Inflation driven by single geometric tachyon with D-brane orbiting around NS5-branes

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We investigate models in which inflation is driven by a single geometrical tachyon. We assume that the D-brane as a probe brane in the background of NS5-branes has non-zero angular momentum which is shown to play similar role as the number of the scalar fields of the assisted inflation. We demonstrate that the angular momentum corrected effective potential allows to account for the observational constraint on COBE normalization, spectral index ns and the tensor to scalar ratio of perturbations consistent with WMAP seven years data.

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I. INTRODUCTION

The efforts to investigate the time dependent backgrounds in string theory have recently attracted attention in context with the study of tachyon condensation [1]. The effective field theoretic set up to describe the dynamics of the rolling tachyon is provided by Dirac-Born-Infeld (DBI) action [2]. It was observed that in the process of tachyon condensation, the unstable brane or the brane-antibrane pair can decay to form a new stable D-brane.

It is interesting to note that the equation of state parameter of the rolling tachyon field varies from zero to minus one which gave rise to hope that the open string tachyon on the unstable D-brane can play the role of inflaton [3,4] (see also Refs.in [5] on the related issue). This idea was also generalized to the radion field in the case of a brane moving towards an antibrane, and vice versa [6]. But since the effective potential of the rolling tachyon computed in the perturbative string theoretic framework does not contain adjustable parameters, it is not surprising that the model fails to be compatible with the requirement of slow-roll and COBE normalization. Efforts were made to address these problems via warped compactification [7,8] of the string theory. Inspite of several attempts to overcome the problems in open string tachyon cosmology, it seems unlikely that this tachyon field is responsible for inflation. Attempts to assign the role of dark matter fluid to rolling tachyon are also faced with a serious problem of caustic formation which points towards the incompleteness of the theory.

In string theory, however, there might be many interesting time dependent backgrounds and one of such possibilities was recently investigated in [9]. In this scenario the motion of D-brane, as a probe brane in the background of k coincident NS5-branes, gives rise to an interesting dynamics which can be mapped to DBI action. Indeed, it was shown in the five-dimensional brane world models (codimension-1 brane) [10] and in the (p+3)-dimensional string theory (codimension-2 brane) [11] that the presence of NS-NS type brane is indispensable to obtain flat backgrounds on the transverse dimensions. This suggests that desired brane world models must involve the NS-brane as their background branes and the SM-branes(D-branes) are then placed near the background NS-branes. These models, however, are faced with an instability problem of the D-brane. It is known [4] that a D-brane propagating at some distance from a stack of k parallel NS5-branes becomes unstable. The NS5-branes are much heavier than the D-branes in the regime of small string coupling in this picture. Geometrically this means that the NS5-branes form an infinite throat in space-time and the string coupling increases as we move towards the bottom. Being lighter, the probe brane is gravitationally pulled towards the NS5-branes.

The D-brane preserves half of the supersymmetry which is different from the other half preserved by the NS5-branes in Type-II theory. Consequently, as the probe brane comes nearer the source brane, the supersymmetry of the system is completely broken which gives rise to a tachyonic degree of freedom on the D-brane. Indeed, the radion becomes tachyonic in this case and there is a map between the tachyonic radion field living on the world volume of the probe brane and the rolling tachyon associated with a non-BPS D-brane. Thus the motion of the probe brane in the throat could be described by the condensation of the tachyon. It is also possible to study the motion of the probe brane in the background of a ring of NS5-branes instead of coincident branes Refs.[12,13]. It is observed that the radion field becomes tachyonic when the probe brane is confined to one dimensional motion inside the ring. The condensation of the geometrical tachyon can play important role in cosmology [14,15].
Cosmological applications of these models with open string tachyon have been studied in literature using BPS and non-BPS D-branes. For instance in [10], an assisted inflation using open string tachyons was studied, where N open string tachyons are allowed to roll simultaneously to obtain slow-roll inflation. In [13], on the other hand, the authors studied rolling geometrical tachyon induced on the N D3-branes moving in the vicinity of NS5-branes to show that this system coupled to gravity gives the slow-roll assisted inflation of the scalar field theory. Further assisted inflation scenario from the rolling of N BPS D3-brane into the NS5-branes, on a transverse geometry of $R^3 \times S^1$, coupled to four dimensional gravity has also been studied in [18]. In general, tachyon-inflation models suffer from the large $\eta$-problem as in the conventional models and are not favored for slow-roll inflation. This difficulty may be avoided by allowing a large number of tachyons to simultaneously roll down to assist the inflation as mentioned above. Such an assisted inflation, however, necessarily contains a large number of D3-brane in the theory and consequently it is not adequate if we want to find a theory with a single D3-brane. Thus in that case we may need to think about a different type of the theory which only involves a single D3-brane but can satisfy the requirement of the slow-roll inflation.

As mentioned above, in the configurations with D-brane(s) near NS5-branes the supersymmetry of the system is completely broken and this gives rise to a tachyonic degree of freedom on the D-brane and it becomes a non-BPS D-brane. The D-brane eventually falls into the fivebranes and decays into a pressureless fluid called \"tachyon matter\". Such a decay of the D-brane, however, can be avoided in special cases. In [10], it was demonstrated that for certain values of energy and angular momentum the D-brane orbits around the fivebranes, maintaining a fixed distance from the fivebranes all the times and the decay of the D-brane is suppressed. Indeed, in the case of nonzero angular momentum the effective tachyon potential takes very different as compared to the case of zero angular momentum and can give rise to viable inflationary scenario.

In this paper we shall examine inflationary models in which the inflation is driven by a single geometrical tachyon with an assumption that the probe D-brane has non-zero angular momentum.

II. D-BRANE DYNAMICS NEAR NS5-BRANES

We begin our discussion by briefly reviewing the work presented in [8] and [10]. In the presence of $k$ coincident NS5-branes, the metric, dilaton and NS-NS 3-form fields are given by

$$ds^2 = dx_\mu dx_\nu + H(x^m)dx^m dx^m' = G_{MN}dx^M dx^N,$$

where $x^\mu (\mu = 0, 1, ...5)$ are the coordinates along the world volume of the $k$ coincident NS5-branes, while $x^m (m = 6, 7, 8, 9)$ are the coordinates along the transverse dimensions. Also $h(x^m)$ is a harmonic function,

$$h = 1 + kl^2/r^2,$$

where $r^2 = \sum_{m=6}^{9} x^m x_m$ and $l_s$ is the string length.

Let us consider a Dp-brane moving in the vicinity of the stack of NS5-branes and stretched along the directions $\langle x_1, ...x_p \rangle$ with $p \leq 5$. If we label the world volume of the D-brane by $\xi^\mu$, $\mu = 0, 1, ...p$. Then in the static gauge we have $\xi^\mu = x^\mu$. The dynamics of the world volume fields of the Dp-brane propagating in the above background fields is governed by DBI (Dirac-Born-Infeld) action

$$S_p = -\tau_p \int d^{p+1}x e^{-(\Phi - \phi_0)} \sqrt{-\det (G_{\mu\nu} + B_{\mu\nu})}, \quad (3)$$

where $G_{\mu\nu}$ and $B_{\mu\nu}$ are the pullbacks of $G_{MN}$ and $B_{MN}$:

$$G_{\mu\nu} = \partial x^M/\partial \xi^\mu \partial x^N/\partial \xi^\nu G_{MN}, \quad B_{\mu\nu} = \partial x^M/\partial \xi^\mu \partial x^N/\partial \xi^\nu B_{MN}. \quad (4)$$

In Eq.(4) $x^M = (\xi^\mu, x^m)$, and $x^m$ represent the position of the Dp-brane in the transverse space and they give rise to world volume scalars $X^m(\xi^\mu)$. In this paper we shall assume that the world volume components of the B-field vanish, i.e., $B_{\mu\nu} = 0$ as it generally breaks the isotropy of the Dp-brane world volume. In case of spatially homogeneous and isotropic background, $X^m = X^m(t)$. $G_{\mu\nu}$ reduces to

$$G_{\mu\nu} = \eta_{\mu\nu} + \delta^0_{\mu} \delta^0_{\nu} h(X^n) X^m X^m', \quad (5)$$

and upon introducing polar coordinates, $X^0 = R \cos \theta$ and $X^7 = R \sin \theta$, the action (3) takes the following form

$$S_p = -\tau_p \int dt \sqrt{h^{-1}(R) - (\dot{R}^2 + R^2 \dot{\theta}^2)}, \quad (6)$$

where $h(R)$ is now $h(R) = 1 + kl^2/R^2$ and we have set the volume of the Dp-brane equal to one.

The angular momentum and conserved energy following from (6) take the forms

$$L = \tau_p \frac{R^2 \dot{\theta}}{\sqrt{h^{-1}(R) - (\dot{R}^2 + R^2 \dot{\theta}^2)}}, \quad (7)$$

$$E = \tau_p \frac{1}{h\sqrt{h^{-1}(R) - (\dot{R}^2 + R^2 \dot{\theta}^2)}}, \quad (8)$$

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1 for further examples see [20]
Solving these two equations in terms of $\dot{R}^2$ and $\dot{\theta}^2$, one obtains

\[ \dot{R}^2 = \frac{1}{\varepsilon^2 h^2} \left[ \varepsilon^2 h - \left(1 + \frac{l^2}{R^2}\right) \right], \]

and

\[ \dot{\theta}^2 = \frac{1}{R^4 h^2} \frac{l^2}{\varepsilon^2}, \]

where $l$ and $\varepsilon$ are defined as $l = L/\tau_p$ and $\varepsilon = E/\tau_p$, respectively. Note that the D-brane orbits around the fivebranes, maintaining certain distance from the fivebranes all the time, i.e., $\dot{R} = 0$ provided the conditions

\[ \varepsilon = 1 \quad \leftrightarrow \quad E = \tau_p, \]

and

\[ l = \sqrt{k_l s} \quad \leftrightarrow \quad L = \sqrt{k_l s} \tau_p \]

are satisfied.

We next consider an action,

\[ \tilde{S}_p = -\tau_p \int dt \sqrt{1 + \frac{l^2}{R^2}} \sqrt{h^{-1} - \dot{R}^2} \]

In [19], it was shown that the action $\tilde{S}_p$ in (13) is classically equivalent to $S_p$ in (6) as far as the radial motion is concerned. This in turn means that the tachyonic behavior described by $\tilde{S}_p$ is equivalent to that of the original action $S_p$ as the (geometrical) tachyon $T$ is only a field redefinition of $R$:

\[ dT = \sqrt{h(R)}dR \]

The solution of (14) is [8]:

\[ T(R) = \sqrt{k_l s} + \frac{1}{2} \sqrt{k_l s} \ln \frac{\sqrt{k_l s} + R^2}{\sqrt{k_l s} - \sqrt{k_l s}} \]

and in terms of $T$, $\tilde{S}_p$ can be rewritten as

\[ \tilde{S}_p = -\int dt \tilde{V}(T) \sqrt{1 - T^2} = \int dt L(t), \]

where $\tilde{V}(T)$ is given by

\[ \tilde{V}(T) = \tau_p \frac{\sqrt{1 + \frac{l^2}{R^2}}}{\sqrt{h(R)}}. \]

It should be noted that $\tilde{V}(T)$ becomes flat for $l = \sqrt{k_l s}$ and in this case the tachyon does not roll at all, thereby indicating that the tachyonic degree of freedom induced on the D-brane disappears and the D-brane returns to the stable brane.

### III. Inflation from a Single Geometrical Tachyon

In this section we shall examine the cosmological implications of the geometrical tachyon in the inflationary universe. As mentioned earlier, the assisted inflation was proposed to overcome the large $\gamma$ problem. The assisted inflation consists of introducing a large number of D3-branes or equivalently a large number of scalar fields. This idea was extended to inflationary model using N-geometrical tachyons [17], where the slow roll could be achieved by taking $N$ to be sufficiently large. In what follows, however, we shall focus on inflationary models which use a single geometrical tachyon. It would be convenient to represent the DBI action (16) in the following form,

\[ \tilde{S}_p = -\int d^4 x \sqrt{-g} \sqrt{1 + \alpha' (\partial_{\mu} \tilde{\Phi})^2}, \]

where $\tilde{\Phi} \equiv -T/\sqrt{\alpha'}$, and $g_{\mu\nu}$, a four-dimensional metric introduced on the D3-brane is taken to be the FRW metric:

\[ ds^2 = -dt^2 + a^2(t)[dr^2 + r^2 d\Omega_2^2]. \]

The action (18) taken with the Einstein-Hilbert action

\[ S_E = \frac{M_p^2}{2} \int d^4x \sqrt{-g} R, \]

lead to the following evolution equations,

\[ \ddot{\tilde{\Phi}} = -(1 - \alpha' \dot{\tilde{\Phi}}^2)(M_p^2 \frac{\dot{\Phi}}{V} + 3H \dot{\tilde{\Phi}}), \]

\[ H^2 = \frac{8\pi G}{3} \frac{\dot{\tilde{V}}(\tilde{\Phi})}{\sqrt{1 - \alpha' \dot{\tilde{\Phi}}^2}} \]

where $H(t) \equiv \dot{a}/a$ is the Hubble parameter. Upon using (slow-roll) condition $\alpha' \dot{\tilde{\Phi}}^2 \ll 1$, (22) can be cast in the standard form

\[ H^2 = \frac{8\pi G}{3} \left| V_{eff}(\psi) + \frac{\dot{\psi}^2}{2} \right|, \]

where $\psi$ is some normalized version of $\tilde{\Phi}$ and $V_{eff}$ is the corresponding potential. In what follows we shall present the explicit form of the effective potential.

We now turn our attention to inflation driven by geometrical tachyon. The slow roll parameters are given by

\[ \epsilon = \frac{M_p^2}{2} \left( \frac{V_{eff}}{V_{eff}} \right)^2, \quad \eta = M_p^2 \frac{V''_{eff}}{V_{eff}} \]

In order to draw required number of e-folds, $\epsilon$ and $|\eta|$ have to be small. These parameters are also related to
the spectral index $n_s$ of the scalar density fluctuations in the early universe as

$$n_s - 1 \simeq -6\epsilon + 2\eta,$$  
(25)

The observations from the CMB measurements [21] require that $n_s \approx 0.95$ which also implies that $\epsilon \ll 1$, $|\eta| \ll 1$. In addition to this, $\epsilon$ and $\eta$ are constrained by the observation on amplitude of the primordial density perturbations,

$$\delta_s = \frac{1}{\pi \sqrt{15}} \frac{V_{eff}^{3/2}}{M_p^3 V'_{eff}} \lesssim 10^{-5}$$  
(26)

at horizon crossing. In the discussion to follow, we shall investigate the particular cases of the effective potential imposing conditions on the angular momentum $l$ of the D3 brane.

(a) case $l < \sqrt{k}l_s$

First we consider the case, $l < \sqrt{k}l_s$. From (17) we observe that $\dot{V}(R)/\tau_p \to 1$ as $R \to \infty$, while $\dot{V}(R)/\tau_p \to l/\sqrt{k}l_s$ as $R \to 0$. So, for $l < \sqrt{k}l_s$ the tachyon rolls from $R \sim \infty$ ($T \sim \infty$) to $R \sim 0$ ($T \sim -\infty$) (note that $\dot{V}(R)$ is a monotonic function), and the inflation is expected to occur at $R \gg \sqrt{k}l_s$.

To find $V_{eff}(\psi)$, we rewrite Eq.(17) in the form,

$$\dot{V}(\Phi) = \tau_3[1 + \left(\frac{l^2 - kl_s^2}{2\alpha_s \phi^2} + O\left(\frac{1}{\phi^4}\right)\right]$$  
(27)

for $R \gg \sqrt{k}l_s$ (or $\Phi \gg 1$). Substituting (27) into (22), and assuming $\alpha_s \phi \ll 1$, we find that,

$$H^2 = \frac{8\pi G}{3} \left[\frac{M_s^4}{(2\pi)^3 g_s} \left(1 + \left(\frac{l^2 - kl_s^2}{2\phi^2}\right) + \frac{1}{(2\pi)^3 g_s} \left(1 + \frac{1}{\phi^2}\right)\right)\right],$$  
(28)

where $\Phi \equiv M_s \Phi$ and we have used $\tau_3 = M_s^4/(2\pi)^3 g_s$. Comparing (28) with (23), one obtains the expression for the effective potential,

$$V_{eff}(\psi) = \frac{M_s^4}{(2\pi)^3 g_s} \left[1 + \left(\frac{l^2 - kl_s^2}{2\phi^2}\right)\right],$$  
(29)

where $\psi$ is defined by $\psi = \Phi/\sqrt{(2\pi)^3 g_s}$. The slow-roll parameters are now given by

$$\epsilon \simeq XY^3, \quad \eta \simeq -6XY^2,$$  
(30)

where

$$X = \frac{(2\pi)^3 g_s}{2k} \frac{M_s^2}{M_p^2} \Delta_1(l), \quad Y = \frac{kM_s^2}{\Phi^2} \Delta_1(l),$$  
(31)

and $\Delta_1(l)$ is defined as, $\Delta_1(l) = (kl_s^2 - l^2)/kl_s^2$.

Let us next estimate the slow-roll parameters. We make the following plausible choice[17]

$$\frac{\Phi^2}{kM_p^2} \sim 10, \quad k = 1.$$  
(32)

Since $0 < \Delta_1(l) < 1$ for $l < \sqrt{k}l_s$, we see from (31) that $Y (= 6\epsilon/|\eta|) \simeq 0.1$, if we take $l$ so as to satisfy $\Delta_1(l) \sim O(1)$. For $X \sim 1$, we have

$$\epsilon \sim 0.002, \quad \eta \sim -0.01,$$  
(33)

which agrees with (25). Also from (29) one finds that (26) can be written as

$$\delta_s = \frac{(2\pi)^2}{\sqrt{75}} \frac{M_s}{M_p} \frac{1}{\sqrt{k\Delta_1(l)}} \frac{1}{XY^{3/2}} \sim 10^{-5},$$  
(34)

which then implies

$$\frac{M_s}{M_p} \sim 10^{-7}.$$  
(35)

It is interesting to note that the tensor to scalar ratio of perturbations, $r = 16\epsilon \simeq 0.032$ is quite low in the model. Finally, $X \sim 1$ together with (35) implies

$$g_s \sim 10^{-16},$$  
(36)

which is the realistic decoupling limit considered in "Little String Theories" [LST] around one TeV [22, 23]. However, (36) is not compatible with the value of $M_s/M_p \sim 10^{-16}$ obtained in these theories [22, 23].

(b) case $l > \sqrt{k}l_s$ with inflation occurring at $R < \sqrt{k}l_s$

Let us now consider, $l > \sqrt{k}l_s$. In this case the tachyon rolls from $R \sim 0$ ($T \sim -\infty$) to $R \sim \infty$ ($T \sim \infty$). Thus for $l > \sqrt{k}l_s$, inflation is expected to occur around $R < \sqrt{k}l_s$ or possibly at $R \sim \sqrt{k}l_s$ if $\sqrt{k}l_s$ is not so large.

In the region, $R \ll \sqrt{k}l_s$, $T(R)$ can be approximated by $T(R) \sim \sqrt{k}l_s$ in $R/\sqrt{k}l_s$; from (17), we obtain

$$\dot{V}(\Phi) \equiv \tau_1 \frac{l}{\sqrt{k}l_s} \left[1 - \frac{1}{2} \Delta_2(l) e^{\frac{\sqrt{\Phi}}{\phi^2}}\right],$$  
(37)

where $\Delta_2(l) \equiv (l^2 - \sqrt{k}l_s^2)/l^2$ and $\tilde{\Phi} = T/\sqrt{\alpha}$. Substituting (37) into (22) and comparing it with (23) one finds,

$$V_{eff}(\psi) = \frac{M_s^4}{(2\pi)^3 g_s} \left[1 - \frac{1}{2} \Delta_2(l) e^{\beta(l) \psi/M_s}\right].$$  
(38)

where $\beta(l) = 2(2\pi)^3 g_s^{-1/2}(l/\sqrt{k}l_s)^{1/2}$ and $\psi = ([2\pi)^{-3} g_s^{-1}(l/\sqrt{k}l_s)^{1/2}\Phi$ with $\Phi \equiv M_s \Phi$. In this case the slow-roll parameters assume the form,

$$\epsilon \simeq XY^2, \quad \eta \simeq -4XY,$$  
(39)
where
\[
X = \frac{(2\pi)^3 g_s M_p^2}{2k} \sqrt{kl_s} \, , \quad Y = \Delta_2(l) e^{\beta(l)\varphi/M_s} \, . \tag{40}
\]

Let us note that \(0 < \Delta_2(l) < 1\) for \(l > \sqrt{kl_s}\) and therefore \(Y = 4\epsilon/|\eta| \to 0\) as \(\beta(l)\varphi/M_s (\simeq \Phi/M_s) \to -\infty\).

Now we estimate \(\epsilon\) and \(\eta\) with the following assumption,
\[
\frac{R}{\sqrt{kl_s}} \sim e^{\frac{\varphi}{\sqrt{kl_s}}} \sim \frac{1}{\sqrt{10}} \, , \quad k = 1, \tag{41}
\]

Eqs.(32) and (40) imply that, \(Y = 4\epsilon/|\eta| \sim 1/\sqrt{10}\) if \(\Delta_2(l) \sim O(1)\). Thus, if we set, \(X \sim 2 \times 10^{-2}\), the slow-roll parameters become
\[
\epsilon \sim 0.002, \quad \eta \sim -0.024, \tag{42}
\]

which agrees with (25). Turing to (26), one can show that \(\delta_s\) can be written as
\[
\delta_s \simeq -\frac{1}{2\pi\sqrt{75}} M_s \frac{1}{XY} \tag{43}
\]

in the present case, \(\delta_s \sim 10^{-5}\) implies
\[
\frac{M_s}{M_p} \sim 10^{-6}. \tag{44}
\]

And \(X \sim 10^{-2}\) together with (44) gives
\[
g_s \sim \tilde{l} \times 10^{-16} \, , \quad (\tilde{l} \equiv l/\sqrt{kl_s}) \, . \tag{45}
\]

So, differently from the case (a), \(g_s\) can not be determined in the case (b).

(c) case \(l > \sqrt{kl_s}\) with inflation occurring at \(R \sim \sqrt{kl_s}\)

We finally consider the case, \(l > \sqrt{kl_s}\) again, but assume that inflation occurs around \(R \sim \sqrt{kl_s}\) this time. To find \(\tilde{V}(\tilde{\Phi})\), we first set
\[
R \sim \sqrt{kl_s} + \eta, \quad (\eta \ll \sqrt{kl_s}). \tag{46}
\]

Then we find from (15)
\[
T \sim \sqrt{2kl_s}[1 + \frac{1}{\sqrt{2}} \ln(\sqrt{2} - 1) + x - \frac{1}{4}x^2], \tag{47}
\]

where \(x \equiv \eta/\sqrt{kl_s}\). Using then Eq.(17), we find that
\[
\tilde{V}(T) \simeq \tau_p \left(1 + \frac{(l/\sqrt{kl_s})^2}{2}\right)^{1/2} \left[1 - \frac{1}{2} \Delta_3(l)x + \frac{1}{2} \Delta_3(l) \frac{[k l_s^2 + \frac{1}{2}(l^2 - k l_s^2)]}{(l^2 + k l_s^2)} x^2\right], \tag{48}
\]

where \(\Delta_3(l) \equiv (l^2 - k l_s^2)/(l^2 + k l_s^2)\). Eq.(47) can be solved by
\[
x = C(\tilde{\Phi}) + \frac{1}{4} C^2(\tilde{\Phi}), \tag{49}
\]

where \(C(\tilde{\Phi})\) is defined by
\[
C(\tilde{\Phi}) = \frac{1}{\sqrt{2k}} (\tilde{\Phi} - \tilde{\Phi}_0) \tag{50}
\]

\[
\tilde{V}(\tilde{\Phi}) = \tau_p \left(\frac{1 + \tilde{l}^2}{2}\right)^{1/2} \left[1 - \frac{1}{2\sqrt{2k}} \Delta_3(l)(\tilde{\Phi} - \tilde{\Phi}_0) + \frac{1}{8k} \Delta_3(l)(1 + \tilde{l}^2)^{-1}(\tilde{\Phi} - \tilde{\Phi}_0)^2\right]. \tag{51}
\]

where \(\tilde{l} \equiv l/\sqrt{kl_s}\). Finally from (22) and (23) one finds
To estimate the unknown parameters, first consider the ratios of scales, given by \( \eta \sim \frac{(2\pi)^3 g_s M_p}{2\sqrt{2k} M_s^2 (1 + \hat{l}^2)^{3/2}} \), and
\[
\delta_s \sim -\frac{2\sqrt{2k}}{\pi \sqrt{s}} \frac{M_p^2}{(2\pi)^3 g_s M_s^2} \frac{1}{(1 + \hat{l}^2)^{1/2}} \Delta_3^{-1}(l). \tag{54}
\]
Also from (53) one finds
\[
\frac{4\epsilon}{\eta} = (1 + \hat{l}^2) \Delta_3(l). \tag{55}
\]
To estimate the unknown parameters, first consider the case \( l \sim \sqrt{k} l_s \), where \( \Delta_3(l) \to 0 \) and \( (1 + \hat{l}^2) \to 2 \). But in this case, \( 4\epsilon/\eta \) goes to zero (also note that \( \eta \) is positive in the case (c)) and the condition (25) can never be satisfied.

We next consider, \( l \gg \sqrt{k} l_s \), which gives rise to \( 4\epsilon/\eta \approx \hat{l}^2 \gg 1 \), and as a result, \( \eta \) is negligibly small as compared to \( \epsilon \). Thus using Eq.(25) we have the following estimates,
\[
\epsilon \sim 0.01, \quad \eta \sim 0, \tag{56}
\]
and from \( \epsilon \sim 20 \times g_s (M_p/M_s)^2 / \hat{l} \sim 10^{-2} \) and \( \delta_s \sim -3 \times 10^{-4} (M_p/M_s)^3 \hat{l} / g_s \sim 2 \times 10^{-5} \) we obtain the ratio of scales,
\[
\frac{M_p}{M_s} \sim 3 \times 10^4, \quad \hat{l} \sim \frac{l}{\sqrt{k} l_s} \sim g_s \times 10^{12}. \tag{57}
\]

We again note that the tensor to scalar ratio of perturbations given by \( r \sim 0.16 \) in the present case is consistent with observations. Eq.(58) may be thought as the equation for the number \( N \) of the scalar fields appearing in the model of assisted inflation [17]. In [14], it was shown that the condition \( \delta_s \sim 2 \times 10^{-5} \) implies \( N \sim (2\pi)^3 g_s \times 10^{10} \). In the case, \( l \gg \sqrt{k} l_s \) and the factor \( \hat{l} \) plays the similar role as the number of the scalar fields (or the number of D3-branes) of the assisted inflation [17].

Before closing this section, we wish to discuss two important aspect of a realistic scenario that might come out of model under consideration. The first is related to back reaction on the D3-brane worldvolume. We stress that the background NS5-branes are much heavier than D3-branes as the mass of NS5-branes is of the order of \( \frac{1}{g_s^2} \), while that of D-branes goes like \( \frac{1}{g_s} \) thereby in weak coupling limit, the NS5-branes are much heavier than the D3-branes and hence the dynamics is almost entirely governed by the gravitational attraction of the D3-branes into the core of NS5-branes. Thus for sufficiently weak coupling, the back reaction of these probe D3-branes can be ignored.

Let us now comment on the 4-dimensional physics of our model. The typical relationship between 4-dimensional Planck mass and the string mass is given by,
\[
\frac{M^2}{M_s^2} = \frac{v_0}{((2\pi)^7 g_s^2)} \tag{58}
\]
where \( v_0 = \text{volume}/(l_s)^6 \) is the overall volume factor of the compact CY3-fold. Presumably, to control stringy corrections, the CY3 must have a large volume, \( v_0 >> 1 \). We consider the asymmetric CY3 with the volume factor \( v_0 \sim 10^3 \) which is also preferred by a small value of string coupling as discussed in [13]. For instance, in the case (c), \( g_s \) can be estimated as \( \sim 10^{-5} \) from the value of \( M_p/M_s \) in (57). With that choice there is no backreaction on CY3-fold and there is no more than one D3-brane per unit CY measured in the units of string length.

**IV. CASE (C) AGAIN**

In Sec.III we have considered three cases of inflationary models where the D3-brane is moving around a stack of NS5-branes with certain radial and also a nonzero angular velocity. Among these models, of particular interest is the case (c) where the D3-brane is moving away from the fivebranes and inflation occurs at the radial distance, \( R \sim \sqrt{k} l_s \); the distance from the fivebranes as small as the string length. In fact, if the D-brane has an energy of the magnitude \( \varepsilon \sim l/\sqrt{2} \), \( R \sim \sqrt{k} l_s \) becomes a turning point of the radial motion, and the situation is that the inflation takes place at the minimum distance \( R \sim \sqrt{k} l_s \) from the fivebranes. For such a D-brane, the angular velocity at \( R \sim \sqrt{k} l_s \) is about \( \tilde{\theta} \sim 1/\sqrt{2k} l_s \) (see (10)), and therefore the time it takes for the D-brane to make a one turn will be about
\[
t \sim l_s \times 10, \tag{59}
\]
which becomes \( t \sim 10^{-38} s \) using Eq.(57).

Let us now estimate the time scale \( H^{-1} \). From (23) and (52) we see that
\[
H^{-1} \sim (2\pi)^{3/2} \frac{M_p}{M_s} \frac{g_s}{l_s^{1/2}}, \tag{60}
\]
which upon using (57) gives
\[
H^{-1} \sim \frac{M_p}{M_s^2} \times 10^{-5} \simeq 10^{-1} l_s \simeq (10^{15} \text{GeV})^{-1},
\] (61)
and therefore the expansion time \( \Delta t (\equiv 100H^{-1}) \) should be about
\[
\Delta t \sim l_s \times 10.
\] (62)

Eq. (61) shows that \( \Delta t \) is of the same order as the time it takes for the D-brane to circle around the fivebranes. So the natural scenario associated with the case (c) may be described as follows. A D-brane comes into a stack of NS5-branes with a certain energy \( \varepsilon \sim l/\sqrt{2} \) and a nonzero angular momentum \( l \). This D-brane approaches the fivebranes until the radial distance becomes \( R \sim \sqrt{kl_s} \), then it turns back to \( R \to \infty \) after circling the fivebranes about once within the distance \( R \sim \sqrt{kl_s} \) and during the revolution, the inflation occurs on the D-brane. In this process, the D-brane loses so much of its energy that it cannot escape to infinity. It gets captured by the fivebranes and starts orbiting around them. Also the orbiting D-brane keeps losing its energy and the angular momentum, and finally finds its stable state corresponding to \( \varepsilon \sim 1 \) and \( \hat{l} \sim 1 \), for which \( \tilde{V}(T) \) becomes \( \tilde{V}(T) \sim \tau_p \). This scenario is quite distinguished from the conventional scenarios based on the ordinary tachyon condensation.

In the case of the ordinary non-BPS D-brane, the tachyon rolls from the maximum of the tachyon potential \( V(T) = \tau_p \) to the minimum \( V(T) = 0 \) and in this process the minimum of the potential describes a tachyon vacuum corresponding to \( \varepsilon = 0 \), where there are no physical open string states. In the case (c), however, the end of the process corresponds to nearly a BPS D-brane (see [24]) with \( \varepsilon \sim 1 \) and \( \hat{l} \sim 1 \), on to which the open strings can end.

V. SUMMARY

In this letter, we have investigated inflationary models which use rolling geometrical tachyon as inflaton induced on a D3-brane moving in the vicinity of NS5-branes. These models are distinguished from those of the assisted inflation in the sense that they include a single geometrical tachyon to drive inflation. We have shown that the configuration with a single D3-brane with large angular momentum is effectively similar to the configuration with a large number of coincidence D3-branes as far as an inflation driven by rolling geometrical tachyon(s) is concerned and this enables us to obtain a scenario which can satisfy the requirement of slow-roll without introducing a large number of D3-branes.

Allowing for nonzero angular momentum \( l \) configurations, leads to an important feature in the tachyon potential, \( V(T) = \tau_p \sqrt{1 + (l^2/R^2)/(\sqrt{h(R(T)})} \). The potential \( V(T) \) crucially differs from the conventional geometrical tachyon potential \( V(T) \): It asymptotically approaches the value \( l \tau_p \) as \( R \to 0 \) where \( l \equiv l/\sqrt{kl_s} \), while it approaches \( \tau_p \) as \( R \to \infty \). Hence in case, \( l > \sqrt{kl_s} \) (or \( \hat{l} > 1 \)), the tachyon rolls from the maximum \( \tilde{V}(T) = \hat{l} \tau_p \) to the minimum \( \tilde{V}(T) = \tau_p \) \( (R = \infty) \). As a result, inflation occurs around \( R \sim 0 \), and thereafter the D-brane moves away and finds its stable orbit at \( R \sim \infty \). We have shown that amongst the various cases discussed in the present paper, the most interesting possibility corresponds to \( l > \sqrt{kl_s} \). In this case, inflation occurs around \( R \sim \sqrt{kl_s} \) and the end of the tachyon condensation corresponds to a BPS D-brane where the open strings can end. This is quite different from the conventional geometrical tachyon scenario where tachyon rolls from the maximum of its potential, \( V(T) = \tau_p \) \( (R = \infty) \) to the minimum, \( V(T) = 0 \) \( (R = 0) \). As a consequence, the D-brane finally decays into a tachyon vacuum as it approaches \( R = 0 \) where there are no physical open string states.

It is interesting to note that in the scenario discussed here, the estimate of expansion time \( H^{-1} \) corresponds to GUT scale which is, of course, related to the assumptions we made in the text. For instance, in the case (c), namely, \( l > \sqrt{kl_s} \), we assumed that the inflation occurs at the turning point \( R \sim \sqrt{kl_s} \) of the radial motion of the D-brane. However, it is also possible to consider other configurations with different turning points allowing us to generate different values of the expansion scale.

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