“Phantom” Inflation in Warped Compactification

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In this paper, in a class of warped compactifications with the brane/flux annihilation, we find that the inflation may be driven by a flat direction identified as that along the number $p$ of $\mathbb{D}3$-branes located at the tip of the Klebanov-Strassler throat. The spectrum of adiabatic perturbation generated during inflation is nearly scale invariant, which may be obtained by using the results shown in the phantom inflation, since in a four-dimension effective description the evolution of energy density along the $p$ direction is slowly increasing.

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The results of recent observations are consistent with an adiabatic and nearly scale invariant spectrum of primordial perturbations, as predicted by the simplest models of inflation. The inflation stage is supposed to have taken place at the earlier moments of the universe $[1,2,3]$, which superluminally stretched a tiny patch to become our observable universe today. During the inflation the quantum fluctuations in the horizon will be able to leave the horizon and become the primordial perturbations responsible for the structure formation of observable universe $[4,5]$. Recently, how embedding the inflation scenario into string theory has received increasing attentions, e.g. see Ref. $[6,7,8,9,10]$ for reviews.

The inflation can be generally regarded as an accelerated or superaccelerated stage, and so may defined as an epoch when the comoving Hubble length decreases. This also may occurs namely during a null energy condition violating expansion which corresponds to e.g. the phantom inflation $[11,12,13]$, see also Ref. $[14]$ for comments. The phantom inflation is only a phenomenological description for such a superaccelerated stage, it dose not equal to the existence of a phantom. Here we will search for a realistic implement of the phantom inflation in possible high energy physical theory, e.g. the string theory. Recently, it has been shown in Ref. $[15]$ that in a class of string compactification with the warped metric, when there are $p \ll M \mathbb{D}3$-branes inside the KS throat, which is a deformed conifold with $M$ units of RR 3-form flux around the internal $S^3$, where $p$ is the number of $\mathbb{D}3$-branes, a metastable and nonsupersymmetric bound state with positive vacuum energy density will form. This false vacuum may be responsible for the inflation, which has awaken a string version of old inflation $[16]$. When $p$ is beyond some critical value $p_c$, the false vacuum will disappear and the system will roll down to the supersymmetric stable vacuum, which signs the end of inflation. In the language of NS5-branes, what happens is that the $\mathbb{D}3$-brane “blows up” into a NS5-brane wrapping a $S^2$ in the internal $S^3$. The resulting state for $p \ll M$ is a metastable state which breaks supersymmetry. When $p$ is close to $M$, the radius of $S^2$ will approach and surpass the radius of the equator of $S^3$, which leads that the NS5-brane becomes perturbatively unstable, and will directly rolls down to a supersymmetric vacuum. In this note, we find that under some conditions there may be a flat direction along the number $p$ of $\mathbb{D}3$-branes in this scenario, which may lead to a nearly scale invariant spectrum of adiabatic perturbation. The spectrum can be obtained by using that similar to the phantom inflation $[11,12]$, since in the four dimension effective description the climbing up of energy density along $p$ direction corresponds to an evolution with the null energy condition violation.

We begin with a 10 dimension CY manifold with the warped KS throat $[17]$. The metric of warped throat may be taken as

$$ds^2 = \frac{1}{\sqrt{f(r)}} ds^2_{(4)} + \sqrt{f(r)}(dr^2 + r^2 ds^2_{(5)})$$

for $r < r_s$, where $ds^2_{(5)}$ is the angular part of the internal metric and $r$ is the proper distance to the tip of the throat. When $r > r_s$, this metric can be glued to the metric of the bulk of the compact space, usually taken to be a CY manifold. The warp factor $f(r)$ has a minimal value at $r_0$, which is determined by $\beta \equiv \frac{\psi_0}{\psi_0} \sim e^{-\frac{\pi}{\psi_0}}$, and when $r_0 \sim r < r_s$, $f(r)$ may be approximately regarded as $f(r) = \left(\frac{R^4}{R^4_s}\right)^{\frac{1}{3}}$, where $R^4 = \frac{2\pi^2}{g_s N}\alpha'^2$, $N$ is equal to the product of the fluxes $M$ and $K$ for the RR and NSNS 3-forms, respectively, and $g_s$ is the string coupling and $\alpha'$ is set by the string scale.

It has been shown in Ref. $[12]$ that when putting $p \ll M \mathbb{D}3$-branes at the tip of the KS throat, the system will be relaxed to a nonsupersymmetric NS5-brane “giant graviton” configuration, in which the NS5-brane warps a $S^2$ in $S^3$, and carries $p$ unites flux for 2-form, which induces the $\mathbb{D}3$-charge. The $S^2$ is inclined to expand as a spherical shell in $S^3$. This process may be parameterized by an angle $0 \lesssim \psi \lesssim \pi$, where $\psi = 0$ corresponds to the north pole of $S^3$ and $\psi = \pi$ is the south pole. The angular position indeed appears as a scalar in the world volume action, which describes the motion of the NS5-brane across the $S^3$. The total effective potential can be given by

$$V_{\text{eff}}(\psi) = M\beta^4 T_3 \left(\frac{b_0^2 \sin^2 \psi}{\pi^2} + \tilde{V}^2(\psi) + \tilde{V}(\psi)\right)$$

with $b_0 \simeq 0.9$, where $\tilde{V}(\psi) = \frac{\psi - \sin (2\psi)/2}{\pi}$ and $T_3$ is the $\mathbb{D}3$-brane tension. This potential is plotted in Fig.1.
with respect to $\psi$ and $p/M$. The increase of potential energy along $p$ direction is actually discrete, since $p$ is the positive integer. However, since $M$ is quite large, thus $p/M$ may be approximately regarded as continual one.

In the regime with $p/M < 0.08$, the system sits on a metastable state corresponding to a static NS5-brane wrapping a $S^2$ in $S^3$, which is classically stable false vacuum. It only may decay to a supersymmetric state by quantum tunneling, however, as has been shown that this probability is exponentially suppressed $^{13}$. In Fig.1, it can be clearly seen that this metastable bound state corresponds to $\psi = 0$, thus the energy density of this metastable vacuum is $V_{\text{eff}}(\psi = 0) = 2p\beta^3 T_3$ given by Eq. (2). It is this vacuum energy that drives inflation. Thus the Hubble expansion is given by

$$h^2 = \frac{2p\beta^3 T_3}{3}, \quad (3)$$

where $8\pi/m_5^2 = 1$ is set. While the true minimum is at $\psi = \pi$, in which the potential energy is 0.

In the regime $p/M \gtrsim 0.08$, the false vacuum disappears, which means that the nonsupersymmetric configuration of $p$ D$3$-branes becomes classically unstable and will relaxes to the supersymmetric minimum by a classical rolling, which denotes that the D$3$-branes cluster to form the maximal NS5-branes, and then rolls down towards the bottom of the potential$^1$, at the north pole $\psi = \pi$, in which the potential energy is 0. Thus in this case, the inflation will be expected to end. The result of this process is $M - p$ D3-branes instead of the original $p$ D$3$-branes appearing in the tip of KS throat, while the 3-form flux $K$ is changed to $K - 1$, which is referred to as brane/flux annihilation in Ref. $^{13}$.

The point here is that the inflation occur only when $p/M < 0.08$, and during this period some D$3$-branes may be expected to continually enter into the throat which will continually increase the value of $p$, and when $p/M \gtrsim 0.08$, the metastable minimum of the potential along $\psi$ will disappear and the potential becomes a monotonic decreasing function of $\psi$, and the exit of inflation is induced by a subsequent rolling along $\psi$, see the dashed line in Fig.1. Thus this inflation model is actually quite similar to the hybrid inflation in Ref. $^{19}$, which is referred to as a water fall field in the hybrid inflation$^2$. However, the distinguish is that here the energy density driving the inflation is increasing before $p$ approaches the critical value.

FIG. 1: The figure of the potential $^{2}$ up to an overall scale $M^3 T_3$. In the left lower panel, the solid line is the effective potential with $p/M = 0.03$. When D$3$-branes are pulled into the throat continuously, the metastable minimum will rise inch by inch. The dashed line is the effective potential with $p/M = 0.1$. The right up panel is the figure of the effective potential with respect to $\psi$ and $p/M$. The dashed line denotes the evolution of energy density along $p$ direction during inflation and the rolling down along $\psi$ direction after inflation. The total process is quite similar to the hybrid inflation $^{19}$. However, the distinguish is that here the energy density driving the inflation is increasing before $p$ approaches the critical value.

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1 It should be mentioned that this classic rolling may also support a slow roll inflation under some conditions, which has been studied in Ref. $^{18}$.

2 In Ref. $^{18}$, $\psi$ was regarded as inflaton, which is different from the point here. The curvature perturbation in their model is given by the curvaton mechanism $^{23, 24}$, since the curvature of the potential of $\psi$ around the metastable minimum is much larger than the Hubble rate during inflation, which implies that the curvature perturbation associated to the quantum fluctuations of $\psi$ is heavily blue tilt and thus suppressed on large scale.

3 It is interesting to note that the super inflation phase with the null energy condition violation also appears in loop quantum gravity, which have been studied in Refs. $^{20, 21}$ and also $^{22}$, or other gravity theories with high order corrections.
Those $\overline{D}3$-branes walking into the throat should be originally in the compactification bulk at $r > r_s$. Their initial energy at $r = r_s$ is obtained by multiplying the brane tension with the redshift factor $r_s/R$. It is required in an effective theory that the $\overline{D}3$-brane should be a small perturbation of the system and the vacuum energy should be determined by the $\overline{D}3$-branes in the throat.

To this purpose, it must be required that the initial energy $\Delta p T_3^3(r_s/R)^4$ of $\overline{D}3$-brane entering into the throat in units of Hubble time should far less than $2p\beta^4T_3$. Thus $\Delta p(r_s/R)^4 \ll p\beta^4$ or equally $\epsilon(r_s/R)^4 \ll \beta^4$ need to be satisfied, which implies that a mild warping of the throat or a large value of $p$ should be considered, which may be actually obtained easily. For instance, we may take $K \sim M \gg \Delta p/\beta^4$. When entering the throat, the $\overline{D}3$-branes can feel a net radial force proportional to the 5-form flux, which will attract them to the tip at $r = r_0$. This force is a sum of gravitation and 5-form contributions. Note that For the $D3$-branes, both terms cancel.

When the brane arrives at $r = r_0$, the interaction between $\overline{D}3$-branes will be involved. However, here for a simplified discussion, we only study its net effect, i.e. the energy density driving the inflation is added with $p$ units.

In addition, we also assume here that all the moduli have been fixed, as has been done in Ref. [25]. Thus to model the evolution required, we introduce a phantom energy, whose simplest implementation is $\bar{\eta}$-dependence perturbation $\delta \eta$.

The primordial perturbation can be generated during this evolution induced by the increase of number of $\overline{D}3$-branes along $p$ direction. To calculate it simply but effectively, we adopt the work hypothesis in which the null energy condition violating evolution behaves like the phantom energy, whose simplest implementation is a scalar field with the reverse sign in its kinetic term. Thus to model the evolution required, we introduce a phantom field $\varphi$ here. Instead of the rolling down of normal scalar field along the potential, the phantom field will be driven up along its potential, as has been analytically and numerically shown in e.g. Refs. [26, 27], and at later time for some potentials, analogous to the slow roll regime of normal scalar field, the phantom field will enter into a “slow climb” regime, where $\varphi^2 \ll \mathcal{V}(\varphi)$. Thus we can have $h^2 \simeq \mathcal{V}/3$, where $\mathcal{V}$ is the potential of phantom field, which can be distinguished from $V(\psi)$. In this case, if we approximately regard $|\epsilon|$ as constant, which here means that $\Delta p/p$ is nearly unchanged, we can have $\mathcal{V}(\varphi) = 2p\beta^4 T_3\exp\left(-\sqrt{\frac{\Delta p}{p}}\varphi\right)$. (4)

where the prefactor is determined by Eq.(3), and $p_i$ is the value of $p$ at the beginning time of phantom inflation. In Eq.(4), $-\varphi$ direction corresponds to the increasing direction of $p$ in Fig.1. In the slow climb approximation $|\epsilon| \simeq \frac{(V')^2}{2} \frac{12}{|\chi|}$, where the prime is the derivative with respect to $\varphi$, thus combining it with Eq.(4), we have $|\epsilon| \simeq \frac{2\beta^4}{|\chi|}$, which is consistent with the requirement of our model. Thus in this sense the upclimbing of energy density induced by the increase of $p$ may be depicted well by the evolution of a phantom field.

The calculations in this case have been done in Ref. [11, 12]. It can be noted that here Fig.1 is nearly same as Fig.1 in Ref. [12] with a replacement $\psi$ with $\sigma$, in which to have an exit from the phantom inflation, a normal scalar field $\sigma$ is introduced to implement a water fall reheating, like that occurring in the hybrid inflation. Here we will reillustrate the calculations in Ref. [12], however, renew in a way independent of the phantom potential, which will be helpful for the analysis of spectral index tilt here. We will work in the longitudinal gauge. In the momentum space, the equation of motion of gauge invariant variable $u_k$, which is related to the Bardeen potential $\Phi$ by $u_k \equiv a\Phi_k/\varphi'$, is given by

$$u_k'' + \left(k^2 - \frac{\beta(q)}{\eta^2}\right)u_k = 0$$

(5)

where the prime is the derivative with respect to the conformal time $\eta$ and $\beta(q) \simeq \epsilon - \left(\frac{d\ln|\epsilon|}{dN}\right)^2/2$, which may be obtained by taking $|\epsilon| \simeq 0$ in Eq.(21) of Ref. [11], where the higher order terms like $\frac{d\ln|\epsilon|}{dN}$ and $\frac{d^2\ln|\epsilon|}{dN^2}$ have been neglected since for inflation $\frac{d\ln|\epsilon|}{dN} \ll 1$, and $N$ is the e-folding number of mode with some scale $\sim 1/k$ which leaves the horizon before the end of phantom inflation and may be defined as

$$N = \ln\left(\frac{a_h h_e}{a_h}\right),$$

(6)

where the subscript ‘e’ denotes the quantity evaluated at the end of the phantom inflation. Here since $h$ is nearly constant, we can have $N \simeq \ln\left(\frac{a_h}{a}\right)$. Note that for $\frac{d\ln|\epsilon|}{dN} \ll 1$, $\beta$ is near constant for all interesting modes $k$. Thus Eq.(6) is a Bessel equation and its general solutions are the Bessels functions with the order $v = \sqrt{\beta + \frac{1}{2}} \simeq \frac{1}{2} + \beta$, since $|\beta| \ll 1$.

In the regime $k\eta \rightarrow \infty$, the mode $u_k$ are very deep in the horizon. Thus Eq.(6) can be reduced to the equation of a simple harmonic oscillator, in which $u_k \sim e^{-ik\eta}/(2k)^{3/2}$ can be taken as the initial condition. In the regime $k\eta \rightarrow 0$, the mode $u_k$ are far out the horizon, and become unstable and grows. In long wave limit, the expansion of the Bessel functions to the leading term of $k$ gives

$$k^{3/2}u_k \simeq \frac{\sqrt{\pi}}{2^{3/2}/2\sin\left(\pi v\right)}\Gamma(1 - v)\left(\frac{-k\eta}{2}\right)^{\frac{3}{2} - v}.$$
Thus we can have $k^3 |u_k|^2 \sim k^{1-2\nu} \sim k^{-2/3}$, which gives the spectral index of $\Phi$

$$n_\Phi - 1 \simeq -2\epsilon + \frac{d \ln |\epsilon|}{dN}, \quad (8)$$

This spectrum is actually given by the constant mode of Bardeen potential $\Phi$, which may easily be inherited by the constant mode after the reheating, and thus the curvature perturbation $\zeta$ in comoving supersurface will have same spectrum as $\Phi$, i.e. $n_\zeta = n_\Phi$. When $|\epsilon| \ll 1$, we see that the spectrum is nearly scale invariant.

The tilt of spectral index is interesting for the observations. When $\epsilon$ is constant, the spectrum is slightly blue, since $\frac{d \ln |\epsilon|}{dN} = 0$ and $\epsilon \lesssim 0$, which looks like not favored by the observation [29]. However, whether $\epsilon$ is constant is actually dependent of the rate that the D3-branes are sent into the throat. When this rate is proximately constant, we will have $\epsilon$ constant, or $\epsilon$ will be time dependent. Thus in principle the spectrum may be blue or red, dependent of physical details of D3-branes motion in the compactification bulk. To make the spectrum become slightly red, $\frac{d \ln |\epsilon|}{dN}$ is negative and larger than $2|\epsilon|$ is required, which means that $|\epsilon|$ must be increased with the decreasing of $N$. Thus it require that in per Hubble time the number $\Delta \Phi$ of D3-branes entering the throat should be more and more with the time. We show this point by a detailed example, which is plotted in Fig.2. This result indicates that under certain condition the spectrum may fit the observations well [29]. Note that the spectral index is to large extent determined by the random motion of D3-branes, thus to match the observation, it seems that some fine tuning for the motion of D3-branes must be required, however, in turn, if the primordial perturbation is actually generated by such an inflation model, in some sense the spectral index observed will encode the ordered motion of D3-branes.

The amplitude of curvature perturbation $\zeta$ may be given by

$$P_\zeta = \frac{k^3}{2\pi^2} |\zeta_k| \simeq \frac{1}{2\pi^2 |\dot{\phi}|} \cdot \frac{1}{2\pi} \frac{|\dot{\phi}'|}{a}^2$$

$$\simeq \frac{1}{2|\epsilon|} \left( \frac{h}{2\pi} \right)^2, \quad (9)$$

where $h = \dot{\phi}^2/2$ and Eq.(7) have been used, in addition we also use $\zeta \simeq \Phi/\epsilon$ since $|\epsilon| \ll 1$. This may be seen by using the relation between $\zeta$ and $\Phi$, which is $\zeta \simeq \Phi + \frac{1}{2} \left( \frac{\dot{\phi}^2}{2\epsilon} + \Phi \right)$. Thus when $|\epsilon| \ll 1$, it naturally leads to $\zeta \simeq \Phi/\epsilon$. When $|\epsilon| \gg 1$, we have $\zeta \simeq \Phi$, which has been used in the calculations of primordial perturbations generated during a slow expansion with $\epsilon \ll -1$ [11]. Eq.(9) is the same as that of usual inflation models.

The spectrum of tensor perturbation can be calculated in a similar way [12, 30]. The amplitude of its spectrum is given by

$$P_T \simeq \frac{k^3}{2\pi^2} \cdot 4 \left( \frac{\sqrt{2}v_k}{a} \right)^2 \simeq 8 \cdot \left( \frac{h}{2\pi} \right)^2, \quad (10)$$

where the gauge invariant variable $v_k$ is related to the tensor perturbation $h_k$ by $v_k = ab_k/\sqrt{2}$. Thus the ratio of tensor to scalar perturbation is $r = P_T/P_\zeta \simeq 16|\epsilon|$. Here in unit of Hubble time $|\epsilon| \simeq \Delta p/2p$, thus $r \simeq 8\Delta p/p$, which can be large, whose detailed value depends on the change of $\Delta p/p$. In general in this model $p \sim 10^5$ to $10^6$ and $1 \lesssim \Delta p \lesssim 10$ or several 10. $\Delta p \gtrsim 1$ is because in probability in unit of Hubble time there should be at least one D3-brane entering into the throat. This actually sets a lower limit for $r$. In our example, $p \approx 500$, thus $r \gtrsim 8/500 \approx 0.016$, which suggests that if $r$ is enough small, for example $r < 0.01$, this model will be ruled out. $\Delta p \lesssim 10$ or several 10 is to make $|\epsilon| \ll 1$ valid in all time. In our example, $r \approx 8 \times 10/500 = 0.16$, which is consistent with recent observation [31], in which $r < 0.20$.

Thus in principle our model predicts a mild range of $r$, i.e. $0.01 \lesssim r \lesssim 0.2$, dependent of the details of model, which may be tested in future observations.

The spectral index of tensor perturbation is nearly scale invariant with the tilt equal to $-2\epsilon$ [12]. Here note that $\epsilon \lesssim 0$, thus the tensor spectrum is slightly blue, which is a distinct feature, since for the usual inflation models, the tensor spectrum is generally slightly red. However, this feature can be consistent with the observations [32].

In principle, it seems that we also may introduce other fields with any kinetic term or fluids to describe this null
energy condition violating evolution, as long as their energy density is increased with the time. However, note that for such fields or fluids, generally the sound speed $c_s^2$ is not exact 1 in their perturbation equation. This means in this case the causal structure of background and perturbation evolution has been altered, which actually is not expected to occur in our model, since here there seems not any motion or evolution leading to such a change of the causal structure. Further, there might exist some fluid which violates the null energy condition and in the meantime whose parameter $w$ of state equation and $c_s^2$ are independent each other, and then $c_s^2$ can be set as 1. However, in this case we will obtain same Eq. (10), where $u_k \equiv \Phi_k/\sqrt{\rho + p}$, and thus same Eqs. (9) and (10). This means that this case is actually equal to that of the simple phantom field. Thus using the perturbation equation derived for the phantom field can be reasonable for our model.

In summary, we note that in a class of warped compactifications with the brane/flux annihilation, there may be a flat direction along the number $p$ of $\overline{D}3$-branes located at the tip of KS throat, which may drive a period of phantom inflation, since in the four dimension effective description the upclimbing of energy density along $p$ direction corresponds to a null energy condition violating evolution. The spectrum of adiabatic perturbation can be nearly scale invariant with a mild tilt dependent of the change of $\overline{D}3$-branes number entering into the throat. This work seems slightly idealistic, however, it may be a start, with which it may be expected that one could envisage more examples in string theory, in which the four dimension effective descriptions may share some similar characters as the phantom inflation. We have neglected some details involving the motion of the $\overline{D}3$-branes and the interaction between them, which maybe able to lead to some different features of primordial spectra. This will be left in the future works.

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