Dumping inflaton energy density out of this world

Kari Enqvist,[1] Anupam Mazumdar,[2] and Abdel Pérez-Lorenzana[3]

[1]*Department of Physical Sciences, University of Helsinki, and Helsinki Institute of Physics, P.O. Box 9, FIN-00014 University of Helsinki, Finland*

[2]*CHEP, McGill University, 3600 University Road, Montréal, Québec, H3A 2T8, Canada*

[3]*Departamento de Física, Centro de Investigación y de Estudios Avanzados del I.P.N Apdo. Post. 14-740, 07000, México, D.F., México.*

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We argue that a brane-world with a warped, infinite extra dimension allows for the inflaton to decay into the bulk so that after inflation, the effective dark energy disappears from our brane. This is achieved by shifting away the decay products into the infinity of the 5th dimension. As a consequence, all matter and CMB density perturbations could have their origin in the decay of a MSSM flat direction rather than the inflaton. We also discuss a string theoretical model where reheating after inflation may not affect the observable brane.

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I. INTRODUCTION

Recent observations [1] strongly support a period of primordial inflation. Besides making the universe flat and homogeneous, inflation is the only known dynamical mechanism which can stretch small quantum fluctuations outside the Hubble horizon. These perturbations act as seeds for the large scale structures in the Universe. However, despite of the success of the inflationary paradigm (for a review, see [2]), we know very little about the inflationary sector. The inflaton potential must nevertheless be very flat with a very small self coupling; likewise, its coupling to other fields must be extremely weak. Such small couplings are hard to come by without fine tuning, which renders inflaton to a gauge singlet whose couplings can be adjusted at our will. Given this, the immediate question is, what is the inflaton decaying into?

Eventually the cold inflationary Universe must reheat with the Standard Model (SM) degrees of freedom, or, as the current theoretical prejudice dictates, with the Minimally Supersymmetric Standard Model (MSSM) degrees of freedom. MSSM contains a number of flat directions along which the renormalizable part of the potential vanishes (for a review, see [3]). During inflation the massless fields corresponding to the flat directions receive scale-invariant perturbations. It would be tempting to associate them with reheating and the generation of CMB perturbations [4–7]. This would require that the flat directions will eventually dominate over the radiation produced by the inflaton decay. In other words, the flat direction would have to act as an MSSM curvaton [8–10]. However, it turns out that such domination is not possible unless the inflaton decays into some hidden degrees of freedom, rather than into MSSM radiation [4,5].

In this respect, brane-world scenarios, where the Universe is regarded as a three dimensional hypersurface embedded in a higher dimensional bulk, bring along new, interesting possibilities. The SM degrees of freedom are assumed to be stuck on a brane while gravity is propagating in the entire bulk. The bulk could also have a nontrivial background geometry which allows for the zero mode of graviton to be trapped at the brane location, such as in the case of an anti de-Sitter (adS) bulk in the Randall-Sundrum type models [11,12]. The value of the Newton's constant requires the fundamental scale to be fairly large, $10^{18}\mathrm{GeV} \geq M_s \geq 10^3$ GeV. In these models the inflaton, again treated as a gauge singlet, could either live on the brane or in the entire bulk. There is also the exciting possibility that the inflaton energy density does not reheat the Universe, but gets deposited in the infinite bulk or onto the adS horizon. In this respect the inflaton potential could be treated as a kind of dark energy for the MSSM brane. It has been shown [6] that if the dark energy of the inflaton field can be transferred into the bulk so that it no longer is visible on our brane, even the simplest MSSM flat directions can act as curvatons and provide both the matter and the density perturbations in the Universe.

The brane-world scenarios are supported by the string theory, where there is a natural explanation for the construction of MSSM like brane with a help of coincident stack of Dp branes (p is the number of spatial dimensions along the brane, p is odd (even) in type IIA (IIB) string theories) attached to some orbifold point [13] [1]. The open strings attached to Dp branes act as sources for gauge fields and gravity again propagates in the entire bulk. In this kind of framework the inflaton may be regarded as a moduli, for example, the physical separation between a Dp and antibrane $\bar{D}p$ brane, or the angular separation between $Dp - Dp$ or $Dp - \bar{D}p$ branes (for a review see [15]). Inflation may end by virtue of tachyon condensation when branes approach close to the string scale [16], or with a help of many tachyons as in the case of assisted

[1]For different constructions, see [14].

inflation [17]. It is, however, not guaranteed that inflaton will reheat the Universe with the MS(SM) degrees of freedom. One could rather argue that it is more likely that the inflaton will reheat the bulk.

The purpose of this paper is to construct a brane-world model with a warped bulk so that it is possible to localize the inflaton energy density away from the MS(SM) like branes. We will argue that the inflaton energy can be dumped away so that after inflation there is effectively no energy density other than that of the excited MSSM flat directions on our brane. For our purposes, the inflaton potential is a form of dark energy which is only responsible for making the Universe, parallel to the brane directions, large, homogeneous and isotropic.

The paper is organized as follows. In Section II we recapitulate some known results of the brane-world models and discuss infinite extra dimensions. In Section III we present a brane-world where the dark energy can be dumped into the bulk instead of being dumped onto the brane and estimate the escape rate of the inflaton decay products from the brane. In Section IV we motivate the model by considerations relating to string theoretical inflation. In Section V we give our conclusions and discuss how the present results can be combined with curvaton-like scenarios involving MSSM flat directions to yield all matter and an adiabatic, scale-invariant spectrum of perturbations.

II. INFINITE EXTRA DIMENSION AND KK DECOMPOSITION

In this section we briefly recapitulate some of the already known results of the brane-world models. We start with the simplest scenario assuming that there is a three dimensional hypersurface, called the brane, which carries MSSM degrees of freedom. The MSSM brane is embedded in a 5 dimensional space (the bulk) with a nonfactorizable metric [11,12] (for a nice review, see [18])

$$ds^2 = a^2(z)\eta_{\mu\nu}dx^\mu dx^\nu - dz^2, \tag{2.1}$$

where $\eta_{\mu\nu}$ is the four dimensional Minkowski metric. We take the extra dimension to be infinite. The brane is located at $z = 0$. The total action for gravity is given by

$$S_g = -\frac{1}{16\pi G_5}\int d^4x dz\sqrt{g^{(5)}}R^{(5)} - \Lambda\int d^4x dz\sqrt{g^{(5)}} - \sigma\int d^4x\sqrt{g^{(4)}}, \tag{2.2}$$

where Λ is a bulk negative cosmological constant which is related to the brane tension σ by the fine tuning relationship

$$\Lambda = -\frac{4\pi}{3}G_5\sigma^2. \tag{2.3}$$

The warp factor present in the metric has the form $a(z) = e^{-k|z|}$, with

$$k = \frac{4\pi}{3}G_5\sigma. \tag{2.4}$$

Here G_5 stands for the true gravitational coupling constant of the five dimensional theory, which defines the fundamental gravity scale as $M_* = (8\pi G_5)^{-1/3}$. The effective four dimensional Newton's constant $M_P = (8\pi G_4)^{-1/2} = 2.4\times 10^{18}$ GeV is given by

$$G_4 = kG_5. \tag{2.5}$$

This setup is inspired by adS/CFT correspondence [19], where the entire bulk has adS geometry. It has an interesting feature along the z axis in that the fifth dimension has a horizon at $z = \infty$. A particle that escapes from the brane and moves along a geodesic travels from $z = 0$ to $z = \infty$ in a finite proper time $\tau = \pi/2k$. In other words, $z = \infty$ is a particle horizon.

Another interesting property is that the spectrum of the KK gravitons has a localized massless mode around the brane, identified as the four dimensional graviton, plus a continuum of modes with wave functions which are oscillatory at large z,

$$h_m(z) = \text{const}\times\sin\left(\frac{m}{k}e^{kz} + \phi_m\right), \tag{2.6}$$

while near the brane location, at $z = 0$, the wave function is suppressed:

$$h_m(z=0) = \text{const}\times\sqrt{\frac{m}{k}}. \tag{2.7}$$

By summing over all the KK modes, including the zero mode of the graviton, one finds a modification of the gravitational interaction between two point particles located on the brane. If their masses are m_1 and m_2 and the particle separation is r, the potential has the form

$$V(r) = -\frac{G_{(4)}m_1m_2}{r}\left(1 + \frac{\text{const}}{k^2r^2}\right). \tag{2.8}$$

For distances $r \gg 1/k$ the correction to the Newtonian gravity is negligible small. Current experiments have tested the validity of Newton's law down to submillimeter distances, which implies that $k > 10^{-3}$ eV [20]. In this paper we will mostly assume that k and M_*, are close to the four dimensional Planck scale.

It is clear that with an infinite fifth dimension the KK modes of other bulk fields also possess a continuum of modes [21], and may also have a quasilocalized mode [22]. Consider for instance the case of a scalar bulk field on the background metric Eq. (2.1). The action is then given by

$$S_\chi = \int d^4x dz\sqrt{g^{(5)}}\left(\frac{1}{2}g^{ab}\partial_a\chi\partial_b\chi - \frac{1}{2}\mu^2\chi^2\right). \tag{2.9}$$

The corresponding KK wave function is then defined as a solution to the field equation

$$[-\partial_z^2 + 4k\,sgn(z)\partial_z + \mu^2 - m^2/a^2(z)]\chi(z;m) = 0, \tag{2.10}$$

where $m^2 = p^\mu p_\mu$ defines both the four dimensional mass and the KK level. Equation (2.10) is supplemented by the boundary condition on the brane $\partial_z \chi(z = 0; m) = 0$, and the normalization condition $\int dz a^2(z) \times \chi(z; m)\chi(z; m') = k\delta(m - m')$. The general solution to Eq. (2.10) is given in terms of Bessel functions of index $\nu = \sqrt{4 + \mu^2/k^2}$, and can be written as [21,22]

$$\chi(z; m) = \frac{1}{N(m)a^2(z)}\left[J_\nu\left(\frac{m}{ka(z)}\right) + A(m)Y_\nu\left(\frac{m}{ka(z)}\right)\right], \tag{2.11}$$

where the normalization factor, N, and the coefficient A are functions of the continuous KK index m. One finds that [21,22] $N(m) = \sqrt{1 + A^2(m)}/\sqrt{m/2}$ with

$$A(m) = -\frac{2J_\nu(m/k) + (m/k)J'_\nu(m/k)}{2Y_\nu(m/k) + (m/k)Y'_\nu(m/k)}. \tag{2.12}$$

For $\mu, m \ll k$ one can approximate the last expression by taking $\nu = 2$ and show that

$$A(m) \approx \frac{\pi}{4}\left(\frac{m}{k}\right)^2. \tag{2.13}$$

Thus, the KK wave function evaluated at the brane is just

$$\chi(0, m) \approx \sqrt{\frac{m}{2}}. \tag{2.14}$$

Note that the functional behavior of the above expression is the same as for the graviton case in Eq. (2.7).

As noted in Ref. [22], this system has in general a resonance around the brane, i.e., there is a quasilocalized mode of nonzero mass; that is, there is no truly bound state in the spectrum. This can be visualized in a simple way: the continuum of modes is determined by the asymptotic form of the field equation at large z, where the mass term μ^2 is negligible compared with $m^2/a^2(z)$, which shows that the spectrum does start at $m = 0$, independently of μ, but there are no bound states within the continuum. By exploring the KK modes one can show that there is a mode that actually has a complex eigenvalue $m = m_0 + i\Gamma$ [22]. Thus, this mode can be considered as a quasilocalized metastable state for which Γ gives the escape rate from the brane into the extra dimension towards infinity. For $\mu \ll k$ one finds $m_0 = \mu/\sqrt{2}$ and

$$\Gamma = \frac{\pi}{16}\left(\frac{m_0}{k}\right)^2 m_0. \tag{2.15}$$

In the following section we will use some of these results when considering the decay life time of the inflaton.

III. INFLATON DECAY INTO THE BULK IN A WARPED BRANE-WORLD SCENARIO

Let us now discuss how after inflation the inflaton may disappear from the brane and leave behind an (almost) empty brane, with a minor impact on the later cosmological evolution. This is a radical point of view that, however, can easily be accommodated within the context of infinite extra dimension models [23]. To be more specific, let us assume that the inflaton is a true 4D brane field, with a homogeneous distribution that dominates the energy density at the early Universe on the observable brane and gives rise to a period of inflation. Then the Friedman equation has a quadratic dependence on the brane density ρ [24] so that the Hubble expansion rate, $H = \dot{R}/R$, where R is the scale factor of the four dimensional Universe is given by,

$$H^2 = \left(\frac{\dot{R}}{R}\right)^2 = \frac{1}{3M_P^2}\rho\left(1 + \frac{\rho}{2\sigma}\right) + \frac{\sigma}{2M_P^2}\left(\frac{r_h}{R}\right)^4, \tag{3.1}$$

and the last term which accounts for the dark energy emitted from the brane into forming a black hole at the adS horizon, where r_h can be interpreted as the black hole horizon [25], we will discuss this issue later on. The key point to note is that the standard Hubble relationship, $H = \sqrt{\rho/3M_P}$, follows only for small densities compared to the brane tension, $\rho \ll \sigma$ [24,26], and the contribution from the last term is sub dominant as long as the scale factor on the brane position, $R \gg r_h$, is greater than the black hole horizon [25].

Once inflation comes to an end, the inflaton will decay, but instead of reheating the brane degrees of freedom, we now assume that it couples to the bulk fields alone, and decays into bulk degrees of freedom. This may happen, for instance, if the inflaton and the bulk fields carry some global quantum number while the brane degrees of freedom do not. All the inflaton energy would be radiated into the empty bulk after the end of inflation in the form of KK modes. These bulk modes carry momentum along the fifth dimension, so that they would simply fly away into the empty bulk, towards infinity, taking the inflaton energy away from the brane. The energy density of the inflaton will be gradually dumped into the bulk before becoming vanishingly small. A small fraction, however, may act as a dark energy on the brane. We will comment on this below.

It is interesting that whether the inflaton density is larger than brane tension or not becomes irrelevant for the purposes of the present discussion. Inflation could well take place in the non standard regime of the theory, without leaving any visible trace on the subsequent thermal evolution of the Universe [27].

The inflaton decaying into the bulk can be extremely efficient. To demonstrate this let us consider the coupling of the inflaton to a bulk scalar field φ, which in the complete 5D theory can be written as

$$\sqrt{g(z)}h\phi(x)\varphi(x, z)\varphi(x, z)\delta(z), \tag{3.2}$$

where h is the corresponding coupling constant. After introducing the KK decomposition of the bulk field and integrating out the extra dimension one gets the effective coupling of the inflaton to the KK modes as

$$h[\chi(0, m)\chi(0, m')]\phi(x)\varphi_m(x)\varphi_{m'}(x); \tag{3.3}$$

where $\chi(0, m)$ are the z dependent wave functions of a KK modes of mass m, given in Eq. (2.11), evaluated at the brane position. The KK mass dependence of the effective couplings indicate that the inflaton would preferably decay into the heavy modes, i.e., those with the largest momentum along the fifth dimension. This scenario is similar to the one discussed in Ref. [28] for the cooling down of a hot brane by graviton emission, although there the KK gravitons were assumed to be thermally produced.

If the inflaton decays into the continuum of KK modes with masses smaller than m_ϕ, it is straightforward to estimate the total decay rate as

$$\Gamma_\phi = \int_0^{m_\phi}\int_0^{\sqrt{m_\phi^2 - m^2}} \frac{dm}{k}\frac{dm'}{k} h^2 \frac{[\chi(0, m)\chi(0, m')]^2}{m_\phi} \approx \frac{h^2}{32}\left(\frac{m_\phi}{k}\right)^2 m_\phi, \tag{3.4}$$

where the RHS has been estimated in the limit where $\mu, m_\phi \ll k$ using Eq. (2.14). Since the inflaton is heavy, say, about $10^{16} - 10^{13}$ GeV scale, whereas k is close to M_P, the suppression on the decay rate is not large. For smaller inflaton mass the decay rate is suppressed as it is in the standard case. For a wide range of the inflaton mass the inflaton may release all its energy into the bulk fields very efficiently.

Let us now discuss what happens to the energy that has been injected into the bulk. As already mentioned above, since KK modes have a fifth momentum, they will travel away from the brane, moving towards infinity in a finite proper time, $\tau = (\pi/2k)$, see [29], hence $z = \infty$ is indeed the particle horizon. The boundary condition at the horizon are imposed such that nothing comes in from behind the horizon. Further note that the time traverse from the brane to the adS horizon is swifter than the inflaton decay rate as long as $k \gg h^{2/3}(\pi/64)^{1/3} m_\phi$.

The four dimensional Poincaré invariance ensures that the coordinate four momentum p_μ coincides with the physical momentum on the brane, but from the point of view of an observer, away from the brane, $z \neq 0$, the four dimensional momentum gets blue shifted, this means that the modes which are softer on the brane get harder away from the brane. This leads to a conjecture that a black hole might form at the horizon of an adS [25,26].

As the original energy density of the inflaton field is moved away from the brane into the extra dimension, only the tail of the density distribution would be felt by the brane. In a similar vein as in [25,26], we imagine that the inflaton energy emitted from the brane would eventually collapse to a black hole at the end of the space.

As long as the four dimensional scale factor $R \gg r_h$, the dark energy contribution to the Hubble expansion rate is negligible, see Eq. (3.1). But how small r_h/R could be such that the contribution to the dark energy is negligible? Even if we imagine that the entire inflaton energy density is dumped into a black hole formation, then the maximum horizon size would be that corresponding to the inflaton energy. From the brane observer point of view, this should be $r_h \sim [V(\phi)]^{-1/4}$. For typical values, $V(\phi)^{-1/4} \sim 10^{-16}(\mathrm{GeV})^{-1} \sim 10^{-32}$ m, which is negligible compared to present size of the Hubble horizon, even this is smaller once the Universe has inflated and eventually reheated due to the MSSM Higgses, where the reheat temperature of the Universe was found to be $\sim 10^9$ GeV, see [6]. The physical point is that the size of the Hubble horizon is manifold larger compared to the horizon of the black hole.

In fact in our case we can estimate roughly how much is the contribution from the last term due to the decay of the inflaton energy density. The rate of energy loss can be estimated by

$$\frac{\Delta\dot\rho}{\rho} = \mathcal{O}(1)\Gamma_\phi. \tag{3.5}$$

The order one coefficient takes into account the degrees of freedom the inflaton is decaying into, and note that $\Gamma_\phi \geq H$ at the time of decaying. At sufficient late times $1/H \gg 1/k$ or $\rho \ll \max\{V(\phi), \sigma\}$, the total contribution of the dark energy will be roughly given by the loss of energy density from the brane, e.g., $\Delta\dot\rho + \Delta\dot\rho_d = 0$, [25],

$$\Omega_d = \frac{\rho_d}{\rho_d + \rho} \approx \int_{\tau_1}^{\infty} d\tau\left(-\frac{\Delta\dot\rho}{\rho}\right). \tag{3.6}$$

In our case we can estimate the integral, which is roughly given by

$$\Omega_d \approx \frac{h^2\pi}{64}\left(\frac{m_\phi}{k}\right)^3. \tag{3.7}$$

For large values of $k \gg m_\phi$, not only the decay rate of the inflaton to the bulk is efficient, but also the dark energy contribution is miniscule, and for all practical purposes we can neglect the dark energy. To summarize the inflaton energy density, which is moved away into the bulk poses no threat to the later evolution of the Universe, which is after inflation solely governed by the MSSM Higgs condensate, [6].

There are some alternatives to the picture presented here. For example, one could use a bulk field to provide for the inflaton as its quasilocalized (resonant) mode as described at the end of previous section. It is interesting to note that the width of such a mode, given by Eq. (2.15), has the same functional form as Eq. (3.4), so that its escape rate from the brane would be as efficient as the decay of the inflaton in the our present model.

IV. A STRINGY MOTIVATION

In string theories inflationary cosmologies have often been discussed (see, e.g., [15]) by making use of the fact that Dp and $\bar{D}p$ branes can attract each other by virtue of the spacetime supersymmetry breaking [30]. If the attractive potential is sufficiently flat then it can give rise to a slow roll inflation [31]. One of the constraints is that the brane separation and hence the bulk has to be larger than the inverse of the compactification scale, which is indeed a problem [32] and can be viewed as an initial condition problem requiring some kind of a bulk inflation [33] before brane inflation. Such early inflation could be triggered, e.g., by gas of branes [34].

An important issue is the stabilization of the volume, the dilaton and the moduli. In the context of a warped background one typically finds a Klebanov-Strassler kind of solution for the background metric [35]. One can think of this geometry as a stringy generalization of the Randall-Sundrum model, where there is an adS throat (or a conical singular region) where the infrared brane can be regularized by the infrared geometry. The Klebanov-Strassler solution gives rise to a noncompact geometry where the radial internal coordinate can be thought of as z in the 5 dimensional language of Randall-Sundrum warped geometry. The Planck brane is at the ultraviolet region (at large r), while the infrared brane is stuck near the adS throat (at small r), where r parametrizes the internal radial coordinate and the metric is of the form

$$ds^2 = h^{-1/2}(r)dx_\mu dx^\mu + h^{1/2}(r)(dr^2 + r^2 ds_{(5)}). \quad (4.1)$$

A particular realization of such compactification on a Calabi-Yau manifold gives rise to the stabilization of the complex structure moduli [36–38].

In this paper we will not be able to discuss a detailed inflationary model within such stringy framework. We will merely present some plausibility arguments supporting our case as discussed in Section III. To this end, let us consider a scenario which consists of a stack of branes mimicking MSSM gauge group (stack of $D3$ branes and $D7 - \bar{D}7$ branes (the latter ones required for an exact cancellation of the tadpoles at the singularity). In addition, there should be other branes and antibranes to drive inflation along the three spatial directions. The MSSM branes are embedded near the ultraviolet part of the geometry, at large r, far from the adS throat, while the branes, which give rise to inflation, are moving towards the region where the adS throat is located, near the infrared part of the geometry, at small r.

To be concrete, let us assume that all the moduli are stabilized like in [38]. A set of $\bar{D}3$ branes are stuck near the infrared part of the bulk, and inflation occurs because of the slow motion of a $D3$ brane approaching from the ultraviolet regime. The potential of $D3 - \bar{D}3$ is also felt by the MSSM branes and is described by a four dimensional effective field theory on the world volume. The presence of $\bar{D}3$ branes break supersymmetry and give rise to a metastable positive vacuum energy state, which is also being felt along the three spatial dimensions. This acts as a source for inflation [38,39] on the MSSM branes. We have depicted this framework in Fig. 1.

We assume that enough inflation can be obtained along the three spatial dimensions of the MSSM branes. In the brane antibrane scenario inflation ends when the separation becomes close to the string scale, whence the open string tachyon on the world volume condenses [16], resulting in an annihilation of the pair of branes. Similar situation could arise in our case. The rolling tachyon couples to the gauge fields living on the brane through covariant derivatives. The annihilation of branes ultimately gives rise to a long excited closed string along the inflated directions [40]. The long closed string decays very late into lighter closed string modes. However, the important point to note here is that only the bulk degrees of freedom near the adS throat are excited. This is so because the branching ratio of the closed string decaying into the bulk is still greater compared to decay into brane degrees of freedom by virtue of phase space arguments. These modes are actually trapped near the infrared regime which is energetically more favorable due to low energy configuration.

This situation is indeed similar to the 5 dimensional adS Randall-Sundrum model described in Section III. There the bulk quanta were dumped towards the adS horizon to form a black hole. Formation of a long closed string after the end of inflation also has a counterpart in field theory. If the inflaton has a global $U(1)$ charge, it may not decay completely but rather fragments into lumps known as Q-balls [41].

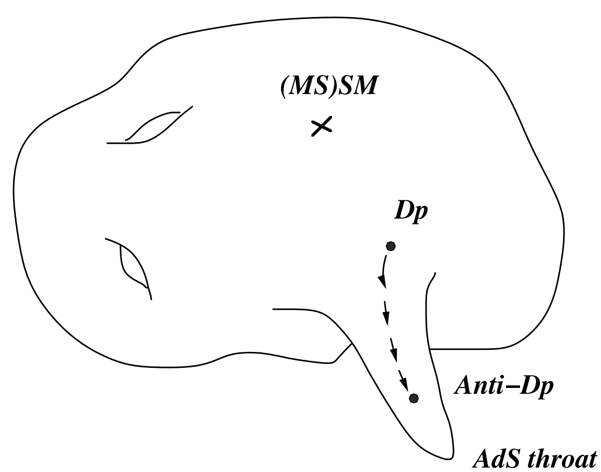

FIG. 1. An illustration of a manifold which has singular points with nontrivial fluxes, and which yields the adS geometry. The cross denotes the point in moduli space where MSSM branes are fixed, while there is a $\bar{D}3$ situated close to the warped geometry near adS throat, and the $D3$ brane is attracted towards the $\bar{D}3$ brane, thereby giving rise to inflation along three spatial directions, which we assume to be parallel to MSSM branes. Reheating occurs near the adS throat. The excited modes from reheating are trapped near the throat.

The main point here is that the MSSM branes are not directly reheated from the decay of the long closed string. They will certainly feel the resulting effective dark energy, but due to the warped metric the dark energy contribution would be small.

V. DISCUSSION

During inflation, massless MSSM fields, or the (almost) flat directions corresponding to certain combinations of squarks, sleptons and Higgses, will be subject to fluctuations and form condensates. Like the ordinary inflaton, the condensates will receive scale-invariant spatial perturbations. Once the inflaton energy has disappeared into the bulk, the potential terms along some flat direction will eventually start to dominate the energy density of the MSSM brane. A potential arises because of nonrenormalizable interactions and the soft supersymmetry breaking mass terms (for a review, see [3]). (In the cosmological context the Higgs coupling $\mu H_u H_d$ does not spoil flatness as μ is much smaller than the relevant field amplitudes).

Once the condensate decays, it will reheat the universe with MSSM degrees of freedom and imprint on the MSSM gas the inflationary perturbations. The flat direction condensate acts as an MSSM curvaton [8–10]. The simplest possibility is the flat direction that consists of the Higgses H_u and H_d, which has been discussed in detail in [6]. The reheat temperature can be estimated to be less than 10^9 GeV [2]. The amplitude of the fluctuations along the $H_u H_d$ direction can match the observed density perturbations in the CMB radiation and the spectrum with a spectral tilt very close to 1, with some weak dependence on the Higgs potential [6].

As an interesting observation we finally point out that in our case we will not have any isocurvature fluctuations, because the entire SM degrees of freedom are created solely by the decay of the Higgses and not by the inflaton. Although our model provides startling distinction from multifield inflationary models, but at present it will be hard to distinguish the predictions of our model from a single field inflationary model.

We have argued that a brane cosmology with a warped, infinite extra dimension allows for the inflaton to decay into the bulk. The inflaton decays efficiently into a continuum of Kaluza Klein modes which carry nonzero momentum along the extra dimension and move away from our four dimensional world, taking inflaton energy with them. In effect, the effective dark energy present on the MSSM brane will disappear into the bulk and be hidden behind the particle horizon at the infinity of the 5th dimension.

Although we have sketched a possible inflationary scenario involving $\bar{D}3$ and $D3$ branes that annihilate at an adS throat, a string theoretical model for reheating remains a challenge. Nevertheless, it seems that at least within brane-world cosmologies it is possible to have the inflaton decay products disappearing into the bulk so that all matter could have its origin in the decay of the MSSM condensate rather than in the inflaton energy density.

[2]In Ref. [41], we discussed the perturbative decay of $H_u H_d$ into the SM fermions and baryons through Yukawa couplings. One might still wonder on other means of reheating, such as gravitational reheating [42] through nonconformal coupling of matter to gravity. However such a nonconformal coupling is not constrained from a fundamental theory in our case, and therefore does not give rise to any concrete prediction on the reheat temperature.

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