Prospects of inflation with perturbed throat geometry

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A R T I C L E   I N F O

A B S T R A C T

We study brane inflation in a warped deformed conifold background that includes general possible corrections to the throat geometry sourced by coupling to the bulk of a compact Calabi–Yau space. We focus specifically, on the perturbation by chiral operator of dimension 3/2 in the CFT. We find that the effective potential in this case can give rise to required number of e-foldings and the spectral index n_S consistent with observation. The tensor to scalar ratio of perturbations is generally very low in this scenario. The COBE normalization, however, poses certain difficulties which can be circumvented provided model parameters are properly fine tuned. We find the numerical values of parameters which can give rise to enough inflation, observationally consistent values of density perturbations, scalar to tensor ratio of perturbations and the spectral index n_S.

1. Introduction

It is well known that the standard model for Big bang cosmology is plagued with some intrinsic problems like horizon and flatness problems. Such problems could be cured if one postulates that the early universe went through a brief period of accelerated expansion, otherwise called the cosmological inflation [1]. The inflationary scenario not only explains the large-scale homogeneity of our universe but also proposes a mechanism, of quantum nature, to generate the primordial inhomogeneities which is the seed for understanding the structure formation in the universe. Such inhomogeities have been observed as anisotropies in the temperature of the cosmic microwave background, thus the theoretical ingredients of inflation are subject to observational constraints. Moreover, the paradigm of inflation has stood the test of theoretical and observational challenges in the past two decades [2,3]. However, in spite of its cosmological successes, it still remains a paradigm in search of a viable theoretical model that can be embedded in a fundamental theory of gravity. In this context, enormous amount of efforts are underway to derive inflationary models from string theory, a consistent quantum field theory around the Planck's scale. Progress in the understanding of the nonperturbative dualities in string theory, to be precise, discovery of D-branes, has further provided an important framework to build and test inflationary models of cosmology.

In past few years, many inflationary models have been constructed in the context of D-brane cosmology. It includes inflation due to tachyon condensation on a non-BPS brane, inflation due to the motion of a D3-brane towards an anti-D3-brane [4–7], inflation due to geometric tachyon arising from the motion of a probe brane in the background of a stack of either NS5-branes or the dual D5-branes [8]. However, these models are based on effective field theory and assume an underlying mechanism for the stabilization of various moduli fields. Thus, they do not take into account the details of compactification and the effects of moduli stabilization and hence any predictions from these models are questionable.

Progress in search of a realistic inflationary model in string theory was made when it was learnt that background fluxes can stabilize all the complex structure moduli fields. In fact it has been shown in Ref. [9] that the fluxes in a warped compactification, using a Klebanov–Strassler (KS) throat [10], can stabilize the axio-dilaton and the complex structure moduli of type IIB string theory compactified on an orientifold of a Calabi–Yau threefold. Further important progress was achieved when it was shown in Ref. [11] that the Kähler moduli fields also can be stabilized by a combination of fluxes and nonperturbative effects. The nonperturbative effects, in this context, arise, via gauge dynamics of either an Euclidean D3-brane or from a stack of D7-branes wrapping super-symmetrically a four cycle in the warped throat. The warped volume of the four cycle controls the magnitude of the nonperturbative effect since it affects the gauge coupling on the D7-branes wrapping this four cycle.

Armed with these results, an inflationary model [12] has been built taking into account of the compactification data (see also Ref. [13]). The inflaton potential is obtained by performing string theoretic computations involving the details of the compactification scheme. In this setup inflation is realized by the motion of a D3-brane towards a distant static anti-D3-brane, placed at the tip

of the throat and the radial separation between the two is considered to be the inflaton field. The effect of the moduli stabilization resulted in a mass term to the inflaton field which is computed in [12]. It turned out, unfortunately that the mass is large and of the order of Hubble parameter and hence spoils the inflation. To avoid this disappointment (see, however, [14]), one needs to search for various sources of correction to the inflaton potential. For example, in Ref. [15], the embedding of the D7-branes as given in [16] was considered. Assuming that at least one of the four-cycles carrying the nonperturbative effects descend down a finite distance into the warped throat so that the brane is constrained to move only inside the throat, it was found that such a configuration leads to a perturbation to the warp factor affecting a correction to the warped four cycle volume. Further, this correction depends on the position of the D3-brane and thus the superpotential for the nonperturbative effect gets corrected by an overall position-dependent factor. Thus the full potential on the brane is the sum of the potential (F-term) coming from the superpotential and the usual D-term potential contributed by the interaction between the D3-brane and the anti-D3-brane. Taking these corrections into account, the volume modulus stabilization has been re-analyzed in Refs. [17–19] and the viability of inflation was investigated in this modified scenario. The stabilization of the volume modulus puts severe constraint which is difficult to solve analytically without invoking approximations. Moreover, the model needs extreme fine tuning.

The above model has been reexamined in Ref. [20] where inflation, involving both the volume modulus and the radial distance between the brane and the anti-brane participate in the dynamics. Making a rotation in the trajectory space, a linear combination of the two fields, which become independent of time and hence can be stabilized, is identified with the volume modulus. The orthogonal combination then becomes the inflaton field. The model again needs severe fine tuning and further it has been observed that when the spectral index of scalar perturbation reaches the scale invariant value, the amplitude tends to be larger than the COBE normalized value by about three order of magnitude, making the model seemingly unrealistic. However, a recent analysis using Monte Carlo method for searching the parameter space shows that COBE normalization as well as the requirement of nearly flat spectrum can be satisfied at the same time [21].

In this Letter, we analyse the possibility of a realistic model for brane inflation remaining within the large volume compactification scheme but incorporating a different correction to the inflaton potential. Infact, the authors of Ref. [22] have recently reported that there can be corrections to the inflaton potential that arise from ultra violet deformations corresponding to perturbation by the lowest-dimension operators in the dual conformal field theory. These contributions to the potential have sensitive dependence on the details of moduli stabilization and the gluing of the throat into the compact Calabi–Yau space. In the next section we briefly review the origin and form of these contributions to the potential. In Section 3, we analyse the inflationary dynamics based on the full potential. Section 4 is devoted to the summery of our analysis and conclusions.

2. The D3-brane potential

In this section we briefly outline the form of the scalar potential on a mobile D3-brane in the set up of brane–anti-brane inflation. The fluxes for the compactification of type IIB string theory on an orientifold of Calabi–Yau theory are chosen such that the internal space has a warped throat region. An example of this background geometry is the deformed conifold described in [10]. The form of the potential for the inflaton field ϕ can be schematically split into three parts; $V(\phi) = V_D(\phi) + V_F(\phi) + V_\Delta(\phi)$ where V_D is the com-

bination of the warped anti-D3-brane tension and the attractive Coulomb potential between the D3- and anti-D3-branes. V_F is the F-term potential, which gives rise to a mass term, related to the moduli stabilization effects. The exact form of V_D and V_F have been known for some time now [12] and we do not elaborate it further. Rather, we draw our attention to the other constituent of the potential, namely V_Δ as proposed in [22]. The authors, using AdS/CFT duality, describe the compactified throat region as an approximate conformal field theory which is cut off at some high mass scale M_{UV}. The throat is taken to be long enough so that the gauge theory is approximately conformal over a wide class of energy scales. The mobile D3-brane is considered to be well-separated from both the ultra violet and infrared regions. Further, it has been assumed that all the moduli are stabilized following Ref. [11]. In the present scenario, the authors observe that there are bulk moduli fields X with F-terms $F_X \sim \xi a_0^2$ where a_0 is the minimal warp factor in the throat and the field X can also be thought of as an open string modulus superfield with it lowest component being the inflaton itself. The value of ξ is fine-tunable depending on the bulk fluxes, choice of sources for supersymmetry breaking or Calabi–Yau geometries. These bulk moduli fields can couple to the fields in the conformal field theory and this coupling leads to a perturbation to the Kähler potential $\mathcal{K}$ of the form:

$$\mathcal{K}_\Delta = c \int d^4\theta \, M_{UV}^{-\Delta} X^\dagger X \mathcal{O}_\Delta \tag{1}$$

where c is a constant and $\mathcal{O}_\Delta$ is a gauge invariant operator of dimension Δ in the conformal gauge theory dual to the throat. M_{UV} is the scale which relates to the UV cutoff of the gauge theory corresponding to the large radial distance limit of the throat geometry. The above perturbation yields a scalar potential of the form:

$$V_\Delta = c M_{UV}^{-\Delta} |F_X|^2 \mathcal{O}_\Delta. \tag{2}$$

The strategy followed in [22] is to find the operators $\mathcal{O}_\Delta$ which are built out of scalar fields that can give rise to a potential on the Coulomb branch of the gauge theory and specially that contributes to the potential which controls the radial motion corresponding to the scalar field ϕ. In this context, an interesting analysis is carried out to identify the leading non-normalizable modes (normalizable in the full compact solution) of the supergravity fields that can perturb the throat geometry, thus perturb the D3-brane potential and arise from coupling to bulk moduli and four dimensional supergravity. The Kaluza–Klein excitations around $AdS_5 \times T^{1,1}$ geometry [23–25] reveals that these modes are linear combinations of the perturbations of the conformal factor of $T^{1,1}$ and that of the four-form gauge field with all four indices along $T^{1,1}$. From the harmonic expansions of these modes, in the basis $SU(2) \times SU(2) \times U(1)_R$ global symmetry, their radial dependence, r^Δ could be deduced and moreover their contribution to the D3-brane potential could be determined (see [22] for details). It is observed that at fixed radial location, the potential is minimized at some angular location and when the brane sits at this angular location the radial potential is negative. Further, the radial potential is minimized at radial distance, $r \to \infty$. Writing the canonically normalized inflation field as $\phi = \sqrt{T_3} r$, the potential is found to be

$$V_\Delta = -c a_0^4 T_3 \left(\frac{\phi}{\phi_{UV}}\right)^\Delta \tag{3}$$

where $c \sim \mathcal{O}(1)$ is a positive constant.

Taking into account the selection rules for the quantum numbers associated with the global symmetry group, the leading corrections to the inflaton potential comes from modes corresponding to $\Delta = 2/3$ or $\Delta = 2$. The field theory dual to the warped deformed conifold geometry is described by an $SU(N + M) \times$

$SU(N)$ gauge theory with bi-fundamental fields. One constructs single trace operators involving these bi-fundamental fields and their complex conjugates. These operators are also labeled by their $SU(2) \times SU(2) \times U(1)_R$ quantum numbers. The dimensions of these operators match with the dimensions of the perturbation modes contributing to the inflaton potential which is verified using AdS/CFT duality. It is observed that the operators having dimension $\Delta = 3/2$ correspond to chiral operators and the same for $\Delta = 2$ correspond to non-chiral operators. Moreover, the chiral operators determine the leading term in the inflaton potential unless they are forbidden by symmetries which are preserved in the string compactification. Infact, it is known that these chiral operators are not present in the Z_2 orbifold of the warped conifold. In such cases, the non-chiral operators determine the leading contribution to the inflaton potential. Thus we have two different models depending on the symmetries of the compactification. We will focus on the details of inflation dynamics for the model where there is no discrete symmetry forbidding the chiral operators. For this case, combining with the V_D and V_F contributions, the full inflation potential is

$$V(\phi) = D\left[1 + \frac{1}{3}\left(\frac{\phi}{M_{pl}}\right)^2 - C_{3/2}\left(\frac{\phi}{M_{UV}}\right)^{3/2} - \frac{3D}{16\pi^2\phi^4}\right] \qquad (4)$$

where $C_{3/2} \sim \mathcal{O}(1)$ and $D \sim 2a_0^4 T_3$ and $M_{UV} \sim \phi_{UV}$. We note that the functional form of the above potential is similar to the single field form of the potential, after the volume modulus field is fixed to the instantaneous minimum, of Ref. [18]. However, the microscopic interpretation of the $\phi^{3/2}$ term is different from the present context. This has been remarked in Ref. [22]. Moreover, the microscopic constraints on the coefficient of this term in the potential plays a crucial role in the discussion of the inflation dynamics. In fact, this constraint led to a two field inflationary scenario of Ref. [20]. In the present model of Ref. [22], these constraints no longer apply and hence it is necessary to analyze the single field dynamics in the prevailing scenario of relaxed constraints.

In the next section we shall discuss the background evolution of the field ϕ with the potential (4).

3. Inflationary dynamics

Let us investigate the possibilities of a viable inflation based upon the potential (4). In what follows it would be convenient to cast the evolution equation for ϕ and the Hubble equation

$$\ddot{\phi} + 3H\dot{\phi} + V_{,\phi} = 0, \qquad (5)$$

$$H^2 = \frac{1}{3M_P^2}\left(\frac{\dot{\phi}^2}{2} + V(\phi)\right) \qquad (6)$$

in the autonomous form

$$\frac{dx}{dN} = \frac{y}{\mathcal{H}}, \qquad (7)$$

$$\frac{dy}{dN} = -3y - \frac{1}{\mathcal{H}}\frac{d\mathcal{V}}{dx}, \qquad (8)$$

$$\mathcal{H}^2 = \frac{\alpha^2}{3}\left(\frac{1}{2}y^2 + \mathcal{V}(x)\right) \qquad (9)$$

where $x = \phi/M_{UV}$, $y = \dot{x}/M_{UV}^2$, $\mathcal{H} = H/M_{UV}$, $\mathcal{V} = V/M_{UV}^4$ and N designates the number of e-foldings. The field potential is expressed through the dimensionless variable x as,

$$\mathcal{V} = \mathcal{D}\left(1 + \frac{\alpha^2}{3}x^2 - C_{3/2}x^{3/2} - \frac{3\mathcal{D}}{16\pi^2\alpha^4 x^4}\right) \qquad (10)$$

where $\mathcal{D} = D/M_{UV}^4$, $\alpha = M_{UV}/M_P$. We should bear in mind that $0 < x < 1$ as $\phi < \phi_{UV} \sim M_{UV}$ and the mobile D3-brane is moving
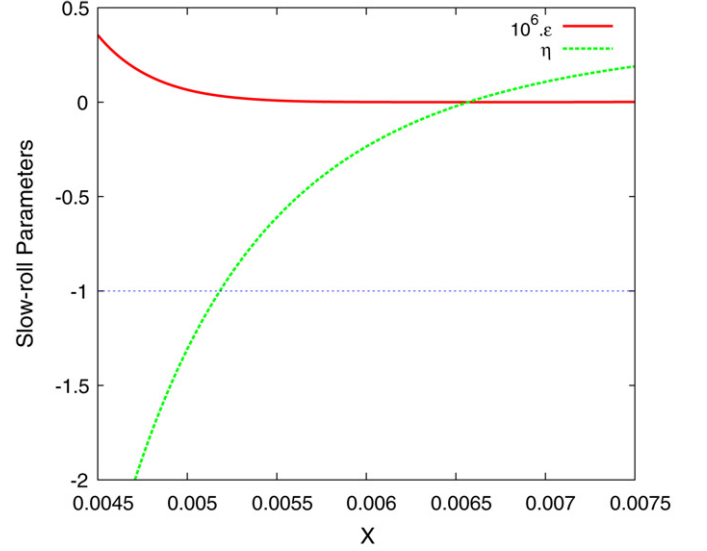

Fig. 1. Plot of the slow roll parameters ϵ and η for $\alpha^{-1} = 2.000$, $C_{3/2} = 0.00982$, $\mathcal{D} = 1.000 \times 10^{-15}$. The slow roll parameters are much smaller than one in the region of interest.

towards the anti-D3-brane located at the tip of the throat corresponding to $x = 0$.

The slow roll parameters for the generic field range are

$$\epsilon = \frac{1}{2\alpha^2}\left(\frac{\mathcal{V}_{,x}}{\mathcal{V}(x)}\right)^2$$

$$\simeq \frac{1}{2\alpha^2}\left[\frac{2\alpha^2}{3}x - \frac{3C_{3/2}}{2}x^{1/2} + \frac{3\mathcal{D}}{4\pi^2\alpha^4 x^5}\right]^2, \qquad (11)$$

$$\eta = \frac{1}{\alpha^2}\frac{\mathcal{V}_{,xx}}{\mathcal{V}(x)} \simeq \frac{2}{3} - \frac{3C_{3/2}}{4\alpha^2}\frac{1}{x^{1/2}} - \frac{15\mathcal{D}}{4\pi^2\alpha^6 x^6}. \qquad (12)$$

Since $|\epsilon| < |\eta|$ in the present case, it is sufficient to consider η for discussing the slow roll conditions. It follows from Eq. (12) that η is always less than one; it decreases as x moves towards the origin. At a particular value of x, the slow roll parameter $\eta = -1$ marking the end of inflation and takes large negative values thereafter for $x \to 0$, see Fig. 1. For generic values of the model parameters, the magnitude of slow roll parameter ϵ is much smaller than one (see Fig. 1). Hence we have to worry here about η alone.

In the case under investigation, the field x rolls from $x = 1$ towards the origin where $\bar{\text{D}}$3-brane is located. Thus the field potential should be monotonously increasing function of x. Depending upon the numerical values of the model parameters, the field potential (10) may be monotonously increasing (decreasing) or even acquiring a minimum for $0 < x < 1$. The parameters should therefore be chosen such that the potential has the right behavior,

$$\mathcal{V}_{,x} \propto \left(\frac{2}{3}\alpha^2 x - \frac{3}{2}C_{3/2}x^{3/2} + \frac{3\mathcal{D}}{16\pi^2\alpha^4 x^5}\right) > 0. \qquad (13)$$

Before we get to numerics, let us emphasize some general features of the model. The last term in Eq. (13) is always positive; hence monotonicity of $\mathcal{V}(x)$ is ensured provided, $C_{3/2} \lesssim 4\alpha^2 x^{1/2}/9$ which imposes a constraint on the coefficient, $C_{3/2}$. To avoid minimum, we require smaller and smaller values of $C_{3/2}$ as we move towards the origin before the last term in Eq. (13) could take over. It turns out that $\mathcal{D}^{1/4} \sim 10^{-4}$ for observational constraints to be satisfied pushing $C_{3/2}$ towards numerical values much smaller than one. It is then possible to make the potential flat near the origin, see Fig. 2. Since the field range viable for inflation is narrow, the potential should be made sufficiently flat to derive required number of e-foldings. In particular, it means that the slow roll parameter ϵ
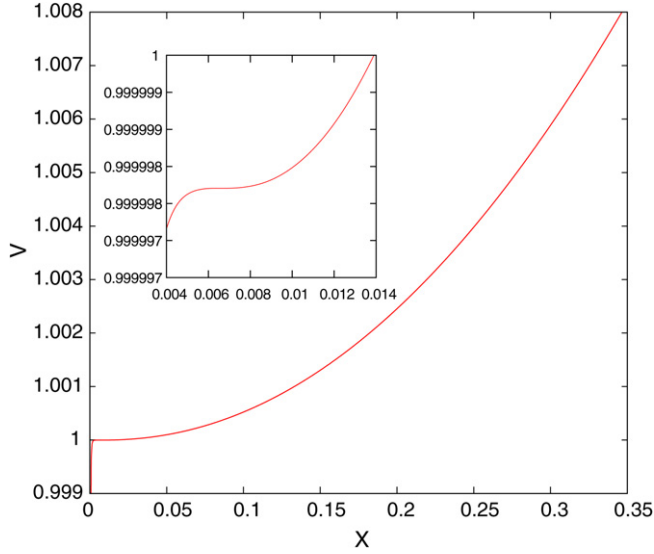
Fig. 2. Plot of the effective potential for the same values of parameters as in Fig. 1. The insert shows the flat part of effective potential near the origin which derives inflation.
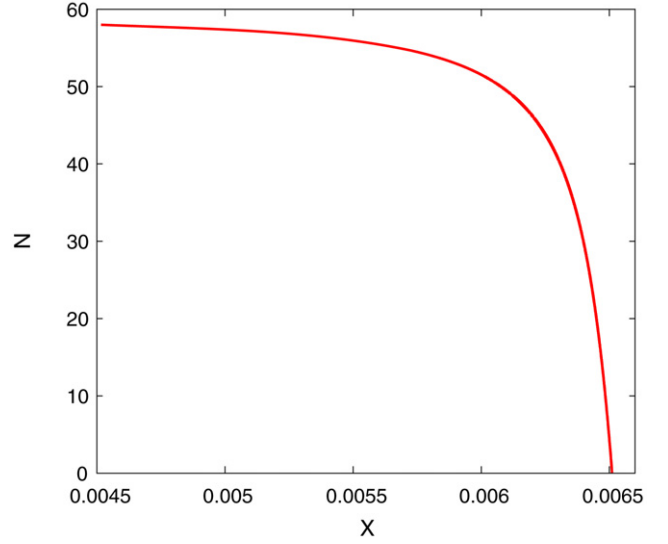

Fig. 3. The number of e-folds N versus the field x from the beginning to the end of inflation. The model parameters are same as used in Fig. 1.
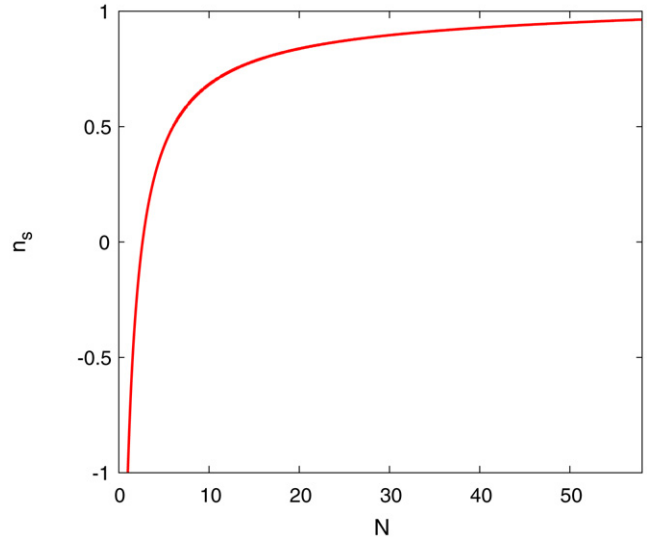

Fig. 4. Evolution of spectral index n_s versus the number of e-folds starting from the end of inflation. n_S reaches the observed value after 60 e-folds.

is very small leading to low value of tensor to scalar ratio of perturbations. This feature, however, becomes problematic for scalar perturbations. Indeed, since $\delta_H^2 \propto \mathcal{V}/\epsilon$ and $\mathcal{V} \sim \mathcal{D}$, smaller values of ϵ lead to larger values of density perturbations. Last but not the least, we emphasize a peculiarity of the potential associated with the last term in (4). The constant $\mathcal{D}$ does not only appear as an over all scale in the expression of the effective potential, it also effects the behavior of $V(x)$ in a crucial manner, for instance it changes the slow roll parameters, which also makes it tedious to set the COBE normalization. If it were not so, it would have been easy to satisfy the COBE constraint. This makes the search difficult for viable parameters.

In what follows we discuss our numerical results. We have evolved the equations of motion numerically by varying the model parameters $\mathcal{D}$, α and $C_{3/2}$. Since $\alpha \equiv M_P/M_{UV} < 1$ and $C_{3/2} \lesssim 4\alpha^2 x^{1/2}/9$ gives $C_{3/2} \lesssim 1/9$ for $\alpha = 1/2$. Avoidance of local minimum of the effective potential as the D3-brane moves towards the origin requires smaller and smaller values of $C_{3/2}$ before the Coulomb part of potential takes over. Since $\mathcal{D}$ does not quite define the scale of the potential, it becomes necessary to vary all the parameters to meet the observational constraints. It is possible to make the potential flat near the origin and since the field range viable to inflation is small for generic cases, the parameters should be suitably adjusted allowing sufficient number of e-folds. In this case, it is easy to obtain the required number of e-folds, flat power spectrum and low value of tensor to scalar ratio of perturbations. However, it is little tedious to get the COBE normalization,

$$\delta_H^2 \simeq \frac{1}{150\pi^2 M_P^4} \frac{V}{\epsilon} = \frac{\alpha^4}{150\pi^2} \frac{\mathcal{V}}{\epsilon} \qquad (14)$$

right as ϵ is small in this case. In case, $\mathcal{D} = 1.00 \times 10^{-15}$, $\alpha^{-1} = 2.000$, $C_{3/2} = 0.00982$, we have displayed the plot of effective potential in Fig. 2 which is flat near the origin. The slow roll parameters are small, specially, $\epsilon \ll 1$ in this case, see Fig. 1. With these parameters, it is possible to get the required number of e-folds, see Fig. 3. The evolution of the spectral index from the end of inflation is shown in Fig. 4; n_S reaches the observed value $n_S \simeq 0.96$ at 60 e-folds. Unfortunately, the amplitude of density perturbations, $\delta_H^2 \simeq 10^{-8}$, is large in this case. It is not possible to set the COBE normalization by merely changing α and $C_{3/2}$; it is necessary to vary $\mathcal{D}$ which in turn requires variation of other two parameters for obtaining the desired values of N and n_S. We looked for other

choices of parameters to satisfy the COBE normalization taking smaller values of $\mathcal{D}$ which further narrows the field range of interest leading to still smaller values of ϵ. We could find parameters,

$$\mathcal{D} = 1.211 \times 10^{-17}, \qquad \alpha^{-1} = 2.11991, \qquad C_{3/2} = 0.0062284$$

which can give rise to sufficient inflation, observed values n_S and δ_H^2 (see Fig. 5). In case, we start evolving from $x = x_{int} = 0.003305$, we can easily generate 60 e-foldings, and observed value of the spectral index $n_S \simeq 0.96$ (see Figs. 6 and 7). As shown in Fig. 8, we also achieve the observed value of density perturbation, $\delta_H^2 \simeq 2.4 \times 10^{-9}$, at cosmologically relevant scales observed by COBE. We note we used the exact expressions for the slow roll parameters in numerical simulations.

It is possible to change around the quoted values of parameters and still satisfy the observational constraints. In Fig. 9, we have plotted δ_H^2 versus the field x where inflaton rolls towards the origin starting from $x = x_{int} = 0.003305$. Inflation ends after 60 e-folds around $x = x_{end} = 0.0024$. Fig. 9 shows that $\delta_H^2 \simeq 2.4 \times 10^{-9}$ at the commencement of inflation. We note that δ_H^2 peaks in the neighborhood of $x = x_{int} = 0.003305$ and little deviation from the
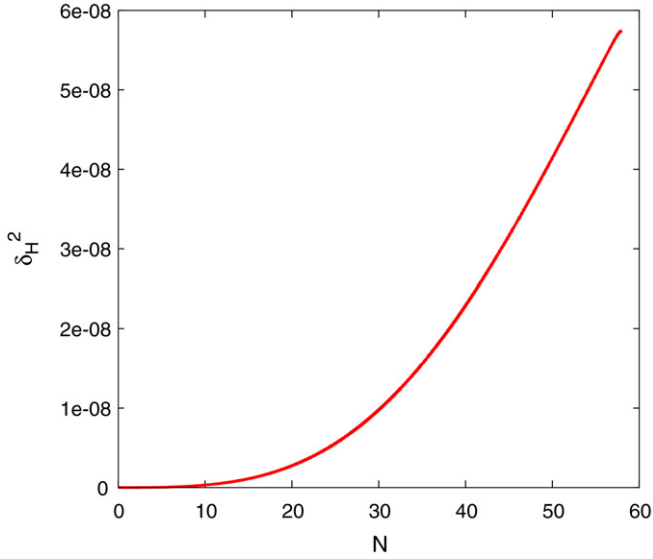
Fig. 5. Evolution of δ_H^2 versus N using the same parameters as in Fig. 1. δ_H^2 is larger than the observed values in this case.
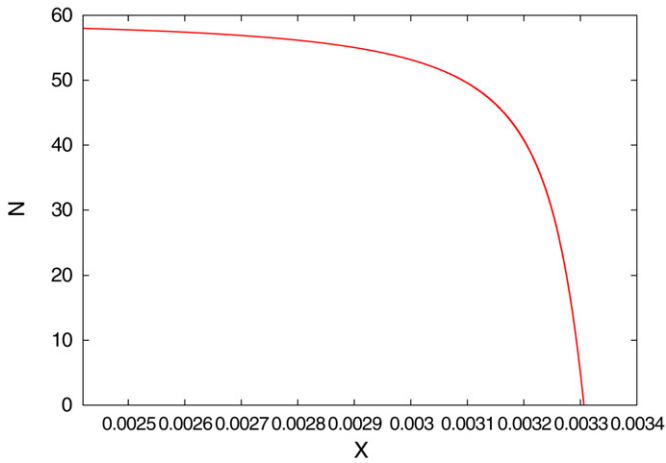

Fig. 6. The number of e-folds N versus the field x (in case of $\mathcal{D} = 1.210 \times 10^{-17}$, $\alpha^{-1} = 2.11991$, $C_{3/2} = 0.0062284$) as inflaton moves from $x_{\text{int}} = 0.003305$ towards the origin. Inflation ends around $x = x_{\text{end}} = 0.00240$ after 60 e-folds.
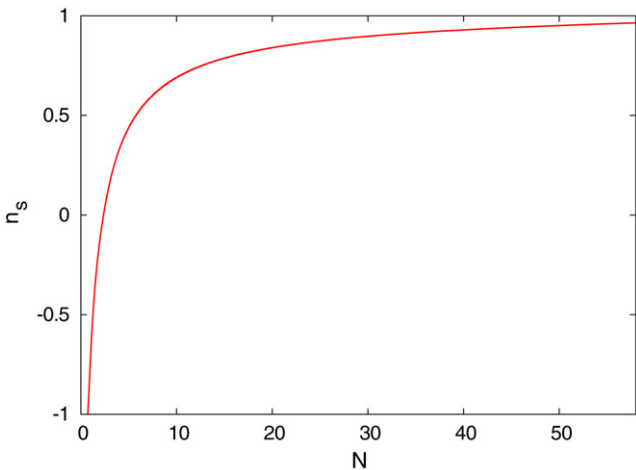

Fig. 7. Evolution of spectral index n_S versus the number of e-folds, for the same parameters values as in Fig. 6, starting from the end of inflation. The spectral index $n_S \simeq 0.96$ for $N = 60$.
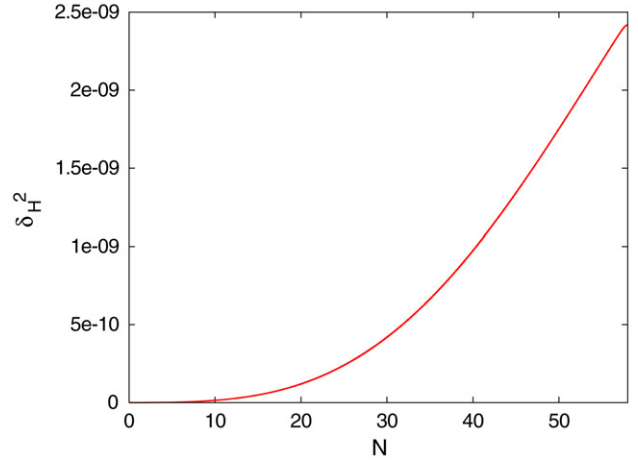

Fig. 8. Evolution of the amplitude of density perturbation δ_H^2. versus N beginning from the end of inflation. $\delta_H^2 \simeq 2.4 \times 10^{-9}$ for $N = 60$. The model parameters are same as in Fig. 6.
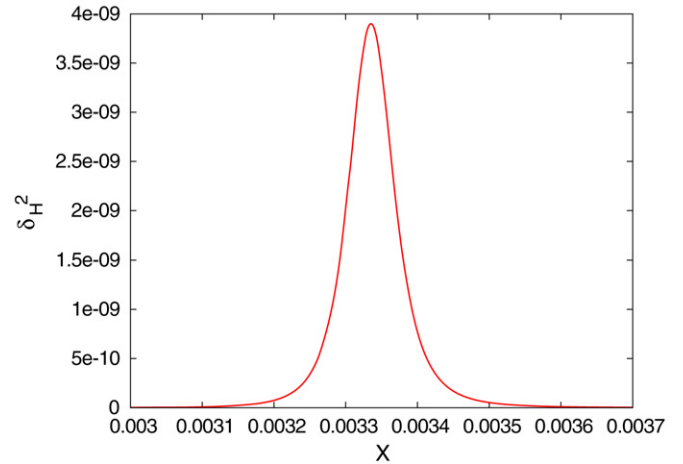

Fig. 9. Plot of δ_H^2 versus the field x for parameter values same as in Fig. 6. Inflation commences when $x = x_{\text{int}} = 0.003305$.

initial condition can easily put the density perturbations out side its observed limit. The parameters should be fine tuned to satisfy the observational constraints, for instance, $C_{3/2}$ required to be fine tuned to the level of one part in 10^{-7}. The changes at the seventh decimal puts physical quantities out side their observational bounds or the potential can acquire local minimum and spoil all the nice features of the model. Similarly other parameters also need to be fine tuned. For instance, if we take $\mathcal{D} = 1.210 \times 10^{-17}$ instead of $\mathcal{D} = 1.211 \times 10^{-17}$, the field gets into the fast roll region before it could derive 60 e-folds and COBE normalization is spoiled. The parameter α^{-1} also requires fine tuning of the order of one part in 10^{-5}. It is quite possible that the systematic search of parameters based on Monte Carlo method discussed in Ref. [21] might help to alleviate the fine tuning problem. It is, nevertheless, remarkable that the correction to throat geometry sourced by coupling to bulk allows not only to solve the well know η problem of D-brane cosmology but also helps in satisfying all the observational constraints given by the WMAP5.

4. Conclusions

In this Letter we have analysed the possibilities of inflation in a warped background with an effective potential (4). The model includes three parameters, $\mathcal{D}$, $\alpha = M_{\text{UV}}/M_P$ and $C_{3/2}$. The fact that D3-brane moves towards the tip of the throat and can reach close to it and not vise-versa imposes constraints on the model param-

eters, namely, $C_{3/2} \lesssim 0.1$ for a viable range of α. The COBE normalization demands that typically, $\mathcal{D}^{1/4} \sim 10^{-4}$ which makes $C_{3/2}$ much smaller than one. Inflation becomes possible near the origin where potential can be made sufficiently flat by appropriately choosing the parameters. Numerical values of model parameters can easily be set to obtain enough inflation and observationally consistent value of n_S. For instance, in case of, $\alpha^{-1} = 2.000$, $C_{3/2} = 0.00982$, $\mathcal{D} = 1.000 \times 10^{-15}$, we find that $N \simeq 60$ and $n_S \simeq 0.96$. The tensor to scalar ratio is very low in this case. The problem is caused by COBE normalization; the amplitude of density perturbations is large in this case, $\delta_H^2 \simeq 6 \times 10^{-8}$. This is related to the fact that the constant $\mathcal{D}$ does not give the over all scale of the potential but crucially effects the slow roll parameters and the spectral index. Infact, for several other choices of $\mathcal{D}$, we can obtain the required values of N and n_S keeping the low value of tensor to scalar ratio of perturbations. However, it is tricky to satisfy the COBE normalization. Our search led to the viable numerical values of the parameters and we have shown that for $\mathcal{D} = 1.211 \times 10^{-17}$, $\alpha^{-1} = 2.11991$, $C_{3/2} = 0.0062284$; the total number of e-foldings after horizon crossing N is around 60 and $n_S \simeq 0.96$ at the cosmologically relevant scales. For this choice of parameters, the COBE normalization can be set correctly and we find that $\delta_H^2 \simeq 2.4 \times 10^{-9}$. It is really interesting that in this case the tensor to scalar ratio is very low, $r = 16\epsilon \simeq 10^{-11}$, at the horizon crossing. It is possible to vary the parameters around the given values and still satisfy the observations constraints provided that we fine tune the parameters $\mathcal{D}$, α and $C_{3/2}$. Our search in the parameter space was carried out manually which allowed us to demonstrate that the scenario under consideration can be made consistent with the findings of WMAP5.

Finally, we should comment on the phenomenological aspects of model discussed in Ref. [18] and compare it with the analysis presented here. It was demonstrated in Ref. [20] that the single field models based upon the scenario of Ref. [18] can give rise to only few e-folds given the constrains on the model parameters. However, the two field model (i.e. the volume modulus also participates in the inflation dynamics) performs much better in this case. But it is difficult to satisfy all the observational constraints simultaneously [20], specially constraints of the spectral index and the COBE normalization are problematic. In the scenario of Ref. [22] (where the microscopic theory is different from Ref. [18]) analyzed here, the constraints on model parameters are relaxed and hence the phenomenology of single field inflation dynamics becomes possible. For such a single field model, we have demonstrated that it is possible to satisfy all the observational constraints.

Acknowledgements

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