Warped reheating in brane–antibrane inflation

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Abstract. We examine how reheating occurs after brane–antibrane inflation in warped geometries, such as those which have recently been considered for type IIB string vacua. We adopt the standard picture that the energy released by brane annihilation is dominantly dumped into massive bulk (closed-string) modes which eventually cascade down into massless particles, but argue that this need not mean that the result is mostly gravitons with negligible visible radiation on the Standard Model brane. Rather, in geometries having strongly warped throats we argue that this energy can instead be predominantly dumped into massless modes which are localized at the throat’s tip, because of the exponential growth of the massive Kaluza–Klein wavefunctions there. The result is preferentially to reheat those branes which lay in the most strongly warped throats. We argue that the efficiency of this process removes a conceptual obstacle to the construction of multi-throat models, wherein inflation occurs in a different throat to the one in which the Standard Model brane resides. Such multi-throat models are desirable because they may resolve a difficulty in reconciling inflation with the supersymmetry-breaking scale on the Standard Model brane, and because they may allow cosmic strings to be sufficiently long-lived to be observable during the present epoch.

Keywords: cosmology with extra dimensions, string theory and cosmology, inflation
1. Introduction

There has been significant progress over the past few years towards the construction of *bona fide* string-theoretic models of inflation. The main progress over early string-inspired supergravity [1] and BPS-brane based [2] models has come due to the recognition that brane–antibrane [3,4] and related [5,6] systems can provide *calculable* mechanisms for identifying potentially inflationary potentials. Even better, they can suggest new observable signatures, such as the natural generation of cosmic strings by the brane–antibrane mechanism [3,7]. The central problem to emerge from these early studies was that of understanding how the many string moduli get fixed, since such an understanding is a prerequisite for a complete inflationary scenario.

Recent developments are based on current progress in modulus stabilization within warped geometries with background fluxes for type IIB vacua [8]–[10]. Both brane–antibrane inflation [11] and modulus inflation [12] have been embedded into this context, with an important role being played in each case by branes living in strongly warped ‘throat-like’ regions within the extra dimensions. These inflationary scenarios have generated considerable activity [13,14] because they open up the possibility of asking in a more focused way how string theory might address the many issues which arise when building inflationary models. For instance, one can more fully compute the abundance and properties of any residual cosmic strings which might survive into the present epoch [16]. Similarly, the possibility of having quasi-realistic massless particle spectra in warped, fluxed type IIB vacua [17] opens up the possibility of locating where the known elementary particles fit into the post-inflationary world [18], a prerequisite for any understanding of reheating and the subsequent emergence of the hot big bang.

Even at the present preliminary level of understanding, a consistent phenomenological picture seems to require more complicated models involving more than a single throat (in addition to the orientifold images)\(^4\). This is mainly because for the single-throat models the success of inflation and particle physics phenomenology place contradictory demands.

\(^4\) Two-throat models are also considered for reasons different to those given here in [15].
on the throat’s warping. They do so because the energy scale in the throat is typically required to be of order $M_i \sim 10^{15}$ GeV to obtain acceptably large temperature fluctuations in the CMB. But as was found in [18], this scale tends to give too large a supersymmetry-breaking scale for ordinary particles if the Standard Model (SM) brane resides in the same throat. This problem appears to be reasonably generic to the KKLT-type models discussed to date, because these models tend to have supersymmetric anti-de Sitter vacua until some sort of supersymmetry-breaking physics is added to lift the vacuum energy to zero. The problem is that the amount of supersymmetry breaking required to zero the vacuum energy also implies so large a gravitino mass that it threatens to ruin the supersymmetric understanding of the low-energy electroweak hierarchy problem.

No general no-go theorem exists, however, and there does appear to be considerable room for trying to address this issue through more clever model-building. Reference [19] provides a first step in this direction within the framework of ‘racetrack’ inflation [12]. Another possibility is a picture having two (or more) throats, with inflation arising because of brane–antibrane motion in one throat but with the Standard Model situated in the other (more about this proposal below). By separating the scales associated with the SM and inflationary branes in this way, it may be possible to reconcile the inflationary and supersymmetry-breaking scales with one another.

Besides possibly helping to resolve this problem of scales, multi-throat models could also help ensure that string defects formed at the end of inflation in the inflationary throat have a chance of surviving into the present epoch and giving rise to new observable effects [16]. They are able to do so because if the Standard Model were on a brane within the same throat as the inflationary branes, these defects would typically break up and disappear by intersecting with the SM brane.

At first sight, however, any multi-throat scenario seems likely to immediately founder on the rock of reheating. Given the absence of direct couplings between the SM and inflationary branes, and the energy barrier produced by the warping of the bulk separating the two throats, one might expect the likely end-point of brane–antibrane annihilation to be dumping energy only into closed-string, bulk modes, such as gravitons, rather than visible degrees of freedom on our brane. In such a universe the energy which drove inflation could be converted almost entirely into gravitons, leaving our observable universe out in the cold.

It is the purpose of the present work to argue that this picture is too pessimistic, because strongly warped geometries provide a generic mechanism for channelling the post-inflationary energy into massless modes localized on the throat having the strongest warping. They can do so because the massive bulk Kaluza–Klein (KK) modes produced by brane–antibrane annihilation prefer to decay into massless particles which are localized on branes within strongly warped throats rather than decaying to massless bulk modes. As such, they open a window for obtaining acceptable reheating from brane–antibrane inflation, even if the inflationary and SM branes are well separated on different throats within the extra dimensions.

The remainder of the paper is organized as follows. In section 2 we introduce a simple generalization of the Randall–Sundrum (RS) model [21] containing two AdS$_5$ throats with different warp factors, as a tractable model for the KLMT inflationary scenario [11] with

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$^5$ See [20] for a discussion of issues concerning brane-related reheating within other contexts.
two throats. Here we recall the form of the KK graviton wavefunctions in the extra dimension. This is followed in section 3 by an account of how the tachyonic fluid describing the unstable brane–antibrane decays into excited closed-string states, which quickly decay into KK gravitons. Section 4 gives an estimate of the reheating temperature on the SM brane which results from the preferential decay of the KK gravitons into SM particles. Our conclusions are given in section 5.

2. Tale of two throats

We wish to describe reheating in a situation where brane–antibrane inflation occurs within an inflationary throat having an energy scale of \(M_i\), due to the warp factor \(a_i = M_i/M_p\), where \(M_p\) is the 4D Planck mass. This throat is assumed to be separated from other, more strongly warped, throats by a weakly warped Giddings–Kachru–Polchinski (GKP) manifold [8] whose volume is only moderately larger than the string scale, so \(M_s \lesssim M_p\). In the simplest situation there are only two throats (plus their orientifold images), with the non-inflationary (Standard Model) throat having warp factor \(a_{sm} \ll a_i\).

There are two natural choices for the SM warp factor, depending on whether or not the SM brane strongly breaks 4D supersymmetry. For instance, if the SM resides on an anti-D3 brane then supersymmetry is badly broken and the SM warp factor must describe the electroweak hierarchy à la Randall and Sundrum [21], with \(a_{sm} \sim M_W/M_p \sim 10^{-16}\). Alternatively, if the SM resides on a D3 or D7 brane which preserves the bulk’s \(N = 1\) supersymmetry in 4D, then SUSY breaking on the SM brane is naturally suppressed by powers of \(1/M_p\) because it is only mediated by virtual effects involving other SUSY-breaking anti-D3 branes. In this case the electroweak hierarchy might instead be described by an intermediate-scale scenario [22], where \(a_{sm} \sim M_{int}/M_p \sim (M_W/M_p)^{1/2} \sim 10^{-8}\).

A potential problem arises with the low-energy field-theory approximation if \(a_{sm} < a_i^2\), because in this case the string scale in the SM throat, \(M_{sm} \sim a_{sm} M_p\), is smaller than the inflationary Hubble scale \(H_i \sim M_i^2/M_p \sim a_i^4 M_p\) [23]. In this case string physics is expected to become important in the SM throat, and stringy corrections may change the low-energy description. The intermediate scale is somewhat more attractive from this point of view, since for it the field-theory approximation may be justified if the underlying GKP geometry’s volume is somewhat larger than the string scale. In this case one has instead \(M_s \sim f M_p\), for \(f < 1\), and so \(M_i \sim f a_i M_p\), \(M_{sm} \sim f a_{sm} M_p\), and \(H_i \sim a_i^2 f^2 M_p\). This implies \(H_i/M_{sm} \sim f (a_i^2/a_{sm})\), and so small \(f\) can ensure \(H_i \ll M_{sm}\) even if \(a_i^2 \sim a_{sm}\).

To proceed we use the fact that within the GKP compactification the geometry within the throat is well approximated by

\[
d s^2 = a^2(y)(d\tau^2 - dx^2) - dy^2 - y^2 d\Omega_5^2, \tag{2.1}
\]

where \(y\) represents the proper distance along the throat, \(a(y) = e^{-k|y|}\) is the throat’s warp factor and \(d\Omega_5^2\) is the metric on the base space of the corresponding conifold singularity of the underlying Calabi–Yau space [24]. Of most interest is the 5D metric built from the observable four dimensions and \(y\), which is well approximated by the metric of five-dimensional anti-de Sitter space.

A simple model of the two-throat situation then consists of placing inflationary brane–antibranes in a throat at \(y = -y_i\) and putting the Standard Model brane at \(y = +y_{sm}\), as is illustrated in figure 1. Our analysis of this geometry follows the spirit of [25]. Since
most of the interest is in the throats, we simplify the description of the intervening bulk geometry by replacing them with a Planck brane at $y = 0$, with the resulting discontinuity in the derivative of $a(y)$ chosen to reproduce the smoother (but otherwise similar) change due to the weakly warped bulk. This approximation is illustrated in figure 2, with the smooth dashed curve representing the warp factor in the real bulk geometry and the solid spiked curve representing the result using an intervening Planck brane instead.

Of particular interest in what follows are the massive Kaluza–Klein modes in the bulk, since these are arguably the most abundantly produced modes after brane–antibrane annihilation. For instance, focusing on the five dimensions which resemble AdS space in the throat, a representative set of metric fluctuations can be parameterized as $h(x, y)$ in the line element,

$$
ds^2 = a^2(y)(dt^2 - dx^2 + h_{\mu\nu} dx^\mu dx^\nu) - dy^2.
$$

(2.2)

In the static AdS background, the KK modes have spatial wavefunctions of the form

$$
h(x, y) = \sum_n \phi_n(y) e^{ip \cdot x}.
$$

(2.3)

with $p \cdot x = -E_n t + p \cdot x$, and $\phi_n(y)$ satisfying the equation of motion

$$
-\frac{d}{dy}\left(e^{-4k|y|} \frac{d\phi_n}{dy}\right) = m_n^2 e^{-2k|y|} \phi_n.
$$

(2.4)

Here $m_n^2 = p \cdot p$ is the mode’s 4D mass as viewed by brane-bound observers.
Exact solutions for $\phi_n(y)$ are possible in the Planck-brane approximation [21, 25, 26], and are linear combinations of Bessel functions times an exponential:

$$\phi_n(y) = N_n e^{2k|y|} \left[ J_2 \left( \frac{m_n}{k} e^{k|y|} \right) + b_n Y_2 \left( \frac{m_n}{k} e^{k|y|} \right) \right]$$  \hspace{1cm} (2.5)

where for low-lying KK modes ($m_n \ll k$) one has

$$b_n \approx \frac{\pi m_n^2}{4k^2}$$  \hspace{1cm} (2.6)

while for heavy KK modes ($m_n/k \approx 1$) one has

$$b_n \approx -0.47 + 1.04 \left( \frac{m_n}{k} \right).$$  \hspace{1cm} (2.7)

$N_n$ is determined by the orthonormality condition, which ensures that the kinetic terms of the KK modes are independent of $a_{\text{sm}}$:

$$\int_{-y_{\text{sm}}}^{y_{\text{sm}}} dy e^{-2k|y|} \phi_n \phi_m = \delta_{nm}.$$  \hspace{1cm} (2.8)

These wavefunctions are graphed schematically in figure 3.

For strongly warped throats it is the exponential dependence which is most important for the KK modes. Because of the exponential arguments of the Bessel functions in equation (2.5), the presence of the Bessel functions modifies the large-$y$ behaviour slightly. Due to the asymptotic forms $J_2(z) \propto z^{-1/2}$ for large $|z|$, and similarly for $Y_2(z)$, we see that $\phi_n(y) \sim e^{3k|y|/2}$ for $m_n e^{k|y|} \gg k$. It is only this behaviour which we follow from here on. Taking the most warping to occur in the SM throat we find that $m_n$ is approximately quantized in units of $M_{\text{sm}} \equiv a_{\text{sm}} M_p$, which is either of order $M_W \sim 10^3$ GeV or $M_{\text{int}} \sim 10^{10}$ GeV depending on whether or not supersymmetry breaks on the SM brane. Keeping only the exponentials we find that orthonormality requires

$$N_n^{-2} \sim \int_{-y_{\text{sm}}}^{y_{\text{sm}}} dy e^{-2k|y|} (e^{3k|y|/2})^2 \sim (k a_{\text{sm}})^{-1},$$

and so

$$\phi_n(y) \sim (a_{\text{sm}} k)^{1/2} e^{3k|y|/2}.$$  \hspace{1cm} (2.9)
showing that these modes are strongly peaked deep within the throat. This is intuitively easy to understand, since being localized near the most highly warped region allows them to minimize their energy most effectively.

Thus, even the most energetic KK modes still have exponentially larger wavefunctions on the TeV brane, with the more energetic modes reaching the asymptotic region for smaller $y$. This is illustrated in figure 4, which shows $\ln |\phi_n(y)|$ versus $y$ in the representative case of a throat having warp factor $a = e^{-10}$, for a series of KK states with masses going as high as $M_n$ ($n = 20,000$). As the figure shows, the wavefunctions grow exponentially toward the TeV brane, with the onset of the asymptotic exponential form setting in earlier for larger mode number. This behaviour is central to the estimates which follow.

Among the KK modes it is the zero modes which are the exceptional case because their wavefunction is constant, $\phi_0 \sim \sqrt{k}$, and so they are not exponentially peaked inside the throat. It is the strong exponential peaking of the lightest massive KK modes relative to the massless modes which is central to the reheating arguments which follow.

In our simplified model, the presence of two throats is not much more complicated than the original RS model. Mathematically it is the same, except that RS identified the two sides $y \leftrightarrow -y$ through orbifolding. Instead we interpret them as two separate throats with different depths defined by the brane locations $-y_l$ and $+y_{sm}$. One can imagine doubling this entire system on $S_1$ and orbifolding as shown in figure 5, to define the boundary conditions on the metric and its perturbations at the infrared branes. In this figure, the orbifold identification acts horizontally so that the inflation and SM branes are distinct fixed points.

### 3. Brane–antibrane annihilation

In brane–antibrane inflation the energy released during reheating is provided by the tensions of the annihilating branes. Although this annihilation process is not yet

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Figure 4. Unnormalized wavefunctions for highly excited KK gravitons with KK numbers $n = 1, 100, 1000$ and $20,000$. 

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6 For the lowest-lying modes having the smallest nonzero masses it can happen that the asymptotic form of the Bessel functions is not yet reached even when $y = y_{sm}$, in which case the exponential peaking is slightly stronger than discussed above.
completely understood, present understanding indicates that the energy released passes through an intermediate stage involving very highly excited string states, before generically being transferred into massless closed-string modes. The time frame for this process is expected to be the local string scale.

For instance many of the features of brane–antibrane annihilation are believed to be captured by the dynamics of the open-string tachyon which emerges for small separations for those strings that stretch between the annihilating branes\(^7\). In flat space and at zero string coupling \((g_s = 0)\), the annihilation instability has been argued to be described by the following tachyon Lagrangian \([28]\):

\[
\mathcal{L}_T = -2\tau_0 e^{-|T|^2/l_s^2} \sqrt{1 - |\partial_\mu T|^2}
\]

where \(T\) is the complex tachyon field, \(\tau_0\) is the tension of either of the branes and \(l_s\) is the string length scale. During inflation, when the brane and antibrane are well separated, \(\dot{T} = 0\) and the pressure of the system \(p_i\) is simply the negative of the tension of the two branes, \(p_i = -\rho_i\), while afterward \(\dot{T} \rightarrow 1\) and \(p_i \rightarrow 0\). In this description the pressureless tachyonic fluid would dominate the energy density of the universe and lead to no reheating whatsoever.

However, for nonvanishing \(g_s\), the time evolution of the tachyon fluid instead very quickly generates highly excited closed-string states \([29,30]\). For \(Dp\) systems with \(p > 2\) the rate of closed-string production in this process is formally finite, whereas it diverges for \(p \leq 2\) (and so passes beyond the domain of validity of the calculation). This divergence is interpreted as meaning that for branes with \(p \leq 2\) all of the energy liberated from the initial brane tensions goes very efficiently into closed-string modes. For spatially homogeneous branes with \(p > 2\) the conversion is less efficient and so can be dominated by other, faster processes. In particular, it is believed that these higher-dimensional branes will decay more efficiently inhomogeneously, since they can then take advantage of the more efficient channels which are available to the lower-dimensional branes. For example a \(D3\) brane could be regarded as a collection of densely packed but smeared-out \(D0\) branes, each of which decays very efficiently into closed strings. Since the decay time

\(^7\) See, however, \([27]\) for a discussion of an alternative mechanism for which the relevant highly excited strings are open strings, but for which the annihilation energy nonetheless eventually ends up in massless closed-string modes.
is of order the local string scale, $l_s = 1/M_s$, the causally connected regions in this kind of
decay are only of order $l_s$ in size, and so have a total energy of order the brane tension
$\tau_0 \sim M_s^4/g_s$.

These flat-space calculations also provide the distribution of closed-string states as a
function of their energy. The energy density deposited by annihilating $D3$ branes into
any given string level is of order $M_s^4$, and so due to the exponentially large density of
excited string states the total energy density produced is dominated by the most highly
excited states into which decays are possible. Since the available energy density goes like
$1/g_s$, the typical closed-string state produced turns out to have a mass of order $M_s/g_s$,
corresponding to string mode numbers of order $N \sim 1/g_s$. On the other hand, the
momentum transverse to the decaying branes for these states turns out to be relatively
small, $p_T \sim M_s/\sqrt{g_s}$ [29,30], and so the most abundantly produced closed-string states
are nonrelativistic.

How do these flat-space conclusions generalize to the warped type IIB geometries
which arise in string inflationary models? If the annihilating 3-branes are localized in the
inflationary throat, then the tension of the annihilating branes is of order
$\tau_0 \sim (a_i M_s)^4/e^{\phi_i}$, where $\phi_i$ denotes the value taken by the dilaton field at the throat’s tip. The highly
excited closed-string states that are produced in this way live in the bulk, with the energy
density produced being dominated by those whose masses are of order $a_i M_s/e^{\phi_i}$. Once
produced, these closed-string bulk modes decay down to lower energies and, as might
be expected from phase-space arguments, most of them typically drop down to massless
string states very quickly. An important exception to this would arise for those states
carrying the most angular momentum at any given string mass level, since these must
cascade more slowly down to lower energies in order to lose their angular momentum [31].
However these seem unlikely to be produced in appreciable numbers by brane–antibrane
annihilation.

We are led in this way to expect that the annihilation energy is distributed relatively
quickly amongst massless string states, or equivalently to KK modes of the higher-
dimensional supergravity which describes these states. Although the initial massive string
modes would be nonrelativistic, with $M \sim a_i M_s/e^{\phi_i}$ and $p_T \sim a_i M_s/e^{\phi_i}/2$, the same need
not be true for the secondary string states produced by their decay, whose masses are
now of order the KK mass scales. Consequently these states may be expected not to
remain localized in the inflationary throat, and so if the extra dimensions are not too
large compared with the string scale these modes would have time to move to the vicinity
of the SM throat before decaying further. Once there, they would be free to fall into the
potential wells formed by the throats as their energy is lost by subsequent decays into
lower energy levels.

This physical picture is supported by the exponential peaking of the KK-mode
wavefunctions in the most deeply warped throats. In order to estimate the efficiency
with which energy can be transferred amongst KK modes, we can use the approximate
behaviour of the wavefunctions given in the previous section to keep track of powers of
the throat’s warp factor, $a_{sm}$. For instance, consider the trilinear vertex among three KK
states having mode numbers $n_1$, $n_2$ and $n_3$ which is obtained by dimensionally reducing
the higher-dimensional Einstein–Hilbert action, $\sqrt{g}R$. Keeping in mind that $\sqrt{g}g^{\mu\nu} \propto a^2$
and that $\psi_n \propto a_{sm}^{1/2}/a^{1/2}$ for the nonzero modes (equation (2.9)), we find that the trilinear
vertex involving $0 \leq r \leq 3$ massive KK modes (and $3 - r$ massless KK modes) has the
following representative estimate:

\[
\mathcal{L}_{\text{int}} \sim \int_{-y_i}^{y_{sm}} dy \sqrt{g} g^{\mu \nu} g^{\alpha \beta} g^{\rho \sigma} \partial_\mu h_{\alpha \beta} \partial_\nu h_{\rho \sigma}
\]

\[
\sim \int_{-y_i}^{y_{sm}} dy e^{-2k|y|} \eta^{\mu \nu} \psi_{n_1}^1(x, y) \partial_\mu \psi_{n_2}^2(x, y) \partial_\nu \psi_{n_3}^3(x, y)
\]

\[
\sim \psi_{n_1}^1(x) \psi_{n_2}^2(x) \psi_{n_3}^3(x) p_2 \cdot p_3 a_{sm}^{r/2} \int_{-y_i}^{y_{sm}} dy e^{-2k|y|} (e^{3k|y|/2})^r
\]

\[
\sim \psi_{n_1}^1(x) \psi_{n_2}^2(x) \psi_{n_3}^3(x) \left( \frac{p_2 \cdot p_3}{k} \right) a_{sm}^{\eta}
\]

where

\[
\eta = 2 - r \quad \text{if } r \geq 2, \quad \text{and} \quad \eta = \frac{r}{2} \quad \text{if } r = 0, 1. \quad (3.3)
\]

Notice for this estimate that since derivatives in the compactified directions are proportional to \(g^{mn}\) rather than \(g^{\mu \nu}\), they suffer from additional suppression by powers of \(a = e^{-ky}\) within the throat. Here \(m, n\) label the internal directions perpendicular to the large \((3+1)\)-dimensional Minkowski space.

Thus a trilinear interaction amongst generic KK modes \((r = 3)\), even those with very large \(n\), is proportional to \(1/(a_{sm} k) \sim 1/M_{sm}\), and so is only suppressed by inverse powers of the low scale. Similarly, \(r = 2\) processes involving two massive KK modes \(B\) and \(B'\), and one massless bulk mode \(Z M\)—such as the reaction \(B \rightarrow B' + Z M\)—are \(\propto 1/k \sim 1/M_p\) and so have the strength of 4D gravity inasmuch as they are Planck suppressed. The same is also true of the \(r = 0\) couplings which purely couple the zero modes amongst themselves\(^8\).

Finally, those couplings involving only a single low-lying massive mode and two zero modes \((r = 1)\)—such as for \(B \rightarrow Z M + Z M'\)—are proportional to \(a_{sm}^{1/2}/k \sim (M_{sm}/M_p)^{3/2}\) and so are even weaker than Planck-suppressed ones.

Similar estimates may also be made for the couplings of the generic and the massless KK modes to degrees of freedom on a brane sitting deep within the most strongly warped throat. Using the expressions \(\phi_0(y_{sm}) \sim 1\) and \(\phi_{n}(y_{sm}) \sim 1/a_{sm}\) for massless and massive KK modes respectively, this gives

\[
\mathcal{L}_i = M_p^{-1} \left( h_{\mu \nu}^{(0)} \phi_0(y_{sm}) + \sum_n h_{\mu \nu}^{(n)} \phi_{n}(y_{sm}) \right) T_{sm}^{\mu \nu}
\]

\[
\sim \left( h_{\mu \nu}^{(0)} M_p + \sum_n h_{\mu \nu}^{(n)} M_{sm} \right) T_{sm}^{\mu \nu}. \quad (3.4)
\]

We see here the standard Planck-suppressed couplings of the massless modes (such as the graviton) as compared with the \(O(1/M_{sm})\) couplings of the massive KK modes.

The picture which emerges is one for which the energy released by brane–antibrane annihilation ends up distributed among the massive KK modes of the massless string states. Because the wavefunctions of these modes tend to pile up at the tip of the most warped (SM) throat, their couplings amongst themselves—and their couplings with states

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\(^8\) The appendix shows that this agrees with the size of the couplings found in the effective 4D supergravity Lagrangian which describes the zero-mode and brane couplings.
localized on branes in this throat—are set by the low scale $M_{\text{sm}}$ rather than by $M_p$. Furthermore, because the $\mathcal{O}(1/M_{\text{sm}})$ couplings to the massless modes on the SM branes are much stronger than the Planck-suppressed couplings to the massless bulk modes, we see that the ultimate decay of these massive KK modes is likely to be into brane states. Thus the final production of massless KK zero modes is likely to be practically nil\(^9\). Although we make the argument here for gravitons, the same warp counting applies equally well to the other fields describing the massless closed-string sector.

In summary, we see that strong warping can provide a mechanism for dumping most of the energy released by the decay of the unstable brane–antibrane system into massless modes localized on branes localized at the most strongly warped throat, regardless of whether the initial brane–antibrane annihilation is located in this throat. It does so because the primary daughter states produced by the decaying brane–antibrane system are expected to be very energetic closed strings, which in turn rapidly decay into massive KK modes of the massless string levels. The strong warping then generically channels the decay energy into massless modes which are localized within the most strongly warped throats, rather than into massless bulk modes.

4. Warped reheating

From the previous section we see that a strongly warped internal geometry provides a natural mechanism for channelling the energy released by brane–antibrane annihilation onto those branes which sit deep within the most strongly warped throats, rather than dumping it into massless bulk states like the graviton. Surprisingly, this is true even if the brane–antibrane annihilation is localized within a region which is well separated within the extra dimensions from the most warped throat.

A naive estimate for the expected reheat temperature is obtained by following standard arguments \[33\]. Neglecting dimensionless factors like string couplings and $2\pi$, the total 4D energy density released by the annihilation is of order $\rho_0 \sim \tau_0 \sim M_i^4$, and this energy is dumped into massless string states with a rate set by the local string scale, $\Gamma_s \sim M_i$. Estimating the rate for the resulting KK modes to decay into massless brane states to be also of this order, $\Gamma_{\text{KK}} \sim M_i$, makes the total rate for producing massless states $\Gamma_{\text{tot}} \sim M_i$. During this decay process the initial energy density dilutes due to the expansion of the four large dimensions, according to $\rho(a) = \rho_0(a_0/a)^n$, with the power $n$ depending on the equation of state of the universe during this time. Since the massless modes cannot come into thermal equilibrium until the Hubble scale, $H = \dot{a}/a \sim \rho^{1/2}/M_p \sim H_0(a_0/a)^{n/2}$, falls below $\Gamma_{\text{tot}}$, it cannot occur earlier than $a = a_{\text{eq}}$, where $a_{\text{eq}}$ is defined by the condition $H(a_{\text{eq}}) = \Gamma_{\text{tot}}$ and so $(a_0/a_{\text{eq}})^{n/2} \sim \Gamma_{\text{tot}}/H_0 \sim \Gamma_{\text{tot}} M_p/\rho_0^{1/2}$.

Assuming that equilibrium happens immediately, once $H$ reaches $\Gamma_{\text{tot}}$, the resulting equilibrium temperature is estimated by writing $\rho(a_{\text{eq}}) \sim T^{4}_{\text{RH}}$, and so $T_{\text{RH}} \sim \rho_0^{1/4} (a_0/a_{\text{eq}})^{n/4} \sim (\Gamma_{\text{tot}} M_p)^{1/2}$. Once order-unity factors are included this leads to

$$T_{\text{RH}} \sim 0.1 (\Gamma_{\text{tot}} M_p)^{1/2} \sim 0.1 a_i^{1/2} M_p \sim 10^{-3} M_p,$$

\[4.1\]

\(^9\) Detailed calculations of these rates would be worthwhile, in particular to verify the extent to which resonant enhancement of various bulk modes along the lines found in \[32\] can overwhelm the exponential growth that we track here.
where the last estimate uses $a_i \sim 10^{-4}$. Notice that this is independent of the equation of state parameter $n$, as long as the universal energy dilution is well described by a power law. If reliable, it would also be high enough to adequately reheat the post-inflationary universe to avoid potential problems to which a low reheat temperature can give rise. This can only be regarded as illustrative, because the result is larger than both $M_i$ and $M_{\text{sm}}$; because it is larger than the string scale in the throats it invalidates the 4D field-theoretic calculation on which it is based. A more careful calculation must instead be based on a higher-dimensional, string-theoretic estimate of the energy loss, which goes beyond the scope of this paper.

In conventional inflation models, such a high reheating temperature would be in conflict with the gravitino bound (overproduction of gravitinos, whose late decays disrupt big bang nucleosynthesis). It is interesting in this regard that the KKLT scenario gives a very large gravitino mass, around $m_{3/2} = 6 \times 10^{10}$ GeV [19], which is so large that there is effectively no upper limit on the reheating temperature (see for instance [41]). The disadvantage of such a large gravitino mass is that supersymmetry is broken at too high a scale to explain the weak scale of the SM. If SUSY is this badly broken, one possibility for explaining the weak hierarchy is that the large landscape of string vacua provides a finely tuned Higgs mass, as well as cosmological constant, as has been suggested in [42]. If this were the case, then the degree of warping in the SM model brane would not be crucial for determining the TeV scale, and the existence of an extra throat to contain the SM model brane would be unnecessary. However, given the large number of 3-cycles in a typical Calabi–Yau manifold, each of which can carry nontrivial fluxes, the existence of many throats should be quite generic, and it would not be surprising to find the SM brane in a different throat from the inflationary one.

5. Conclusions

We have argued that for brane–antibrane inflation in strongly warped extra-dimensional vacua—such as have been considered in detail for type IIB string models—there is a natural mechanism which channels the released energy into reheating the Standard Model degrees of freedom. This is because a sizable fraction of the false vacuum energy of the brane–antibrane system naturally ends up being deposited into massless modes on branes which are localized inside the most strongly warped throats, rather than being dumped into massless bulk-state modes.

This process relies on what is known about brane–antibrane annihilation in flat space, where it is believed that the annihilation energy dominantly produces very massive closed-string states, which then quickly themselves decay to produce massive KK modes for massless string states. What is important for our purposes is that the wavefunctions for all of the massive KK modes of this type are typically exponentially enhanced at the bottom of warped throats, while those for the massless KK bulk modes are not. This enhancement arises because the energies of these states are minimized if their probabilities are greatest in the most highly warped regions. This peaking is crucial because it acts to suppress the couplings of the massive KK modes to the massless bulk states, while enhancing their couplings to brane modes in the most warped throats.

The upshot is that most of the energy density of the brane–antibrane system ends up as radiation on the SM brane, regardless of whether or not this brane is physically
separated from the place where the brane–antibrane annihilation actually occurs. Our arguments are generic, and do not require us to make any unusual assumptions about or impose any restrictions on the parameters. From this point of view, efficiently reheating the SM brane after brane–antibrane inflation is unavoidable, as long as there are no other hidden branes lying in even deeper throats than the SM one. This observation is all the more interesting given the attention which multiple-throat inflationary models are now receiving, both due to the better understanding which they permit for the relation between the inflationary scale and those of low-energy particle physics, and to the prospects they raise for producing long-lived cosmic string networks with potentially observable consequences.

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Appendix: The 4D view

In this appendix we compute the low-energy couplings amongst the bulk zero modes and brane modes in the effective 4D supergravity obtained after modulus stabilization à la KKLT [10]. Besides checking the scaling of the kinetic terms obtained by dimensionally reducing the Einstein–Hilbert action, this also allows the study of the couplings in the scalar potential which arise from modulus stabilization and so are more difficult to analyse from a semiclassical, higher-dimensional point of view.

To this end imagine integrating out all of the extra-dimensional physics to obtain the low-energy effective 4D supergravity for a type IIB GKP vacuum having only the mandatory volume modulus (and its supersymmetric friends) plus various low-energy brane modes (such as those describing the motion of various $D3$ branes). The terms in this supergravity involving up to two derivatives are completely described once the Kähler function, $K$, superpotential, $W$, and gauge kinetic function, $f_{ab}$, are specified.

Denoting the bulk-modulus supermultiplet by $T$ and the brane multiplets by $\phi^I$, we use the Kähler potential [34, 10, 11, 35]

$$K = -3 \log[r],$$

where $r = T + T^* + k(\phi, \phi^*)$. For instance, if $\phi^I$ denotes the position of single brane, then $k$ is the Kähler potential for the underlying 6D manifold. This implies that the scalar kinetic terms are governed by the following Kähler metric in field space:

$$K_{TT^*} = \frac{3}{r^2}, \quad K_{IT^*} = \frac{3k_I}{r^2} \quad \text{and} \quad K_{IJ^*} = \frac{3}{r^2} \left[ k_I k_J^* - r k_{IJ^*} \right],$$

(A.2)
with inverse
\[ K^{-T} = \frac{r}{3} \left[ r - k^{L} N k_{L} * k_{N} \right], \quad K^{J-T} = \frac{r}{3} k^{J-L} k_{L} \quad \text{and} \quad K^{J-I} = -\frac{r}{3} k^{J-I}. \quad (A.3) \]

In the absence of modulus stabilization the superpotential of the effective theory is a constant \[36\], \( W = w_0 \), and the supergravity takes the usual no-scale form \[37\], with vanishing scalar potential. If, however, there are low-energy gauge multiplets associated
with any of the \( D7 \) branes of the model, then their gauge kinetic function is \( f_{ab} = T \delta_{ab} \).
For non-Abelian multiplets of this type gaugino condensation \[38,39\] can generate a nontrivial superpotential, of the form
\[ W = w_0 + A \exp[-a T], \quad (A.4) \]
where \( A \) and \( a \) are calculable constants.

With these choices the Kähler derivatives of the superpotential become
\[ D_T W = W_T - \frac{3W}{r}, \quad \text{and} \quad D_I W = -\frac{3k_I W}{r}, \quad (A.5) \]
and so the supersymmetric scalar potential \[40\] becomes
\[ V = \frac{1}{3r^2} \left[ (r - k^{I-J} k_{I} k_{J}) |W_T|^2 - 3(W^* W_T + W W_T^*) \right]. \quad (A.6) \]
Notice that use of these expressions implicitly requires that we work in the 4D Einstein frame, and so are using 4D Planck units for which \( M_p = O(1) \).

If we specialize to the case of several branes, for which \( \{ \phi^I \} = \{ \phi^I_n \} \), with \( i \) labelling the fields on a given brane and \( n = 1, \ldots, N \) labelling which brane is involved, then we typically have
\[ k(\phi^I, \phi^{I*}) = \sum_n k^{(n)}(\phi^I_n, \phi^{I*}_n). \quad (A.7) \]
In this case the Kähler metric built from \( k \) is block diagonal, with \( k_{n,m} = k_{ij}^{(n)} \delta_{mn} \), and so \( k^{I-J} k_{I} k_{J} = \sum_n k^{(n)}_{I-J} k^{(n)}_I k^{(n)}_J \) and so on.

We may now see how strongly the bulk KK zero modes, \( g_{\mu\nu} \) and \( T \), couple to one another and to the brane modes. Setting \( k = 0 \) in the above shows that the couplings of \( T \) and \( g_{\mu\nu} \) to one another are of order unity, and since our use of the standard 4D supergravity formalism requires us to be in the Einstein frame, this implies these are all of 4D Planck strength (in agreement with our higher-dimensional estimates).

Couplings to the branes are obtained by keeping \( k \) nonzero, and in the event that the branes are located in highly warped regions, we must take \( k^{(n)} = O(a^2_n) \) with \( a_n \ll 1 \) denoting the warp factor at the position of brane \( n \).\(^{10}\) In this case the combination \( k^{(n)}_{I-J} k^{(n)}_I k^{(n)}_J \) is also \( O(a^2_n) \).

Suppose we now expand the functions \( k^{(n)} \) in powers of \( \phi \) and keep only the leading powers:
\[ k^{(n)} \approx a^2_n \sum_i \phi^{i*}_n \phi^i_n. \quad (A.8) \]

\(^{10}\) For instance, this power of \( a_n \) reproduces the dependence on \( a_n \) of the factor \( \sqrt{g} g^{\mu\nu} \) obtained by dimensionally reducing the higher-dimensional kinetic terms.
Then, since the $\phi_n$ kinetic terms are $O(a_n^2)$, we see that the canonically normalized fields are $\chi^i_n = a_n \phi^i_n$. Once this is done the leading couplings to $T$ and $g_{\mu\nu}$ are those which involve those parts of $k^{(n)}$ that are quadratic in $\chi^i_n$, and since these are also of order unity, these couplings are also of Planck strength (again in agreement with our earlier estimates).

Alternatively, consider now those couplings which only involve the brane modes. Working to leading order in $a_n^2$, we see that a term in $k^{(n)}$ of the form $(\chi^i_n)^k$ has a strength which is of order $a_n^{-2-k}$. For instance the case $k = 3$ generates cubic couplings from the kinetic Lagrangian of order $a_n^{-1} \chi \partial \chi \partial \chi$, whose coefficient is of order $(a_n M_p)^{-1} = M_{\text{sm}}^{-1}$. These are larger than Planck-suppressed ones, as expected.

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