

Non-Topological Inflation from Embedded Defects

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Non-Topological Inflation from Embedded Defects

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We discuss a new mechanism of obtaining a period of cosmological inflation in the context of string theory. This mechanism is based on embedded defects which form dynamically on higher dimensional D-branes. Such defects generate topological inflation, but unlike topological inflation from stable defects, here there is a natural graceful exit from inflation: the decay of the embedded defect. We demonstrate the idea in the context of a brane-antibrane annihilation process. The graceful exit mechanism suggested here applies generically to all realizations of inflation on D-branes.

I. INTRODUCTION

There has been a lot of interest recently in exploring ways of obtaining a period of cosmological inflation from string theory. A realization of inflation in string theory should address some of the conceptual problems of quantum field theory realizations of inflationary cosmology [1]. Since string theory contains many moduli fields which are massless at the perturbative level, stringy inflation may yield a natural mechanism to produce fluctuations of the required small amplitude [2]. A description of inflation in the context of string theory would also automatically resolve the *trans-Planckian problem* [3] of inflationary cosmology, since it would yield a complete description of the dynamics of fluctuations during the entire inflationary period. Furthermore, since one of the goals of string theory is to provide a nonsingular cosmology, a stringy inflation model should also allow one to resolve the singularity problem [4] of inflationary cosmology.

Recent developments in particle phenomenology have explored the possibility that the our matter fields are confined to a four-dimensional space-time hypersurface in a higher-dimensional bulk space-time [5,6]. In the context of string theory this hypersurface arises naturally as a D-brane, one of possibly many branes that reside in the bulk spacetime [7]. However, for our purposes here, a sufficient starting point is some higher dimensional field theory (coupled weakly to gravity). This higher dimensional theory can be completed in the ultraviolet by embedding in critical string theory, but also by using the $(2,0)$ fixed point [8], little string theory [9,10], or by deconstruction [11]. In order to have a definite mental picture we assume the higher dimensional theory is embedded in string theory, and is realized on D-branes.

This “brane-world” scenario has generated new realizations of cosmic inflation. In one setup [12,13], the inflaton is the separation between a brane-antibrane pair. However, it has been shown [13,14] that in this context

it is not easy to obtain initial conditions which lead to a sufficient period of inflation. Another possibility is to have inflation generated by branes at an angle [15], or by a configuration of defects of different dimensionalities [16]. This has the advantage of possibly providing a mechanism for localization of chiral matter on the inflating locus.

An alternative way to obtain inflation from brane collisions was suggested in [17], based on the realization [18] that topological defects of co-dimension greater or equal to one are formed on the brane worldvolume during the collision of a brane-antibrane pair. In the simplest example, the order parameter of the phase transition which occurs when the brane-antibrane pair meets is given by a complex tachyon field [19] which condenses at a non-vanishing expectation value. The vacuum manifold of the tachyon field ground state values has the topology of S^1 . By causality [20], the values of the condensate in the vacuum manifold are uncorrelated on large length scales, and hence a network of co-dimension 2 defects inevitably will form. In the case of a $D5-\bar{D}5$ brane pair, the resulting topological defect is a $D3$ brane which could be our world. Provided that the thickness of the topological defect is larger than the Hubble radius, the defect core can undergo inflation (*topological inflation* [21]).

A problem with the mechanism of [17] (and of models of 4 dimensional topological inflation [21]) is the *graceful exit problem*: how does inflation end, and a transition to the usual late time cosmological evolution occur?

In this paper, we provide a simple solution to this problem, which generalizes to all above mentioned “brane-world” models of inflation. We point out that the defects which form in brane collisions will often be *embedded defects* (which are non-BPS branes in the string context) (see [22,23] for a discussion of embedded defects in field theory). These embedded defects can be stabilized at high temperatures by plasma effects [24,25], but decay at a sufficiently low temperature, thus providing a natural graceful exit mechanism from inflation. In the context

of field theory inflation, such a scenario was suggested in [26].

We may call the proposed inflationary model “Higher Dimensional (Non-) Topological Inflation” since inflation is taking place in the higher-dimensional brane worldvolume, not only in the directions of the defect (which are the dimensions of our four-dimensional physical space-time). Assuming that there is a mechanism to stabilize the radion in the two (compact) distinguished extra dimensions (the directions of the original brane orthogonal to the defect), we would have found a way of making **two** of the internal dimensions larger than the others, thus making contact to the scenario of [27] *. Alternatively, the two extra dimensions may be taken to be infinite; as one moves away from the core of the defect, inflation proceeds more slowly, and ends earlier. Thus the resulting spacetime is expected to be warped, making contact with the work of [6]. The question of which alternative is realized in specific examples needs a detailed, model dependent analysis, which we hope to return to in the future.

The causality argument (Kibble mechanism [20]) which ensures that in a phase transition leading to the possibility of topological defect formation such defects inevitably will form, with a separation smaller or comparable to the Hubble length, also ensures the formation of stabilized embedded defects in models which admit them, as will be explained later. Thus, there is no initial condition problem for our model of inflation.

In field theory models with a symmetry breaking scale η much smaller than the Planck scale, the width of a defect is typically much smaller than the Hubble length. Thus, topological inflation will not occur. An advantage of inflation in the context of higher dimensions is that there is a new scale, the brane tension, which is naturally higher than the Planck mass scale [†]. As will be shown in Section III, in this context the defect width can easily be larger than the Hubble length, thus making defect-driven inflation possible.

Our proposed mechanism of topological inflation from embedded defects can apply in situations more general than that of a brane-antibrane annihilation process. For example, in the context of brane intersections (the intersection region undergoing topological inflation), the branes may partially unwind around each other, leaving behind some non-inflating components. This process excites the off-diagonal open strings, which have no geometrical interpretation.

*Note that one also obtains two internal dimensions larger than the others in the “dynamical de-compactification” scenario resulting from brane gas cosmology [28,7]. However, in that context there is no reason to expect a large hierarchy of scales.

[†]This is a standard useful property of D-branes: they can be very heavy without causing large back-reaction.

The outline of this paper is as follows. In the next section, we discuss a particular example in which one can realize topological inflation from an embedded defect, namely embedded defects forming in a brane-antibrane annihilation process. Next, we discuss how topological inflation from embedded defects has a natural graceful exit mechanism. Finally, we summarize our results, and discuss further applications of the basic idea of topological inflation from embedded defects in the context of brane physics.

II. EMBEDDED DEFECTS FROM HIGHER DIMENSIONAL D-BRANES

We will start with a brief review of the original scenario of [17]. The scenario of [17] takes as a starting point a pair of parallel $D5$ and $\bar{D}5$ branes approaching each other. As studied in [19], when the pair gets sufficiently close, a tachyon instability sets in. For minimal gauge content of the branes, the tachyon can be described by a complex scalar field ϕ with a standard symmetry breaking (*Mexican hat*) potential

$$V(\phi) = \lambda(|\phi|^2 - \eta^2)^2, \quad (1)$$

where the point $\phi = 0$ corresponds to branes at zero separation, and η is the symmetry breaking scale. The vacuum at $|\phi| = \eta$ corresponds to the closed string vacuum, where the original branes have annihilated.

The vacuum manifold of the tachyon condensates is S^1 , and hence according to the usual Kibble argument [20], the fact that the condensate values must be uncorrelated on scales larger than the Hubble radius inevitably leads to the formation of stable co-dimension 2 defects, in this case $D3$ branes, on the worldvolume of the original $D5$ branes. Unlike most of the discussion of defects on annihilating D-branes, we are interested in the case where the resulting $D3$ brane is not BPS saturated, so that inflation is possible.

In terms of gauge theory on the original brane worldvolume, the above process corresponds to the symmetry breaking

$$U(1) \times U(1) \rightarrow U(1), \quad (2)$$

where one of the $U(1)$ factors on the left hand side of the equation lives on either of the branes, and the unbroken subgroup corresponds to the diagonal $U(1)$ factor. The standard homotopy arguments (see e.g. [29] for a discussion of such arguments in the context of topological defects) yield

$$\Pi_1(\mathcal{M}) \equiv \Pi_1(G/H) = \mathcal{Z}, \quad (3)$$

where \mathcal{M} is the vacuum manifold, and G and H are the

full gauge group and the residual gauge group after symmetry breaking, respectively [‡].

This set-up is not restricted to annihilating D-branes. Suppose we start with any higher dimensional D-brane configuration, and that at initial times the set-up is removed from its vacuum. This vacuum may be the closed string vacuum, but also an open string vacuum, where the branes align in a supersymmetric fashion. In the process of relaxation to the vacuum, the Kibble mechanism will ensure the creation of local defects, for example brane intersections. It is those defects which seed inflation in the higher-dimensional topological context, and which are the subject of our discussion. In particular, the energy scale associated with this initial phase transition can be much smaller than the string scale, so the above (higher dimensional) field theory considerations apply.

Our idea, in this context, is to concentrate on the case of embedded, non-topological, defects. For this end we enhance the gauge symmetry on each of the branes, e.g. from $U(1)$ to $SU(2)$. Indeed, this arises naturally in the context of fivebranes in type I theory[§]. In this case the worldvolume gauge group is $SU(2) \times SU(2) = SO(4)$. The tachyonic scalar fields transform as $(2, 2)$ under $SU(2) \times SU(2)$ ^{**}.

Now, the symmetry breaking pattern during tachyon condensation becomes

$$SU(2) \times SU(2) \rightarrow SU(2). \quad (4)$$

In this case, no stable topological defects (in particular no co-dimension 2 defects) form during this phase transition since the vacuum manifold is $\mathcal{M} = S^3$ and hence $\Pi_1(\mathcal{M}) = 1$. Note that in the context of studies of branes in Type I string theory it is already known that the $D3$ brane is unstable [33].

The situation is similar to what happens in the electroweak theory where (for vanishing Weinberg angle), the symmetry breaking is $SU(2) \rightarrow 1$, as effectively in the above case (4). However, as is well known from analysis of the electroweak theory [23], it is possible to construct *embedded defects*, solutions of the equations of motion which correspond to unstable defects. In the case of the electroweak theory in four space-time dimensions, the defects are *electroweak strings*, an example of embedded strings ^{††}.

[‡]In the usual discussion this topological charge is the central charge appearing in the supersymmetry algebra, and the resulting defect is then BPS saturated. For inflation to occur, this charge should not be that central charge, and the defect should be non-BPS. In any event, the embedded defects discussed later are always non-BPS.

[§]Alternatively, our D-brane set-up can be placed at a Z_2 orbifold singularity [30]

^{**}There are also other matter fields, which are assumed, like all moduli, to have been stabilized by a separate mechanism.

^{††}Another example of embedded strings occurs in the sigma

As was recently realized [24], such embedded defects can be stabilized at finite density by plasma effects. Consider again first the example of the standard electroweak theory. The order parameter of the symmetry breaking phase transition has four real components, two of which are electrically charged, two are electrically neutral. In the presence of a thermal bath of photons, the effective potential will then be lifted in the charged scalar field directions by more than it is lifted in the neutral scalar field directions. Thus, the vacuum manifold corresponding to the theory in a photon plasma is S^1 and not S^3 . This can stabilize a subset of embedded strings, namely the string configurations constructed from the neutral scalar fields (these are the electroweak Z-strings in the case of the standard electroweak theory).

Note that if the embedded defects are thermally stabilized at the temperature just below the phase transition temperature, then the causality argument (Kibble mechanism [20]) applies and states that at least of the order one such defect will form per correlation volume (which in turn is smaller than the Hubble volume). The argument is as follows: during the phase transition, the order parameter relaxes to the vacuum manifold S^3 . However, in the presence of thermal stabilizing effects it will relax to the reduced vacuum manifold S^1 . However, which value in S^1 is taken on is random on scales larger than the correlation length. Hence, defects will form. This process was studied in detail in the case of semilocal strings (related to embedded strings) in [34].

Our idea is to apply this mechanism to brane cosmology. We need to assume that the branes contain a thermal bath of gauge fields which couple in the same way to the $SU(2)$ symmetry breaking field of our brane setup as the photon does to the $SU(2)$ symmetry breaking order parameter in the electroweak theory. Ideally, the usual photon field could play this role. In this case, it would be natural to have only this field in thermal equilibrium at late times, since it would be the only massless gauge field. The situation would then be identical to the electroweak theory discussed above. Given the above assumptions about field content and coupling to the $SU(2)$ symmetry breaking order parameter, it is completely natural from the point of view of cosmology (hot initial state in the very early Universe) to have the field excited in the required way.

In this case, the phase transition corresponding to brane collision will, by the Kibble mechanism [20], inevitably produce a network of co-dimension 2 branes. These defects will remain stable until the matter density on the branes has sufficiently diluted.

Let us now be a bit more specific. Given an $SU(2)$

model description of low energy QCD in the limit of vanishing bare quark masses [31]. In this case, the embedded strings are the pion strings [32].

gauge group on each of the $D5$ branes, the world volume theory of a coincident $D5$ - $\bar{D}5$ brane pair is $SU(2) \times SU(2)$, and any gauge field configuration can be written as a 4×4 matrix of complex entries, the upper diagonal 2×2 entries corresponding to a $SU(2)$ matrix describing open strings which start and end on the first brane, the lower diagonal 2×2 entries describing a $SU(2)$ matrix corresponding to the open strings starting and ending on the other brane. The two off-diagonal 2×2 sub-matrices are $SU(2)$ matrices corresponding to open strings beginning and ending on different branes. It is these latter sectors which contain the tachyon [35]. In the case of $U(1)$ gauge fields on each of the branes, the tachyon is a complex scalar field (one real component from each of the sectors), in the case of $SU(2)$ gauge fields, there are two real components in each sector.

Let us now assume that on each of the branes there is a thermal bath of a $U(1)$ subgroup of $SU(2)$, and the other gauge fields assumed to be vanishing. In this case, the gauge fields excited on the branes will only couple to one of the two real tachyon fields in each sector. Let us combine these two fields into a complex field χ (charged with respect to the excited gauge field), the other two into a complex field ϕ (neutral with respect to the excited gauge field). The effective world volume Lagrangian then becomes

$$\mathcal{L}_{\text{eff}} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \frac{1}{2}\bar{D}_\mu\chi D^\mu\chi + \frac{1}{2}\partial_\mu\bar{\phi}\partial^\mu\phi - V(\phi, \chi) \quad (5)$$

where the D_μ stands for the gauge covariant derivative, the gauge field being the one excited on the brane. Hence, the gauge field directly couples only to the charged scalar field χ and not to the neutral field ϕ . The field strength tensor of the gauge field A_μ excited on the brane is denoted by $F_{\mu\nu}$, but the F^2 term in the Lagrangian will not be important in the following. The tachyon potential V has the typical symmetry breaking form

$$V(\phi, \chi) \equiv V_0(\phi, \chi) = \lambda(|\phi|^2 + |\chi|^2 - \eta^2)^2, \quad (6)$$

where λ is a coupling constant, and η is the ground state expectation value of the tachyon field magnitude. The corresponding vacuum manifold is S^3 .

The above effective Lagrangian is the same which describes the sigma model of low energy QCD in the chiral limit, coupled to an external bath of photons [24,31]. It also describes the standard electroweak theory with only the photon field excited [25]. In the presence of a thermal bath of A_μ gauge fields, we can take the thermal average of the Lagrangian. Terms linear in A_μ vanish, the quadratic term $A_\mu A^\mu$ becomes κT^2 , where T is the temperature of the bath. Thus, an effective mass term for the χ field is generated. The effective potential becomes

$$V(\phi, \chi) = V_0(\phi, \chi) + g^2\kappa T^2|\chi|^2. \quad (7)$$

Hence, the vacuum manifold is lifted in the charged scalar field directions. The vacuum manifold of the effective

potential becomes S^1 . Thus, co-dimension 2 defects in which the neutral scalar field configuration takes on the Nielsen-Olesen [36] form become important as metastable defects.

III. EMBEDDED DEFECTS, INFLATION, AND GRACEFUL EXIT

The idea of topological inflation is quite simple [21]: provided that the core of the defect is comparable or larger than the Hubble radius, then the potential energy density in the defect core will be larger than the tension energy, and thus via standard minimal coupling to gravity, the core will commence inflationary expansion. For defects associated with an energy scale much smaller than the Planck scale, the width is much too small to obtain topological inflation. However, if the tension of the defect is given by an energy scale η which is comparable or larger than m_{pl} , then the condition for topological inflation is satisfied. A quantitative analysis [37] yields as condition for topological inflation

$$\eta > \frac{1}{4}m_{\text{pl}}. \quad (8)$$

Let us give a brief derivation of this result. For a typical symmetry breaking potential of the form (1), the defect width w is of the order [36]

$$w \simeq \lambda^{-1/2}\eta^{-1}, \quad (9)$$

which can be seen by balancing potential and tension energies. In the core, the potential energy density dominates, in the outside regions of the defect the tension energy density is larger. The condition for defect inflation is that the core width is greater than the Hubble length H^{-1} . Applying the Friedmann equation

$$H^2 = \frac{8\pi}{3}G\rho \quad (10)$$

(where ρ is the energy density) to the defect core (where $\rho \simeq V(\phi = 0)$), the condition becomes

$$\frac{\eta}{M_{\text{pl}}} > \frac{1}{2\sqrt{2}}. \quad (11)$$

Thus, for values of η comparable or larger than m_{pl} but coupling constant $\lambda \ll 1$, both the condition (11) for topological inflation and the condition $V(0) < m_{\text{pl}}^4$ for applicability of the Friedmann equations for the evolution of the background space-time are satisfied.

The above conditions can easily be satisfied for defect arising on D-branes in superstring theory. In this context, it has been shown that the tachyon potential, computed in superstring field theory, has the form [38] (see also [39])

$$V(T, \bar{T}) = \tau(\alpha')^2(|T|^2 - (\alpha')^{-1})^2, \quad (12)$$

where τ is proportional to the brane tension and α' is the string tension. Thus, if $\alpha' < m_{\text{pl}}^2$, then the conditions for topological inflation are satisfied. This condition requires the (six dimensional) string coupling to be small.

Topological inflation from stable defects suffers from a graceful exit problem. Inflation in the defect core continues forever, and in order to make a successful transition to late time cosmology one usually has to postulate that our region of space originated from a region close to the original defect boundary. This is already problematic in the case of field theory inflation in four dimensional space-time. Here the problem becomes even more acute, since in the context of brane world models, our matter fields have to be localized on the locus that is to become our physical spacetime.

However, if the inflationary expansion takes place in the core of an embedded defect, the graceful exit problem is naturally resolved: after some finite time, the defect decays, the potential energy density disappears, and inflation stops.

As was shown in [31], embedded defects typically acquire superconducting current. These currents help stabilize the embedded defects down to very low temperatures. At the temperature T_c below which the thermal barrier becomes too low to stabilize the embedded defect, a core phase transition [40] occurs (the charged scalar fields acquire nonvanishing expectation values), but the defect persists [25]. The defect decays only when the gauge field responsible for the stabilization of the defect ceases to be in thermal equilibrium. Let us assume that it is the photon which is responsible for the stabilization of our embedded $D3$ brane. In this case, the decay temperature T_d could be ^{‡‡} as low as

$$T_d \sim 10^{-10} \text{GeV}, \quad (13)$$

the temperature of recombination ^{§§}. If the temperature at the onset of inflation is about 10^{17}GeV , in light of the formula (12) for the tachyon potential in string field theory quite a reasonable value (it is also the upper bound on the scale of inflation in order not to overproduce gravitational waves), then we (optimistically) obtain about 62 e-foldings of inflation, only slightly larger than the minimum number of e-foldings required [41] for inflation to solve the horizon and flatness problems of standard cosmology. Note that for this small number of e-foldings, the physical wavelength of comoving scales which correspond to present day cosmological scales was larger than

^{‡‡}However, this is an optimistic assumption since we are considering an epoch during which the Universe is exponentially expanding - we thank J. Khoury, W. Kinney and T. Baltz for a discussion on this point.

^{§§}Note that after decay, the Universe will reheat to a very high temperature as long as some of the energy of the embedded defect goes into standard particle physics model matter.

the Planck length at the beginning of inflation. Thus, there is no trans-Planckian problem [3] in this model.

As mentioned in the Introduction, in this mechanism of topological inflation from embedded defects produced during 5 and $\bar{5}$ brane annihilations, two of the internal dimensions (namely those parallel to the worldvolume of the initial branes) also inflate. Given the low number of e-foldings of inflation which result in this scenario, the size of the two large extra dimensions is comfortably below the observational bounds (assuming that there is a radion stabilization mechanism which sets after inflation in for these directions).

IV. SUMMARY AND DISCUSSION

In this paper we have proposed a new way of obtaining inflation in the context of brane physics. The mechanism is based on inflation taking place in the core of an embedded defect. The embedded defect is stabilized at high temperatures by plasma effects, as studied in [24], but becomes unstable and decays at lower temperatures, thus providing a natural graceful exit mechanism from inflation.

We have suggested a specific realization of this idea in the context of a brane-antibrane setup, where our world volume corresponds to that of an embedded 3 brane which is generated in the process of tachyon condensation of an annihilating 5 brane - antibrane pair. To obtain an embedded defect in this setup, the gauge groups on the 5 branes must be enhanced from the minimal gauge group $U(1)$.

Clearly this idea is more general than the brane-antibrane setup. For example, in the context of brane intersections [42] an embedded defect can be realized if the gauge content on the branes is enlarged. Some intersections can unwind as the temperature cools, a process described in the field theory as the decay of the embedded defect, and requires one to excite off-diagonal, non-geometrical open strings (see for example [43] for a detailed discussion).

More generally, it is well-known, at least in the BPS sector, that stability of defects depends on parameters of the theory (see e.g. [44]). This is well controlled for supersymmetric states, by using the BPS formula, but is expected to be the case generally. In this case, one can use the topological inflation idea, and provide a natural graceful exit mechanism along the lines of this paper. In this scenario the initial stable defect causes inflation. As inflation proceeds, and effective parameters of the theory change in response to cooling down, causing the initial defect to become unstable. This paper is a concrete example of this more general scenario.

Obviously, our scenario of “topological inflation” from an embedded defect can also be realized in ordinary four dimensional quantum field theory. However, in that context it appears unnatural to obtain defects with the

width required for topological inflation, unless field values larger than the Planck scale are invoked. As we have seen, in the context of the brane setup, the existence of new physical scales (e.g. the brane tensions) makes it easy to obtain defects of sufficient width to support topological inflation.

We have also seen that in our realization of embedded defect inflation in the context of an annihilating 5 and $\bar{5}$ brane pair, inflation can produce a hierarchy in the sizes of the internal dimensions (or a warped geometry), since the two extra spatial dimensions parallel to the worldvolume of the 5 branes also expand exponentially.

Recently, it was argued that de Sitter space vacua in string theory will suffer from conceptual problems [45,46], if the lifetime of de Sitter space lasts longer than the Poincaré recurrence time $t_r \sim e^{S_0}$, where S_0 is the entropy of the space. For example, Kachru et. al provided a stringy realization of a meta-stable de Sitter space which tunnels to flat space-time in a timescale shorter than t_r [47]. Our scenario is more reminiscent of the new inflationary models, in that no tunneling is required to stop inflation. Consequently our inflationary spacetime turns into a hot FRW space-time after ~ 62 e-foldings, a time scale exponentially shorter than t_r .

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