The \textit{N-Tachyon Assisted Inflation}

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\textbf{Abstract}

In continuation of the papers \texttt{hep-th/0505012} and \texttt{hep-th/0508101} we investigate the consequences when \textit{N} open-string tachyons roll down simultaneously. We demonstrate that the \textit{N-Tachyon} system coupled to gravity does indeed give rise to the assisted slow-roll inflation.
1 Introduction

The concept of inflation [1], the universe expanding with an accelerated expansion, is one of the most economic idea in the early universe cosmology. Having an initial inflationary phase solves many puzzles; the isotropy and homogeneity of the large scale structure, observed flatness of the universe and the primordial monopole problem etc. Inflation is also viewed as the most plausible source of the density perturbations in the universe. The conventional inflation models deal with a scalar field rolling down in a slowly varying potential in FRW spacetime [2]. Lately now, it is also being realised in string theory that the string landscape of cosmological vacua [3] does indeed possess many inflationary solutions. In fact there is a problem of being plenty. The string models which directly address the problem of inflation are the brane-inflation [4] models in which brane-antibrane force drives the modulus separating them. But the inflation is not the slow-roll one. The brane-antibrane inflation in warped compactification scenario [5], however, lacks precise knowledge of the inflaton potential in presence of fluxes. The recent idea of $N$ axion inflation can be found in [6, 7]. While there are open-string tachyon models which can also provide inflation, see [8, 9, 10, 11, 12], in general, these tachyon models are also plagued with the same large $\eta$ problem as the conventional models and are not favoured for slow-roll inflation, see [13]. There are number of various other attempts to provide stringy inflation models. These are for instance inflationary models based on branes intersecting at special angles [14], and also the D3/D7 models with fluxes where the distance modulus plays the role of inflaton field [15, 16]. The race-track inflation driven by closed string modulus could be found in [17].

A naive idea of assisted inflation was proposed in [18] in which $N$ scalar fields with identical potentials are considered. In doing so the Hubble parameter becomes of $O(\sqrt{N})$ thus increasing the gravitational frictional effects which effectively slow down the rolling of the scalar fields.\(^1\) In recent works [10, 12] we studied tachyon inflation models based on $N$ unstable D3-branes coupled to gravity. These $N$-tachyon models are found to be capable of providing solutions with slow-roll inflation provided $N$ is taken sufficiently large. In fact this requires a critical density of the branes on the compact Calabi-Yau 3-fold [12]. In this paper we refresh that work with a new proposal that our model [12] with $N$ non-BPS branes is actually an example of assisted inflation involving $N$ tachyons. Again the number density has to be large enough for these models to be cosmologically viable. The paper is organised as follows. We review the aspects of assisted inflation in the section-II. In section-III we study the $N$ tachyon assisted inflation. The conclusions are in section-IV.

\(^1\)The heterotic M-theory assisted inflation models are studied in [19].
2 Review: Assisted Inflation

The assisted inflation idea was proposed in [18] to overcome the large $\eta$ problem in scalar field driven inflation models. We review the main aspects of that work here. We consider a scalar field with potential $V(\phi)$ minimally coupled to Einstein gravity [2]

$$\int d^4x \sqrt{-g} \left[ \frac{M_p^2}{2} R - \frac{1}{2} (\partial \phi)^2 - V(\phi) \right]$$

(1)

where four-dimensional Planck mass $M_p$ is related to the Newton’s constant $G$ as $M_p^{-2} = 8\pi G$. Considering purely time-dependent field in a spatially flat FRW spacetime

$$ds^2 = -dt^2 + a(t)^2 \left( dr^2 + r^2 (d\theta^2 + \sin^2 \theta d\phi^2) \right)$$

(2)

the classical field equations can be written as

$$\ddot{\phi} = -V_{,\phi} - 3H \dot{\phi}$$

$$H^2 = \frac{8\pi G}{3} \left( \frac{\dot{\phi}^2}{2} + V \right)$$

(3)

where $H(t) \equiv \dot{a}/a$ is the Hubble parameter. This simple model has proved to be a prototype for explaining the mechanism of inflation in early universe. For example, if we take a quadratic potential $V = m^2 \phi^2 / 2$ and let the field roll down from some large initial value, the field will roll down to its minimum value $\phi = 0$ and so spacetime will inflate [2]. But the inflation has to be a slowly rolling one in order to fit with cosmological observations. That is the field $\phi$ must vary slowly such that there is a vanishing acceleration, $\ddot{\phi} \sim 0$. Under the slow-roll conditions, the time variation of $\phi$ gets related to the slope of the potential as

$$\dot{\phi} \simeq -\frac{V_{,\phi}}{3H}$$

So the potentials with gentle slop are preferred for a good inflation.

The standard slow-roll parameters are [20]

$$\epsilon \equiv \frac{M_p^2 V''}{2 V}, \quad \eta \equiv \frac{M_p^2 V'''}{V}$$

(4)

where primes are the derivatives with respect to $\phi$. For the slow-roll inflation both $\epsilon$ and $\eta$ will have to be small. Also these parameters are in turn related to the spectral index, $n_s$, of the scalar density fluctuations in the early universe as

$$n_s - 1 \simeq -6\epsilon + 2\eta.$$  

(5)

A nearly uniform power spectrum observed over a wide range of frequencies in the density perturbations in CMBR measurements [21], however, requires $n_s \simeq .95$. It can be achieved only if

$$\epsilon \ll 1, \quad \eta \ll 1.$$
These are some of the stringent bounds from cosmology which inflationary models have to comply with.

For single scalar field the power spectrum of the scalar curvature perturbations can be written as [20]

\[ P_R = \frac{1}{12\pi^2 M_p^6} \frac{V^3}{V'} \]  

(6)

The amplitude (or size) of these fluctuations are governed by

\[ \delta_H = \frac{2}{5} \sqrt{P_R} = \frac{1}{5\sqrt{3\pi M_p^3}} \frac{V^{3/2}}{V'} \leq 2 \times 10^{-5} \]  

(7)

The inequality on the right side of equation (7) indicates the COBE bound on the size of these perturbations at the beginning of the last 50 e-folds of inflation.\(^2\)

It can be easily seen that the models with quadratic potentials are not useful for inflation as they are plagued with so called \(\eta\)-problem. From the above we find that \(\eta \sim 2M_p^2/\langle \phi \rangle^2\). Therefore \(\eta\) will be small only if \(\phi\) has trans-Planckian vacuum expectation value at the time when inflation starts. But allowing the quantum fields to have trans-Planckian vevs will spoil the classical analysis and will involve quantum corrections. However, one can also see recent developments in [7].

An effective resolution of the large \(\eta\) problem can come from assisted inflation model [18]. The model involves large number of scalar fields \(\phi_i (i = 1, 2, \ldots, N)\). Let us demonstrate it here for a simple quadratic potential \(m^2 \phi^2/2\). We take the case where all scalars have the same mass \(m\) (i.e. potential). The scalar fields are taken to be noninteracting but are minimally coupled to gravity. The equations of motion are

\[ \ddot{\phi}_i = -m^2 \phi_i - 3H \dot{\phi}_i, \]

\[ H^2 = \frac{8\pi G}{3} \sum_{i=1}^{N} (V(\phi) + \frac{\dot{\phi}_i^2}{2}) . \]  

(8)

A particular solution exists where all fields are taken equal \(\phi_1 = \phi_2 = \ldots = \phi_N = \Phi\). The simplified equations become

\[ \ddot{\Phi} = -m^2 \Phi - 3H \dot{\Phi} \]

\[ H^2 = \frac{8\pi G}{3} N(V(\Phi) + \frac{\dot{\Phi}^2}{2}) \]  

(9)

One finds that eqs.(9) are having an effective Newton’s constant \(G.N\) when compared with eqs.(3). So we easily calculate that

\[ \epsilon = \eta = \frac{2M_p^2}{N\Phi^2} . \]

\(^2\)The number of e-folds, \(N_e\), during inflationary time interval \((t_f - t_i)\) are estimated as \(N_e = \frac{1}{3} \int_{t_i}^{t_f} H dt\). Our universe requires 50-60 e-folds of expansion in order to explain the present size of the observed large scale structure.
So if $N$ is sufficiently large enough $\Phi$ need not have trans-Planckian vev. It can also be seen that since $H \sim \mathcal{O}(\sqrt{N})$ it can provide slow-roll inflation since $[18]$

$$n_s - 1 = 2 \frac{\dot{H}}{H^2} \sim \mathcal{O}(\frac{1}{N}).$$

(10)

If larger is the value of $N$, more flat will be the observed spectrum. As an estimate, for

$$\frac{2M_p^2}{N\Phi^2} \sim .01$$

and the scalar mass

$$m \sim 10^{-3}M_p,$$

one can produce good slow-roll inflation. In fact, in Ref.[18] it was shown that the same phenomenon occurs even for exponentially fast potentials like $V = V_0 \exp(-\alpha \phi_i)$.

3 N-Tachyon Inflation

In this section we shall demonstrate that the assisted inflation can also occur for open string tachyons. It is rather natural for tachyons, unlike in the scalar field model above, to have the same masses $-M_s^2/2$ and also the identical potentials.

We consider $N$ unstable D3-branes distributed over a compact $CY_3$ with the six-volume parametrised as

$$V_{(6)} = (l_s)^6v_0$$

(11)

where $l_s = \sqrt{\alpha'} = 1/M_s$ is the string length. In the dilute ‘brane gas’ approximation (far separated branes in a large six-volume ($v_0 \gg 1$)) the unstable branes will have leading tachyon modes coming from stretched open strings attached to themselves only. Thus we are ignoring the tachyonic modes which come from the strings attached between a pair of far separated branes. Effectively we are assuming that the branes are non-interacting but interact through the gravity only. So the system is reduced to exactly $N$ tachyon fields with gravitational interaction. The effective low energy tachyon action can be written from $[22][23]$

$$-\sum_{i=1}^{N} \int d^4x V_i(T_i) \sqrt{-\det(g_{\mu\nu} + \alpha'\partial_{\mu}T_i \partial_{\nu}T_i)}$$

(12)

where $T_i$’s are the $N$ tachyon fields, $V_i(T_i)$ is the tachyon potential for $T_i$, and $g_{\mu\nu}$ is the pull-back of the spacetime metric. There is no summation over index $i$ within the integrand.

From the Sen’s conjectures the tachyon potential, $V(T)$, is symmetric and has a central maximum at $T = 0$ with a pair of global minima at $T \rightarrow \pm \infty$. The exact form of tachyon potential is not quite well known, although it is known to behave as

$$V(T) \sim e^{-\frac{T}{\sqrt{\alpha'}}}$$
near the tachyon vacuum (non-perturbative open string vacuum) at $T = \infty$. While near the perturbative vacuum at $T = 0$, the region of our interest, the potential can precisely be written as

$$V(T) = \sqrt{2} \tau_3 (1 - \frac{T^2}{4})$$

where $\tau_3 = \frac{1}{(2\pi)^3 g_s \alpha'^2}$ is the BPS D3-brane tension, $g_s$ being the string coupling constant. The potential (13) is valid for $|T| \ll 2$ only. It is like an inverted oscillator and the small excitations near $T = 0$ are tachyonic with negative (mass)$^2 = -\frac{M_s^2}{2}$.

In the flat spacetime background, the equations of motion for the purely time-dependent (homogeneous) tachyon fields can be written from (12)

$$\ddot{T}_1 = -(1 - \alpha' \dot{T}_1^2) \frac{M_s^2 dV_1}{V_1 dT_1},$$
$$\ddot{T}_2 = -(1 - \alpha' \dot{T}_2^2) \frac{M_s^2 dV_2}{V_2 dT_2},$$
$$\vdots$$

(14)

These equations are completely decoupled which is our primary assumption. So we can find a rolling tachyon solution where all $N$ tachyon fields roll down simultaneously. We can have $T_1(t) = T_2(t) = \cdots = T_N(t) = \phi(t)$ provided we also take $V_1(T_1) = V_2(T_2) = \cdots = V_N(T_N) = V(\phi)$. This is guaranteed from the known form of tachyon potential (13) valid in the perturbative open string vacuum. Under this simultaneous ansatz, the effective action (12) will simply become

$$-N \int d^4x V(\phi) \sqrt{-g} \sqrt{1 + \alpha' (\partial_\mu \phi)^2}. \tag{15}$$

Note that this was the action which we called as the fat unstable 3-brane in \cite{10,12}. We shall now couple this system to four-dimensional Einstein gravity

$$\frac{M_p^2}{2} \int d^4x \sqrt{-g} R. \tag{16}$$

The Planck mass $M_p$ is related to ten-dimensional string mass as

$$M_p^2 = \frac{v}{g_s} M_s^2, \tag{17}$$

where our convention has been

$$\kappa_{(10)}^2 = (2\pi)^7 g_s^2 \alpha'^4, \quad v \equiv v_0/(2\pi)^7. \tag{18}$$

For the effective four-dimensional analysis to remain valid, we must take string coupling $g_s \ll 1$ and also $v_0 \gg 1$ in order to suppress the stringy corrections. From the relation (17), this always brings the string mass $M_s$ much lower than $M_p$, which is a common assumption in large volume ($v_0 \gg 1$) string compactifications.
In order to study cosmological solutions we take FRW ansatz and consider purely
time dependent tachyon fields. From (12) and (16) the combined gravity and tachyon
field equations, keeping the simultaneous ansatz as above, can be written as

\[ \ddot{\phi} = -(1 - \alpha' \dot{\phi}^2) \left( M_s^2 \frac{V_{\phi}}{V} + 3H \dot{\phi} \right) \]
\[ H^2 = \frac{8\pi G}{3} (1 - \frac{\phi^2}{4}) \sqrt{\frac{2\tau_3 N}{1 - \alpha' \dot{\phi}^2}} \] (19)

We again emphasize that these equations are valid only near the top of the tachyon
potential where \( \phi^2 \ll 4 \) and \( \dot{\phi}^2 \ll M_s^2 \), these follow from the tachyon potential (13) which
is valid only near \( T \sim 0 \). Using these approximations and keeping upto quadratic terms
in the field \( \phi \), the equations (19) can be approximated as

\[ \ddot{\phi} = \left( M_s^2 \phi - 3H \dot{\phi} \right) + O(\phi^3) \]
\[ H^2 = \frac{8\sqrt{2} \pi G \tau_3 N}{3} (1 - \frac{\phi^2}{4} + \frac{\alpha' \dot{\phi}^2}{2}) + O(\phi^4) . \] (20)

Note that the field \( \phi \) is dimensionless so far in our analysis. We can restore its canonical
mass dimension by the rescaling \( \phi \rightarrow \sqrt{\alpha'} \phi \). In which case, dropping the higher order
terms altogether, we get

\[ \ddot{\phi} = \frac{M_s^2 \phi}{2} - 3H \dot{\phi} \]
\[ H^2 = \frac{8\sqrt{2} \pi G N}{3} \left( V_{\text{eff}} \phi + \frac{\phi^2}{2} \right) \] (21)

where we defined

\[ V_{\text{eff}} = M_s^4 - M_s^2 \phi^2 \frac{4}{4}, \quad \tilde{N} \equiv \frac{\sqrt{2} N}{(2\pi)^3} \] (22)

These are the actual equations relevant for an assisted inflation, find the similarity with
eqs. (9). Comparing eqs. (9) and (21) we determine that the mass of the inflaton field
is the tachyon mass \(-M_s^2/2\). But unlike conventional quadratic potential \( m^2 \phi^2/2 \), the
potential \( V_{\text{eff}} \) has a maximum at \( \phi = 0 \). In any case, for very large value of \( N \), we will
have \( H = O(\sqrt{N}) \) and the slow-roll assisted inflation is possible. This is our main result.

4 Slow-roll parameters

In order to know the slow-roll parameters, it will be useful to define a new field,

\[ \psi \equiv \left( \frac{\tilde{N}}{g_s} \right)^{\frac{1}{2}} \phi \] (23)
In terms of $\psi$ the equations in (21) become

\[ \ddot{\psi} = \frac{M^2}{2} \psi - 3H \dot{\psi}, \]
\[ H^2 = \frac{8\pi G}{3} \left( V + \frac{\psi^2}{2} \right) \] (24)

where

\[ V = \frac{\tilde{N}}{g_s} M_s^4 - M_s^2 \frac{\psi^2}{4}. \] (25)

We must caution the reader here, though it appears that there is only one scalar field in eqs. (24), one must keep in mind that actually there are $N$ of them. The energy density for single unstable brane is $\sqrt{2} \tau_3$. Viewing (24) as a single scalar field system the inflationary slow-roll parameters remain

\[ \epsilon \equiv \frac{M_p^2 V'^2}{2 V^2}, \quad \eta \equiv \frac{M_p^2 V''}{V} \]

with derivatives with respect to the new field $\psi$. These become

\[ \epsilon = \frac{M_p^2}{2} \left( \frac{-M_s^2 \psi}{2V} \right)^2 \approx \frac{1}{2 N g_s} \left( \frac{\phi}{2M_s} \right)^2 \]
\[ \eta \approx -\frac{1}{2 N g_s} \] (26)

where we have used eq. (24) and the fact that $\phi \ll 2M_s$. Thus, already $\epsilon \ll 1$ provided $\tilde{N}/v = O(1/g_s)$. This is a quality reminiscent of topological inflation where inflation happen very close to the top of a potential. But in order to make $\eta \ll 1$, we need to have

\[ \frac{\tilde{N}}{v} \gg \frac{1}{g_s}. \] (27)

Note that unlike trans-Planckian vevs in scalar field inflation models, in the case of tachyon field the inflation happens when $\langle \phi \rangle \ll 2M_s \ll M_p$.

We now determine the size of scalar density fluctuations

\[ \delta_H = \frac{1}{\sqrt{75\pi M_p^3 V,\psi}}. \] (28)

Evaluating this quantity using eqs. (25) and (24), we find

\[ \delta_H = \frac{\sqrt{2} M_s g_s N}{\sqrt{75\pi M_p (2\pi)^3 v}} \frac{2M_s}{\phi} \] (29)

Thus, by taking

\[ \frac{M_s}{M_p} = \frac{g_s}{\sqrt{v}} \approx 10^{-6}, \quad \frac{\sqrt{2} g_s N}{(2\pi)^3 v} \approx 10 \] (30)
we estimate from (29)

\[ \delta_H \sim 3.7 \times 10^{-7} \left( \frac{2M_s}{\phi} \right) \]  

(31)

That is, the size of amplitudes will be \( \sim 1.8 \times 10^{-5} \) if we choose \( \frac{\phi}{2M_s} = .02 \). In summary, we can get the size of density fluctuations within the COBE bound \( \delta_H \leq 1.9 \times 10^{-5} \) at the start of the last 50 e-folds of inflation. Note that the inflation will end before \( \langle \phi \rangle \sim 2M_s \).

**Critical number density:**

From equations (27) and (30) it is clear that we need to have at least

\[ \frac{\sqrt{2}N}{(2\pi)^3 v} \geq \frac{10}{g_s} . \]  

(32)

Using the definition \( v \equiv v_0/(2\pi)^7 \) we get

\[ \frac{N}{v_0} \geq \frac{5\sqrt{2}}{(2\pi)^4 g_s} . \]  

(33)

This is what we call the critical number density of the unstable D3-branes on the compact manifold for which \( v_0 \gg 1 \). Anything less than that will not work. Note that \( v_0 \) is the absolute measure of the CY3 volume in \( (l_s)^6 \) units. So if we take reasonably smaller value of the string coupling, \( g_s \sim .1 \), we get an estimate

\[ \frac{N}{v_0} \geq .045 . \]  

(34)

The eqs. (33) and (34) do seem to define for us a very dilute gas of D3-branes on the CY3. Eq. (34) shows that we need to have less than one D3-brane per unit six-volume of the compact manifold measured in string length units. This justifies our earlier assumption that inter-brane interactions could be ignored and the \( N \) tachyon action eq. (12) can be taken as a direct sum of the action for individual tachyons.

**Some Remarks:**

In the above we took \( g_s \sim .1 \) while any value in the range \(.01 \leq g_s \leq .1 \) is allowed. But if \( g_s \) is taken further smaller than \(.01 \), the ratio \( N/v_0 \) will get far away from a dilute gas approximation since we would like to keep not more than one 3-brane per unit Calabi-Yau six-volume. It perhaps suggests us that we should then include the warping of the spacetime in our analysis something in line with the works [5, 11]. It was shown in Ref. [11] that the warping effects can produce slow-roll inflation. There are important issues of moduli stabilisation during inflation, the moduli fields could interfere with an otherwise successful inflation. We have ignored these effects in this paper as those should be addressed separately.
Let us also estimate the value of the Hubble parameter at the beginning of the inflation ($\psi \sim 0$). Using eqs. (24) and (30)

$$H(0)^2 \approx \frac{8\pi G \bar{N}}{3} M_s^4 \approx \frac{10}{3} M_s^2.$$  \hspace{1cm} (35)

It means that $H(0) \sim M_s$, which is good because for a low energy analysis to remain valid $H \leq M_s$. Also from (30), we get the string scale of $10^{13} GeV$ much smaller than $M_p$. This is an usual feature of large volume compactifications. For $g_s \sim .01$, we estimate the volume parameter $v \sim 10^8$. The absolute $CY_3$ volume $V(6)$ measured in string length units will be $(2\pi)^7 10^8 (l_s)^6$. Equating $V(6) \equiv (2\pi R)^6$, it gives the compactification radius $R \sim 29 l_s$, which makes the compact volume reasonably large to suppress higher order string effects.

5 Conclusion

We have studied in detail the $N$-tachyon inflation in an FRW spacetime. It is shown that when a large number of tachyon fields roll down simultaneously from the top of the tachyon potential, we get an assisted slow-roll inflation. This is also the concrete example of assisted inflation in string theory involving only tachyons. The results of this work match with our previous work [12]. Particularly, we need to have a dilute number density of unstable D3-branes on the compactified manifold such that there are less than one D3-brane per unit string size volume of the compact manifold. We have not considered the effects of warping in the throat regions of the compact manifold. If these effects are considered together with sizable number of unstable D3-branes, we will be assured of slow-roll inflation with desired properties so as to fit the cosmological bounds.

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Note added: After we reported this work we came to know of Ref. [24] where the authors study assisted inflation with exponential tachyon potentials, which is valid for large values of tachyon field $T$. While we have considered the potential with a central maximum which is valid near small $T$. 

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