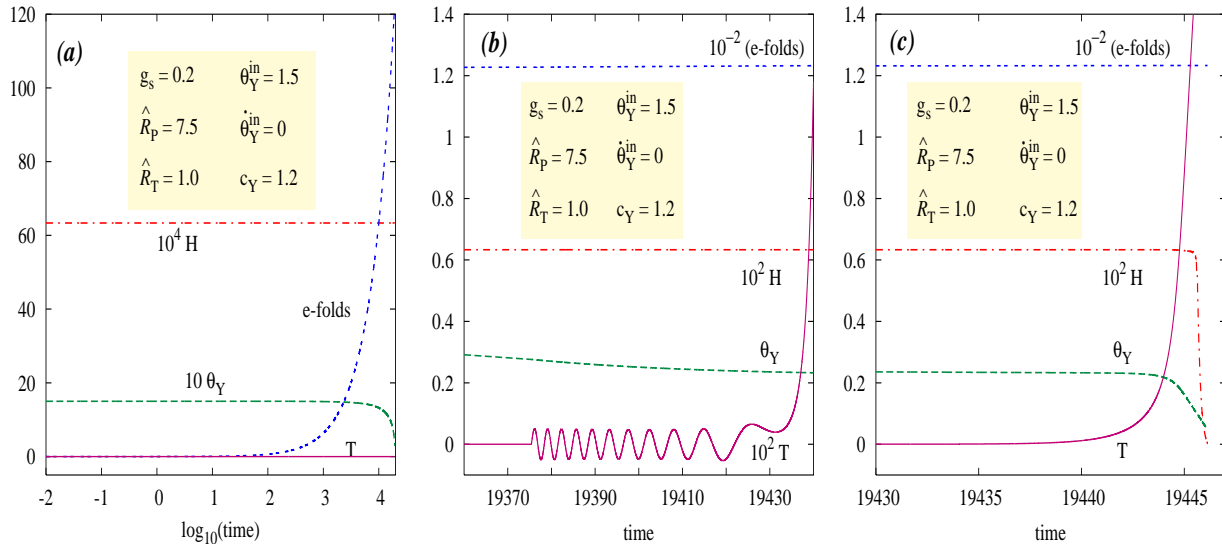


Figure 2: *The time variation of the fields T , θ_Y ($\equiv Y/\hat{R}_P$) and the Hubble parameter H . Also shown is the growth of the number of e-foldings with time. Panels (b) and (c) are zoomed-in views of tiny slices of (a).*

other words, at very early epochs, one should start with $T_{\text{init}} = 0$, $\dot{T}_{\text{init}} = 0$. The classical equations of motion (3.5) immediately dictate that the field should continue to be firmly pinned at $T = 0$. Consequently, as Y rolls down the potential $\Phi(Y)$, it exchanges energy only with gravity. With the effective potential $\mathcal{V}_{\text{eff}}(Y)$ being a slowly varying function of Y (at least for $Y \sim R_B$), one expects that the initial evolution of Y would be a very slow one. With the time evolution of H being governed by eqn.(3.6), this immediately translates to an even slower time variation of H . Both these expectations are borne out by Fig. 2(a). An immediate consequence is that the universe experiences an (approximately) exponential growth rate with the number of e-foldings, defined through

$$N_e(t) = \int_{t_0}^t dt H(t) ,$$

growing almost linearly with t .

Once the field T is excited, the dynamics undergoes a qualitative change. The field T , still with a positive mass-square, can now oscillate about its mean position (see Fig. 2(b)). With Y continuing its roll down towards $Y = 1/\sqrt{2\alpha'}$, the effective mass of T decreases with time, thereby increasing the oscillation time period. As Y falls below the critical value $1/\sqrt{2\alpha'}$, T becomes momentarily massless, then turns tachyonic. This is reflected by a brief slowdown of the time evolution of T , only to be followed by a rapid rolling down the potential. It is interesting to note that even with an evolving T , the Hubble parameter is almost constant in time (at least, in the initial phase). Only after T has

rolled down the potential sufficiently, does $H(t)$ start to suffer a (modest) decrease (see Fig. 2(c)).

4.1 Conditions for slow roll

Let us now discuss the conditions for slow roll. An investigation of equations (3.4) and (3.6) show that the slow roll parameters are quite like the standard ones with \mathcal{V}_{eff} in (3.7) as the potential. In the regime in which we are interested in their values, it is easy to check that the Born-Infeld form of the action does not make much of a difference. Moreover, for most of the inflationary epoch, the evolution of T can safely be neglected. The parameters that characterize slow roll, are

$$\varepsilon \equiv \frac{1}{4T_6 V_{\parallel} \kappa^2} \left(\frac{\mathcal{V}'_{\text{eff}}}{\mathcal{V}_{\text{eff}}} \right)^2, \quad \eta \equiv \frac{1}{2T_6 V_{\parallel} \kappa^2} \left(\frac{\mathcal{V}''_{\text{eff}}}{\mathcal{V}_{\text{eff}}} \right), \quad \xi \equiv \left(\frac{1}{2T_6 V_{\parallel} \kappa^2} \right)^2 \left(\frac{\mathcal{V}'_{\text{eff}} \mathcal{V}'''_{\text{eff}}}{\mathcal{V}_{\text{eff}}^2} \right), \quad (4.3)$$

where, the factors of $2T_6 V_{\parallel}$ can be attributed to non-canonical normalization of the field Y . Inflation is of slow roll type if $\varepsilon, |\eta|, |\xi| \ll 1$. For the case of the RP^3 these expressions read

$$\begin{aligned} \varepsilon &= \frac{g_s}{8(8\pi)^3 (2\pi)^{1/2} \hat{R}_P} \left[\frac{2 \cot \theta_Y - (2\theta_Y - \pi) \csc^2 \theta_Y}{1 + c_{RP} (2\theta_Y - \pi) \cot \theta_Y} \right]^2, \\ \eta &= \frac{1}{32\pi} \frac{\csc^2 \theta_Y (-2 + (2\theta_Y - \pi) \cot \theta_Y)}{1 + c_{RP} (2\theta_Y - \pi) \cot \theta_Y}, \\ \xi &= \frac{\csc^2 \theta_Y}{8(16\pi)^2} \frac{2 \cot \theta_Y - (2\theta_Y - \pi) \csc^2 \theta_Y}{(1 + c_{RP} (2\theta_Y - \pi) \cot \theta_Y)^2} \left[6 \cot \theta_Y - (2\theta_Y - \pi) (2 \cot^2 \theta_Y + \csc^2 \theta_Y) \right], \end{aligned} \quad (4.4)$$

where, $c_{RP} = g_s / (8(2\pi)^{5/2} \hat{R}_P)$, and, for convenience, we have set $2\pi\alpha' = 1$. It is worth pointing out that η and ξ are, for most part, independent of the parameters of the model⁶. Therefore, the slow roll conditions are largely unaffected, leaving us the freedom to vary the parameters.

It is easy to see that $\varepsilon \ll |\eta|$. As Fig. 3(a) shows, $|\eta| \ll 1$ as long as $\theta_Y \gtrsim 0.5$. As can be checked, $\theta_Y \simeq 0.24$ corresponds to the point where T becomes tachyonic and the expansion of the universe starts to slow down, *i.e.*, $H(t)$ starts decreasing with time. The spectral index

$$n_s \equiv 1 - 6\varepsilon + 2\eta,$$

where the right hand side is to be evaluated at approximately 60 e-folds before the end of inflation, is marginally below 1 and eminently consistent with the MAXIMA, BOOMERANG

⁶We thank Nima Arkani-Hamed for emphasizing this point.

and DASI observations [27]. Note that, unlike in canonical hybrid inflation models, we have $\eta < 0$, a consequence of the shape of the potential $\Phi(Y)$. The running of the spectral index

$$\frac{dn_s}{d \ln k} = \frac{2}{3} \left((n_s - 1)^2 - 4\eta^2 \right) + 2\xi \sim 2 \times 10^{-4},$$

has a sign opposite to that preferred by the WMAP data [28]. Note, however, that the extraction is beset with uncertainties, both observational and theoretical, and that our value is very well consistent at the 2σ level.

Yet another relevant quantity is the deceleration parameter q

$$q = -\frac{\dot{H}}{H^2} - 1, \quad (4.5)$$

which we have plotted in Fig. 3(b). That q is almost identically -1 for $\theta_Y \gtrsim 0.24$ is but a restatement of the exponential inflationary phase. On the other hand, $q > 0$ signals the end of the inflationary era and the beginning of the decelerating phase. For the present choice of parameters, this occurs at $\theta_Y \approx 0.2$, *i.e.*, a little after the tachyon develops. In summary, the slow roll expansion culminates in a tachyon driven fast roll one [9, 7] and very quickly evolves into the decelerating phase.

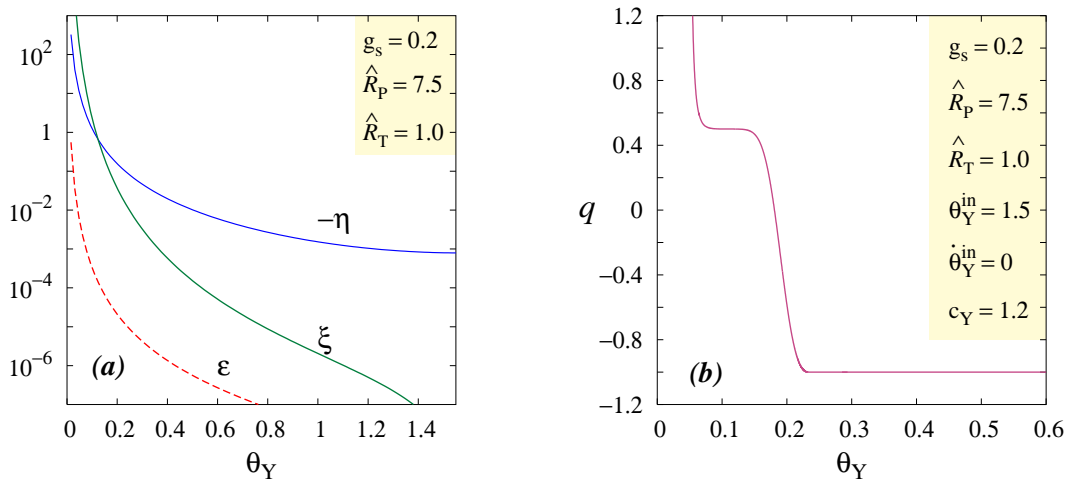


Figure 3: (a) The slow roll parameters ε and η as a function of the normalized field θ_Y . (b) The deceleration parameter as a function of θ_Y .

4.2 Dependence on parameters & Density perturbations

Until now, we have chosen to work only with a particular set of model parameters. It is important that we check for the robustness of our findings with variations in these. To

start with, let us consider the string coupling constant g_s and the radius of the projective space \hat{R}_P . As equations (3.2) and (3.6) imply, the Hubble parameter $H \propto \sqrt{g_s}$, and hence, all other things remaining the same, a smaller (respectively larger) value for the latter would imply a smaller (larger) value for the number of e-foldings. There is, of course, an opposite effect on account of the change in the size of the nontrivial term in the effective potential $\mathcal{V}_{\text{eff}}(Y)$ —see eqn(3.7)— as this determines the rate at which Y rolls down (and hence the length of the inflationary era). On convoluting the two effects, the resultant dependence is rather weak as is reflected in Fig. 4(a), where we display the variation in the number of e-foldings with \hat{R}_P for different values of g_s . It should also be noted that the slow roll parameter $\varepsilon \propto g_s$. However, since ε is so very small during most of the inflationary phase (see Fig. 3), reasonable changes in g_s would not result in any observable effect. The dependence on \hat{R}_P is more subtle though. Since $\kappa^2 \propto 1/\hat{R}_P^3$ (see eqn.(3.2)),

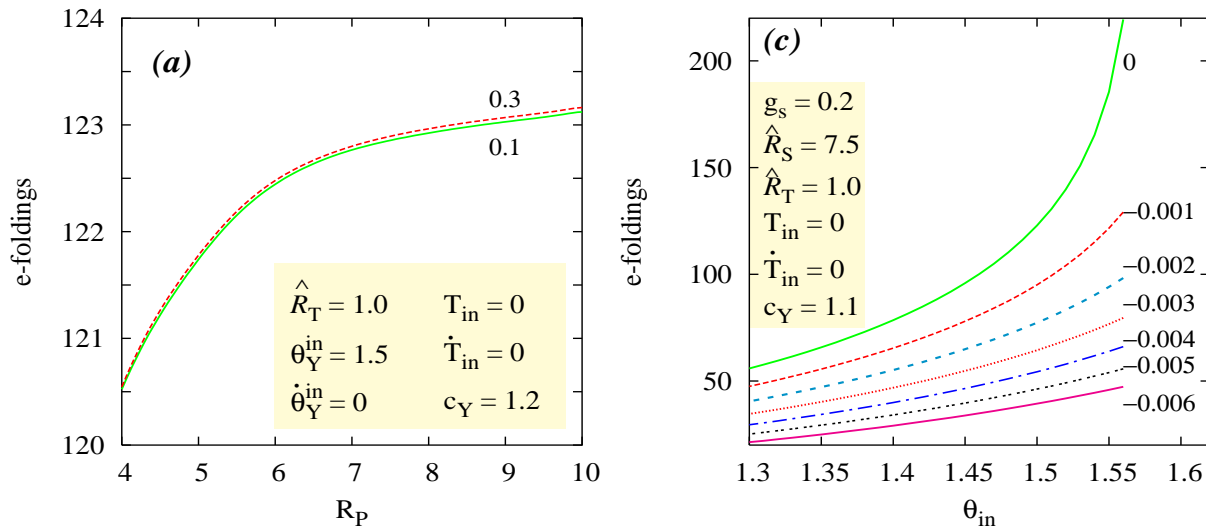


Figure 4: (a) *The dependence of the number of e-foldings on \hat{R}_P , the radius of the projective three-space. The two curves correspond to different values of the string coupling constant g_s .* (b) *The dependence of the number of e-foldings on the initial position of the brane, on RP^3 , with respect to the antibrane. The different curves correspond to different values of the initial inter-brane radial velocity.*

it would be tempting to conclude that $H(t) \propto 1/\hat{R}_P^{3/2}$ for most of the inflationary period and hence $N_e \propto 1/\hat{R}_P^{3/2}$ as well. While the first supposition is largely true, the second is not, as is clear from an inspection of Fig. 4(a). In a large measure, this is due to the fact that the distance to be traversed by the brane increases with \hat{R}_P . Furthermore, increasing \hat{R}_P suppresses the slope of the potential $\mathcal{V}_{\text{eff}}(Y)$ resulting in a slower roll in the initial phase. In effect then, the amount of inflation is nearly independent of \hat{R}_P , though the duration of the inflationary phase is not.

Is \hat{R}_P then unmeasurable? Fortunately, observations guide us here. As the fields roll to their minima, adiabatic density perturbations are caused. The amplitude of these perturbations can be estimated to be [1]

$$\begin{aligned}\delta_H &\simeq \frac{2T_6 V_{\parallel} \kappa^3}{5\pi} \frac{\mathcal{V}_{\text{eff}}^{3/2}(Y)}{|\mathcal{V}'_{\text{eff}}(Y)|} \\ &= \frac{2^{25/4} \sqrt{g_s}}{5\pi^{3/4} \hat{R}_P^{5/2}} \frac{(1 + c_{RP} (2\theta_Y - \pi) \cot \theta_Y)^{3/2}}{|2 \cot \theta_Y - (2\theta_Y - \pi) \csc^2 \theta_Y|}.\end{aligned}\tag{4.6}$$

It would immediately transpire that our result for δ_H (obtained with the default value for the radius \hat{R}_P) is three orders of magnitude larger than the COBE result of $\delta_H \sim 2 \times 10^{-5}$. This, in a sense, was to be expected. With $\delta_H \propto \sqrt{\kappa^4 V / \varepsilon}$ for a generic theory, even a moderately small value of ε would require that the appropriate value of the potential be much smaller than $1/\kappa^4$. With our theory being operative essentially at the Planck scale, this certainly is not true.

On the other hand, the COBE result could immediately be used to fix \hat{R}_P . A better value of δ_H is obtained if \hat{R}_P is scaled up by a factor of ~ 16 . This has an immediate consequence in that the string scale comes out to be $\sim 6 \times 10^{13}$ GeV (assuming $\hat{R}_T = 1$).

Increasing \hat{R}_P has secondary effects. For one, the excitation of T and, hence, the exit from inflation would now occur at even smaller values of θ_Y . Since the slow roll parameters are essentially determined by θ_Y , one might think that these grow significantly and thereby affect the spectral index inordinately. It can easily be checked though that this is not the case and that the differences (in either the value of n_s or its running) would be barely noticeable at the scale of COBE normalization.

Let us now discuss the dependence on the initial separation and the relative velocity of the branes on RP^3 , which have been kept fixed until now⁷. Placing them strictly at antipodal points with zero initial velocity would imply an equilibrium, though unstable, configuration. On the other hand, if the branes were to start far from the antipodal points, they would be feeling a relatively strong attractive force right from the beginning. Compounded with the fact that they would now have to traverse a smaller distance as well, it is easy to see that the number of e-foldings should reduce dramatically with the initial interbrane separation (see Fig. 4(b)). For example, the 122 e-foldings that we had obtained starting from $\theta_Y^{in} = 1.5$ increases to 220 if we start from $\theta_Y^{in} = 1.56$ and drops to only 70 for $\theta_Y^{in} = 1.37$. Of course, if the brane had, in addition, an initial radial velocity towards the antibrane, the reduction in the time of flight (and hence in N_e) is only helped (Fig. 4(b)). Note, however, that with a nonzero initial velocity, the dependence of N_e on θ_Y^{in} is weakened. This, again, is only to be expected.

The only parameter that remains to be discussed is \hat{R}_T . As our equations of motion would testify, the dynamics is entirely independent of \hat{R}_T . The only place it enters is in

⁷The Ref. [29] discusses the initial conditions for brane inflation models.

the definition of the four dimensional gravitational coupling constant in terms of the (*a priori* unknown) ten dimensional one.

4.3 Effect of initial fluctuations of T

Let us return to the issue of the fluctuations of the field T at early stages of evolution. We may try to mimic random (quantum) fluctuations by displacing it from its minimum or give it a non-zero velocity to start with. The consequent dynamics is portrayed in Fig. 5. While it may appear that the oscillation frequency is increasing rapidly with time, it is but an artifact of the choice of a logarithmic scale for time. In reality, the frequency *decreases* with time, as it should for a system losing energy. This is also reflected by the decreasing amplitude of the oscillation. As the coupling of T to Y is stronger than that

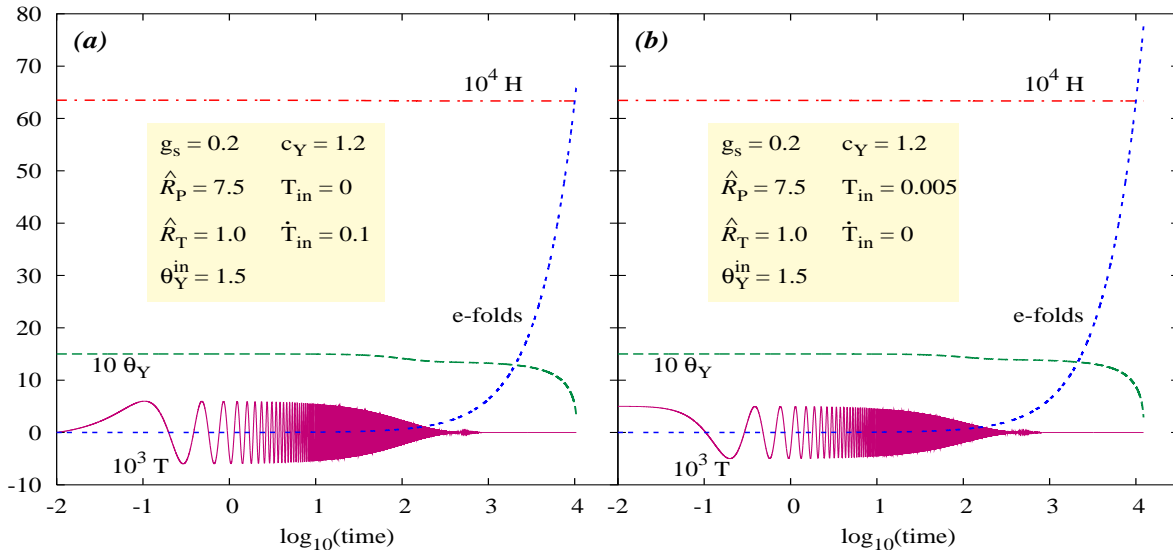


Figure 5: *The time variation of the fields T , θ_Y , Hubble parameter H and the growth of the number of e-foldings. The parameters are the same as in Fig. 2, but initial fluctuations of T are included. (a) T is given only a nonzero initial velocity; (b) T is given only a nonzero initial displacement.*

to the graviton, the energy that T loses is essentially gained by Y , in the form of kinetic energy. Consequently, Y rolls faster, whereas $H(t)$ remains virtually unaltered.

An immediate outcome of this faster rolling of Y is the drop in the time allowed for inflation. With $H(t)$ remaining the same, this translates to a rapid drop in the number of e-foldings as is evinced by Fig. 6. However, as mentioned earlier, it is inconsistent to

consider the fluctuations of the extremely heavy field T within the framework of a low energy effective theory. Moreover, if T is excited, there is no reason to exclude excitations of massive stringy modes, which for large separation are lighter than the tachyon. While a proper resolution of this has to emerge from the full string theory, in a purely field theoretic hybrid inflationary model, this effect may reduce the number of e-foldings.

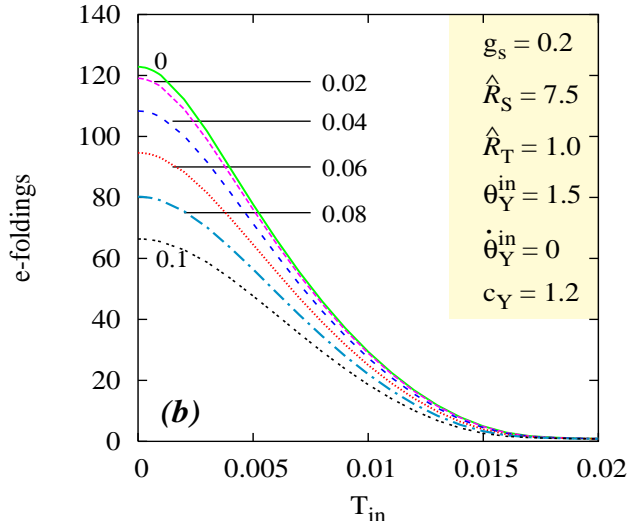


Figure 6: *The dependence of the number of e-foldings on the initial value of the T field. The different curves correspond to different values of the initial velocity \dot{T} .*

5 Discussion

We have studied an hybrid inflationary scenario in which a Dirichlet six- and antisix-brane, partially wrapped on a compact space, annihilate. The geometry we have worked with is Minkowski spacetime times a six dimensional compact space, which is a crude approximation to a Calabi-Yau manifold. We also assume that some mechanism (like flux) stabilizes unwanted closed string moduli and gives minimal supersymmetry. Ours is not a brane world model in the sense that after the process of brane annihilation, we are left with a purely closed string background determined by the Calabi-Yau space (with fluxes). In particular, all the standard model fields arise from the massless excitations of closed strings.

Most of the analysis in the paper is done for the case in which the space transverse to the branes is a projective three space RP^3 . A simpler choice would be the sphere S^3 . As a matter of fact, there is no qualitative difference between the two cases. Indeed most formulas are almost identical except for some numerical factors, although the potentials (2.2) that drive inflation differ in their details. It turns out that, for comparable choice

of parameters, one gets much more (almost an order of magnitude increase in) e-foldings than RP^3 . (Fig. 7 shows the time evolution of some of the relevant quantities.) Notice that this gives one considerably more freedom in varying the initial conditions (interbrane separation and relative velocity), still keeping the amount of inflation within acceptable limits.

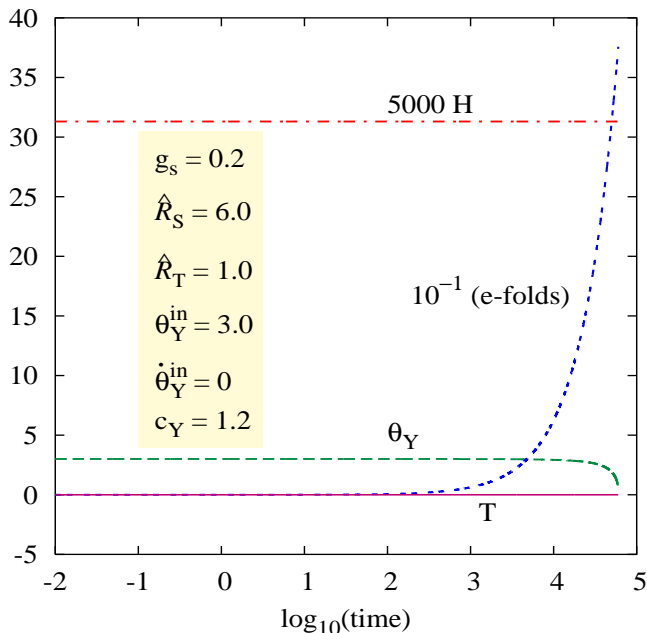


Figure 7: The time variation of the fields T , θ_Y ($\equiv Y/\hat{R}_S$) and the Hubble parameter H for the case of S^3 . Also shown is the growth of the number of e-foldings with time.

One would like to know how the matter and radiation fields take over at the end of inflation⁸. At first sight, it might seem that the runaway potential of the tachyon does not allow for reheating to excite the standard model fields. However, D-branes are excitations of the closed superstrings, hence when they annihilate the energy has to be converted into closed string modes. The details of this process are just beginning to be addressed [25, 30, 31]. However, the basic point is that D-branes are classical sources of closed strings, and D-branes with rolling tachyon are therefore time dependent classical sources. As in any field theory, coupling to time dependent sources leads to particle production. In the case of string theory, there are an infinite number of particles and it turns out that, in a homogeneous process of brane decay, heavy stringy modes of mass of order $1/g_s$ at zero spacetime momenta are preferentially excited. This is due to the fact that the tension of the D-brane is $\sim 1/g_s$. It is also consistent with the fact that these heavy closed string modes have the same thermodynamic properties as the ‘tachyon

⁸The issue of reheating in tachyonic inflation has been discussed in Ref. [11]. However, unlike ours, theirs is a brane world model.

matter' [23, 26]. However, it is not yet clear how and at what time scale these modes will decay into massless closed string states.

We have only considered a homogeneous decay of the complex tachyon. However, before the field T turns tachyonic, its potential flattens out allowing it to oscillate with large amplitudes. Therefore, at different points of space, T may roll down along different directions in field space. Topological defects may form in this process [32]. It will be interesting to study this in detail.

It is natural to wonder how the six-antisix brane pair come into being and start in the way they do. Namely wrapped on supersymmetric T^3 -cycles of a compact Calabi-Yau manifold at nearly antipodal points of a base manifold \mathcal{B} . Let us end by speculating on a seemingly natural way to obtain this within type IIA string theory. Recall that this has space-filling D9-branes which are unstable. More specifically, there is a tachyonic scalar on its worldvolume. It has been shown by Sen [17] that all the branes on type II string theory may be realized by appropriate solitonic configuration involving the tachyon (and gauge) fields of the unstable D9-brane. In particular, a six-brane is an 't Hooft-Polyakov monopole, (and an anti-sixbrane is an anti-monopole). Let us start with an unstable D9-brane of type IIA theory in which the geometry is four dimensional Minkowski times a compact Calabi-Yau space (with fluxes). The tachyon condensation process leads to the formation of a D6-antiD6 pair as a monopole-antimonopole on the compact space \mathcal{B} . This appears to make sense for the following reasons. First, although the unstable D9-brane can decay into any of the lower dimensional branes, the first ones that would be stable are the D6-branes. This is because a Calabi-Yau space does not have a non-trivial five-cycle on which the D8-branes can be wrapped to be stabilized and the D7-branes are anyway unstable in type IIA theory. Secondly, on a compact space \mathcal{B} , charge conservation would require that one always forms branes and antibranes in pairs. Therefore, as the local (radially symmetric) formation of a brane tries to reduce energy, global issues (namely charge conservation) force us to end up with an unstable configuration again. Moreover, the formation of a monopole (sixbrane) at, say, the north pole of $\mathcal{B} = S^3$ will naturally produce an anti-monopole (anti-sixbrane) at the south pole. Of course, small irregularities would ensure that the branes do not form at *exact* antipodal points. From the four-dimensional point of view, the process outlined above will also produce some inflation [7] resulting in a couple of e-foldings. This *pre-inflation* is of fast roll type [9, 7] and might help in setting up appropriate initial conditions [33].

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