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On the cosmological relevance of the tachyon

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Abstract

We analyse of the effective action of the tachyon field on a D-brane, of both bosonic as well as superstring theory. We find that the non-standard kinetic term of the tachyon field requires a correction to the Born–Infeld type Lagrangian. The cosmological significance of the resulting dynamics is explored. We also examine if the rolling tachyon can provide an effective cosmological constant and contrast its behaviour with quintessence.

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1. Introduction

The perturbative spectrum of open (super-)string around a (non-supersymmetric) D-brane contains a scalar which is tachyonic. The dynamics of this tachyon field has many unusual features. Sen conjectured [1] that the potential of this field has a minimum and when the tachyon condenses to its minimum, *all* open string excitations become infinitely heavy and hence inaccessible to the low energy observer. Of course, for such a thing to happen, the Lagrangian describing the tachyon dynamics ought to be of an unconventional form. In particular, the effective open string coupling is inversely proportional to the value of the tachyon potential [2]. Since the potential

vanishes at the minimum, the open string theory there is infinitely strongly coupled.

More recently, Sen [3] has pointed out that the time evolution of the tachyon field may have cosmological significance. Several authors have initiated an exploration of this idea [4–9]. (For other works on the cosmological relevance of the tachyon, see [10,11].) Thanks to unconventional form of the tachyon action, cosmology with tachyon matter can lead to results different from those obtained with a normal scalar field. For example, inflation driven by the tachyon differs from that driven by a conventional inflaton. It is also possible that tachyon matter could be a candidate for cold dark matter [3] as, at the minimum of its potential, it becomes a pressureless gas. Padmanabhan has recently emphasized [9] that it is always possible to construct a potential that leads to a given inflationary scenario. Therefore, it is all the more important to study cosmological consequences of potentials arising out of a fundamental theory such as string theory. With this motivation, we study inflation driven by the rolling

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of the tachyonic scalar on a (non-supersymmetric) D3-brane.

At this stage we should point out the limitations of the scope of our study. We assume, following Ref. [5], that the open string tachyon couples minimally to the graviton, but to *no* other closed string mode. This is rather a drastic assumption since there is no reason to ignore the coupling to, say, the dilaton. Moreover it is not known how to stabilize some of the other closed string moduli, such as the volume of compactification, the dynamics of which might affect that of the tachyon. We will, however, take the pragmatic point of view that some as yet unknown mechanism freezes these unwanted moduli and leaves us with an effective theory of the tachyon minimally coupled to Einstein–Hilbert gravity in four dimensions, as proposed by Gibbons [5].

The plan of the Letter is as follows. In the next section, we review the relevant aspects of tachyon dynamics following Refs. [5,6]. In Section 3, based on a comparison between the Born–Infeld action [2,12] and the tachyon effective action [13–15] obtained in the so-called background independent string field theory (B-SFT) [16], we propose a correction to the BI action used in the recent works [3–9]. The consequences of the extra term are worked out in Sections 4 and 5 for bosonic and superstring respectively. We find that the tachyon does not give enough inflation. Section 6 is an analysis of the slow roll conditions in these models. Perhaps not unexpectedly, these conditions are not met. In the concluding section, we examine whether the rolling tachyon could be an alternative to cosmological constant or quintessence.

2. Dynamics of tachyon matter

The effective field theory on a (non-supersymmetric) D-brane is described by the Born–Infeld Lagrangian

$$\mathcal{L}_{\text{BI}} = -V(T) \left\{ -\det[g_{\mu\nu} + 2\pi\ell_s^2(F_{\mu\nu} + \partial_\mu Y^i \partial_\nu Y^i + f(T)\partial_\mu T \partial_\nu T)] \right\}^{1/2}, \quad (1)$$

where, $1/2\pi\ell_s^2$ is the string tension, g the induced metric, F the gauge field strength and Y^i are scalar fields describing the transverse motion of the brane. Moreover, T is the tachyonic mode on the Dm -brane,

$V(T)$ the potential of T and $f(T)$ is a function of the tachyon. The particular dependence on the potential is a characteristic of D-brane physics [2] compatible with the conjectures of Sen [1]. The tachyon potential thus plays the role of the (inverse) effective coupling of this theory. That the derivatives of the tachyon field appears under the square-root was pointed out in [12] based on the T-duality symmetry of string theory. The arguments in [12] allow for a function $f(T)$ accompanying these derivative terms. We shall see this term is non-trivial as the kinetic term of the effective action following from string field theory does not have a canonical form.

Let us couple this system to gravity [5] and neglect all fields other than the tachyon. Since our objective is to study the cosmological significance of tachyon dynamics, we shall assume a Friedmann–Robertson–Walker metric

$$ds^2 = -dt^2 + a^2(t) \left(\frac{dr^2}{1-kr^2} + r^2 d\Omega_{m-1}^2 \right), \quad (2)$$

as well as a spatially homogeneous and isotropic tachyon field T which depends only on time. Let us specialize to the case $m = 3$ from now on. The dynamics of this system is given by the Lagrangian

$$\sqrt{-\det g} \left(\frac{1}{2\kappa^2} R(g) - V(T) \sqrt{1 - 2\pi\ell_s^2 f(T) (\dot{T})^2} \right), \quad (3)$$

where, the relation

$$\frac{\sqrt{8\pi}}{m_{\text{Pl}}} = \kappa = \frac{(2\pi)^2 g_s \ell_s}{\sqrt{2v}} \quad (4)$$

expresses the inverse reduced Planck mass κ in terms of the string tension and string coupling g_s . In (4), v is a dimensionless parameter corresponding to the volume of the 22- or 6-dimensional space transverse to the 3-brane. It depends on how the four-dimensional spacetime is realized from the 26- or 10-dimensional one.

The Lagrangian (3), with $f(T) = 1$, has recently been analyzed in Refs. [5,6]. The inclusion of $f(T)$ leads to some modifications. The Hamiltonian or

energy density of T is¹

$$\mathcal{H} = \frac{a^3(t)V(T)}{\sqrt{1 - 2\pi\ell_s^2 f(T)\dot{T}^2}} = a^3(t)\rho(T) \quad (5)$$

and its equation of motion is

$$\dot{\rho} = -3H(\rho + p), \quad (6)$$

where, $H = \dot{a}/a$ is the Hubble parameter and $p = -V(T)\sqrt{1 - 2\pi\ell_s^2 f(T)\dot{T}^2}$ is the pressure of the tachyonic fluid. The explicit expressions of p and ρ may be substituted to obtain the second order evolution equation of T .

The gravitational equations, on the other hand, are

$$H^2 = \frac{\kappa^2}{3}\rho = \frac{\kappa^2}{3} \frac{V(T)}{\sqrt{1 - 2\pi\ell_s^2 f(T)\dot{T}^2}}, \quad (7)$$

$$\begin{aligned} \dot{H} &= -\frac{\kappa^2}{2}(\rho + p) \\ &= -\frac{\kappa^2}{2} \frac{V(T)f(T)}{\sqrt{1 - 2\pi\ell_s^2 f(T)\dot{T}^2}} 2\pi\ell_s^2 \dot{T}^2. \end{aligned} \quad (8)$$

Eq. (7) may be rewritten as $\sqrt{2\pi\ell_s^2 f(T)\dot{T}^2} = \sqrt{(1 - \kappa^4 V^2(T)/9H^4)}$.

3. Inflation from tachyon on a bosonic D-brane

The effective action of the tachyon field $T(x)$ determined in the framework of bosonic B-SFT is [13, 14]

$$S_B = \tau_3 \int d^4x (\ell_s^2 e^{-T} \partial_\mu T \partial^\mu T + (T + 1)e^{-T}), \quad (9)$$

where the normalization factor

$$\tau_3 = \frac{1}{(2\pi\ell_s^2)^2 2\pi g_s} = \frac{\pi^3 g_s^3}{8v^2} m_{\text{Pl}}^4 \quad (10)$$

is the tension of the D3-brane [19]. In our convention, the tachyon field $T(x)$ is dimensionless. Hence $V(T) = \tau_3(T + 1)\exp(-T)$ has mass dimension four. It is parametrized by the dimensionless quantities v and g_s .

Upto two derivative terms, the action (9) is exact, in the sense that it incorporates the effect of all open string modes. The potential has a maximum corresponding to the D-brane at $T = 0$ and a minimum at $T = \infty$. (It is also unbounded for negative values of T , but that is a pathology of the bosonic theory.) Notice the e^{-T} factor due to which the kinetic term has a non-standard form. Hence the distance between the maximum and minimum (in field space) is finite.

It is of course possible to do a field redefinition to bring the kinetic term to the canonical form $\partial_\mu \phi \partial^\mu \phi$. The transformation is $\phi = 2e^{-T/2}$, in terms of which the Lagrangian in

$$\mathcal{L}_\phi = \tau_3 \left(\ell_s^2 \partial_\mu \phi \partial^\mu \phi - \frac{1}{4} \phi^2 \left(\ln \frac{\phi^2}{4} - 1 \right) \right). \quad (11)$$

We will come back to the description of the dynamics in terms of the field ϕ later on in this section.

There are an infinite number of corrections to the action (9). In particular, the Born-Infeld action (1) gives a good description of slowly varying fields, i.e., those for which the second and higher derivatives of the field can be ignored. Let us now treat the action (9) as an expansion of an action of the form (1)

$$S_{\text{BI}} = \tau_3 \int d^4x (1 + T)e^{-T} (1 + \pi\ell_s^2 f(T)\partial_\mu T \partial^\mu T). \quad (12)$$

From a comparison of Eqs. (9) and (12), in a region where both are good descriptions, it follows that $f(T) = 1/(\pi(1 + T))$. Put in another way, this factor is the necessary field redefinition to relate the tachyon fields in the Born-Infeld and B-SFT descriptions.

Once again, it is possible to do a field redefinition to soak up the factor of $f(T)$, but this does not map the entire field space $T(x) \in (-\infty, \infty)$ to that of the redefined tachyon $\tilde{T}(x) = \sqrt{1 + T(x)}$. Interestingly, it does, however, map the domain $-1 \leq T(x) \leq \infty$, which is precisely the region accessed by all the classical solutions, viz., the D-branes. Another amusing aspect of this field redefinition is that the potential in terms of $\tilde{T}(x)$ has a Gaussian factor, reminiscent of the superstring case. It should, however, be noted that the late time evolution of the tachyon field in either of these two cases does not seem to match the asymptotic behaviour analyzed by Sen [8]. Of course, a field redefinition could reproduce the correct asymptotics, but it would, necessarily, have to involve derivatives of the

¹ We set $k = 0$ for simplicity.

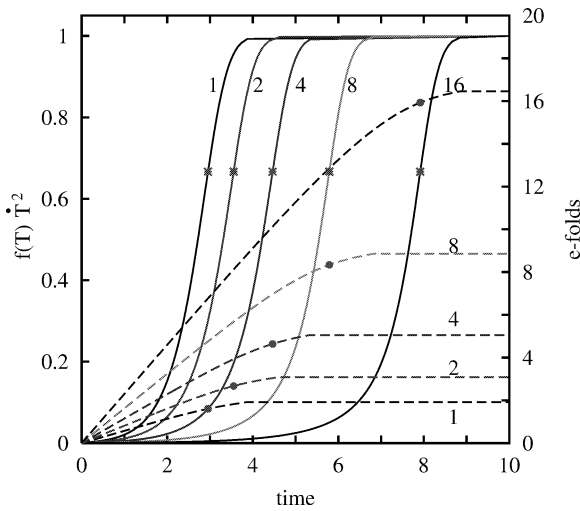


Fig. 1. The kinetic term $f(T)\dot{T}^2$ (solid lines) and e-foldings (dashed lines) vs time t for bosonic string. Numbers labelling the graphs correspond to values of $\kappa^2\tau_3$. In each case, the point at which the inflation ends has been marked.

tachyon field. Since we are mainly interested in the early time evolution where the time derivatives of the tachyon turn out to be small, the aforementioned derivative terms do not play a significant role. Furthermore, a field redefinition of the tachyon would not change the dynamics of the scale factor of the universe. Hence we may conclude that the results we obtain are not very sensitive to the large time behaviour.

One can now solve the equations of motion (7) and (8) numerically to determine the evolution of the tachyon and the Hubble parameter. Following [3], we choose the initial conditions such that the tachyon is infinitesimally displaced to the right of the maximum of the potential² and has zero velocity. The results depend on the tension τ_3 of the D3-brane, and are displayed in Fig. 1. In the same figure, we also plot the number of e-foldings

$$N_e(t) = \ln \frac{a(t_f)}{a(t_i)} = \int_{t_i}^{t_f} dt H(t) \quad (13)$$

² For the numerical computation, we set the string tension $1/2\pi\ell_s^2 = 1$. (The tension of the 3-brane τ_3 is therefore $1/2\pi g_s$.) The actual seed used for the initial value of the tachyon is $T \sim 0.1$, which is compatible with the requirement in [23].

as a function of time. It is clear from the plot that the effective kinetic term $2\pi\ell_s^2 f(T)\dot{T}^2$ initially remains very close to zero, during which the expansion takes place. Subsequently, the tachyon rolls very fast and the kinetic energy saturates to the maximum value of $2\pi\ell_s^2 f(T)\dot{T}^2 = 1$. The exit from the inflationary phase occurs very soon during the fast roll at the point $2\pi\ell_s^2 f(T)\dot{T}^2 = 2/3$.

Finally, we would like to mention that in terms of the field $\phi = 2e^{-T/2}$, which has a canonical kinetic term (11), one obtains $f(\phi) = 2/V(\phi)$. It is easy to check that one gets the same amount of e-folding from the evolution of ϕ . This is but a reaffirmation of the fact that field redefinitions cannot change the physical consequences of the dynamics.

4. Inflation from tachyon on a non-BPS D-brane

The analysis of the previous section can easily be extended to the case of the superstring. Recall, that the effective action of the tachyon field $\mathcal{T}(x)$ on a non-BPS D3-brane³ computed in supersymmetric B-SFT is [15]

$$S_F = \tau_3 \int d^4x (\ell_s^2 \ln 2e^{-T^2/4} \partial_\mu \mathcal{T} \partial^\mu \mathcal{T} + e^{-T^2/4}), \quad (14)$$

where, τ_3 , the tension of the brane is the overall normalization factor [21]. As in the bosonic case, this action is exact upto two derivatives. Using arguments identical to those in the previous section, we can write down the Born–Infeld form of the action

$$S_{BI} = \tau_3 \int d^4x e^{-T^2/4} \sqrt{1 - 2\ell_s^2 \ln 2(\dot{\mathcal{T}})^2}. \quad (15)$$

In other words, $f(\mathcal{T}) = (\ln 2)/\pi$.

As in the previous section, we now solve the equations of motion numerically to obtain the evolution of the tachyon and the Hubble parameter. Once again, the initial conditions are chosen such that the tachyon field starts with a small positive value and zero velocity. The results are displayed in Fig. 2. The qualitative features are similar to those in the bosonic case, although

³ One can also study the complex tachyon on a non-supersymmetric brane–antibrane pair [20]. The qualitative features of this system, however, are expected to be the same as those of a non-BPS brane.

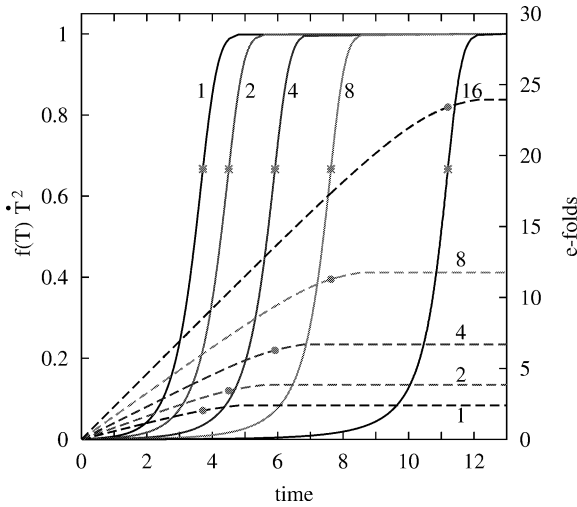


Fig. 2. The kinetic term $f(T)\dot{T}^2$ (solid lines) and e-foldings (dashed lines) vs time t for the superstring. The parameters refer to $\kappa^2\tau_3$.

for similar values of the brane tension τ_3 , the tachyon of the superstring has small velocity for a somewhat longer time, which results in a few more e-foldings.

5. Remarks on tachyon driven inflation

In order to generate enough inflation, it is necessary for the inflaton field to roll slowly enough. This is characterized by two dimensionless parameters ϵ and η . For the conventional inflaton with canonically normalized kinetic energy term, these are given by [22]

$$\epsilon = \frac{m_{\text{Pl}}^2}{2} \left(\frac{V'}{V} \right)^2, \quad \eta = m_{\text{Pl}}^2 \left(\frac{V''}{V} \right). \quad (16)$$

The conditions for slow roll inflation are $\epsilon \ll 1$ and $|\eta| \ll 1$. These formulas are not directly applicable to case of the tachyon, which has a non-standard action (12) or (15). Using the field redefinition discussed in Section 3, it is easy to show that

$$\begin{aligned} \epsilon &= \frac{m_{\text{Pl}}^2}{4\pi\ell_s^2 f(T)V(T)} \left(\frac{V'(T)}{V(T)} \right)^2, \\ \eta &= -\epsilon + \frac{m_{\text{Pl}}^2}{2\pi\ell_s^2 f(T)V(T)} \\ &\quad \times \left(\frac{V''(T)}{V(T)} - \frac{f'(T)V'(T)}{2f(T)V(T)} \right), \end{aligned} \quad (17)$$

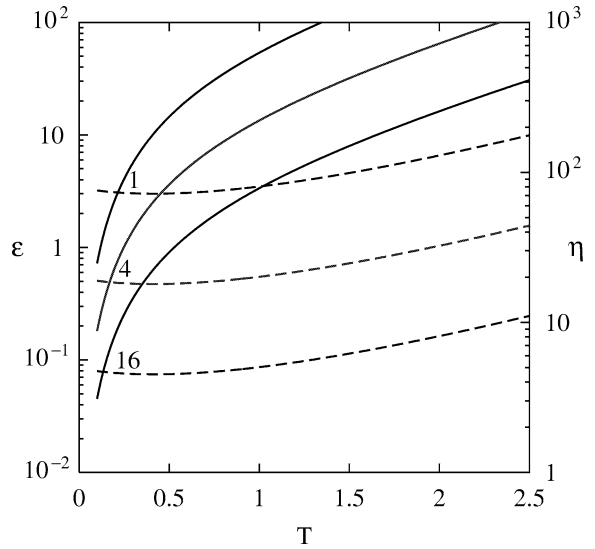


Fig. 3. Slow roll parameters ϵ (solid lines) and η (dashed lines) as a function of tachyon for the bosonic theory. Numbers labelling the graphs correspond to values of $\kappa^2\tau_3$.

are the appropriate parameters.

There is, nevertheless, a difference with conventional inflaton, especially if we are in the weak string coupling limit (small to moderate values of $\kappa^2\tau_3$). For the potential dictated by either the bosonic string or the superstring theory, the consequent inflation is not of slow roll type [17]. For example, while the slow roll parameter ϵ is small at the top of the potential, η is not. In fact, as Fig. 3 demonstrates for the bosonic theory,⁴ η is significantly larger than unity for the entire length of the evolution governed by tachyon dynamics. Thus, while the inflation satisfies one slow roll criterion (at least, initially), it fails to satisfy the other. The universe fails to inflate enough as a consequence of this.

Now this may not be all bad news, for the situation here is similar to the fast roll inflation [23]. It has been argued that such a fast roll inflation could have preceded the conventional slow roll inflation and could, in fact, have set the stage for the latter by suitably adjusting the initial conditions for slow roll. In this sense, it is satisfying to see that the inflation driven by the tachyon does not give rise to sufficient number

⁴ It is easy to see that while ϵ for the superstring tachyon starts arbitrarily close to zero, η , once again, is always large if one is in the weak coupling limit.

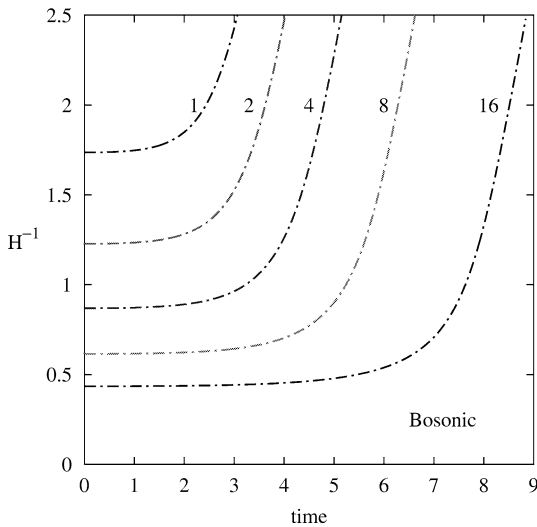


Fig. 4. The inverse Hubble parameter, in the bosonic theory, as a function of time. Numbers labelling the graphs correspond to values of $\kappa^2 \tau_3$.

of e-foldings, a job which will be relegated to the slow roll inflation that follows. After we exit the early epoch of tachyon driven fast roll inflation, the scale factor grows according to a power law (see Fig. 4). However, if a new scalar field takes over at this epoch and leads to slow roll inflation then the scale factor will continue to grow exponentially. Such details, though interesting, are beyond the scope of this work. Suffice is to note that there are enough candidates for the scalar among the open string (transverse scalars on the brane) and closed string moduli. Let us note in passing that in slow roll inflation, the ratio of tensor to scalar amplitudes is proportional to ϵ and observations demand it to be of the order of 10^{-2} or smaller. Since in our case ϵ is of the order of 10^{-1} , this ratio is larger than the observed value.⁵ However, if the tachyon driven fast roll inflation is followed by a slow roll inflation then it is conceivable that it may readjust this ratio to the desired value. One can also imagine modification of the present scenario of inflation on a single brane by considering a stack of multiple branes as considered in [6,11].

It is worth emphasizing another feature of inflation driven by the tachyon field. The tachyon potential is not flat. Indeed, the leading term in the effective

tachyon potentials (9) and (14) from string theory is quadratic at the maximum (as befits a tachyon). The kinetic term of the tachyon field, on the other hand, is non-trivial. This leads to all the difference in their dynamics.

6. Quintessence vs rolling tachyon

Recent observations indicate that there is probably an infinitesimally small positive vacuum energy density in the universe. This could be due to a cosmological constant. An alternative to this idea is that of quintessence, a conventional scalar field χ with a potential $V(\chi)$, such that χ is still rolling towards its minimum (see [24] and references therein). Generically the minimum of the potential is taken to be at $\chi = \infty$. In this section, we examine whether the rolling tachyon could be an answer to this problem.

The equation of state for quintessence is $p = \omega_\chi \rho$ with

$$-1 < \omega_\chi < -\frac{1}{3}. \quad (18)$$

The evolution of a scalar field in a homogeneous and isotropic universe described by the FRW metric (2) (with $k = 0$ for simplicity) is given by (6). This, along with the equation of state, determine the energy density ρ in terms of the scale factor $a(t)$. Eq. (7) can now be solved to get

$$\rho = \rho_0 \left(\frac{a_0}{a} \right)^{3(1+\omega)}, \quad a(t) = a_0 \left(\frac{t}{t_0} \right)^{2/(3(1+\omega))}. \quad (19)$$

It is clear from the power law behaviour of the scale factor that the Hubble parameter H is inversely proportional to time, indicating a curvature singularity at $t = 0$. This is because the curvature tensors are proportional to H and \dot{H} . As time evolves, curvature goes to zero and spacetime becomes flat.

The equations of motion for the tachyon field, when expressed in terms of the energy density and pressure, are identical to that of an ordinary scalar field. Its equation of state, however, is fundamentally different [3]. The parameter ω is not a constant but instead is a time-dependent function. More specifically,

$$\omega(T) = 2\pi \ell_s^2 f(T) \dot{T}^2 - 1, \quad (20)$$

⁵ L. Kofman, private communication.

as follows from the expressions of ρ and p in Section 2. With the initial conditions we have chosen in Sections 3 and 4, namely $T(0) = 0+$ and $\dot{T}(0) = 0$, ω starts with the value -1 . It is easy to see that the energy density at this stage is independent of the scale factor $a(t)$. Substituting this into the Einstein's equation (7) we find that the scale factor grows exponentially,

$$\rho = \rho_0, \quad a(t) = a_0 e^{t/t_0}, \quad (21)$$

with $\rho_0 = \tau_3$ and arbitrary constants a_0 and t_0 . This means that in the early epoch the scale factor grows exponentially giving a constant Hubble parameter H . Thus unlike the quintessence scalar field, tachyon dynamics does not lead to a curvature singularity as $t \rightarrow 0$.

As the system evolves, $\omega(T)$ moves away from -1 . We see from the Figs. 1 and 2 that initially $2\pi\ell_s^2 f(T)\dot{T}^2 \ll 1$ for some time. This in turn means that $\omega \simeq -1$ (but $\omega \neq -1$), is virtually a constant during this epoch. The system evolves according to Eq. (19), i.e., the scale factor grows according to a power law with a large exponent. Close to the end of the inflationary era, $f(T)\dot{T}^2$ grows rapidly, leading to an increase in the value of ω , which eventually settles down to $\omega = 0$.

In summary, we see that the early time development of the tachyon is quite different from that of the conventional quintessence scalar field. In particular, there is no curvature singularity as $t \rightarrow 0$; instead there is a constant scale factor and a flat metric. With the rolling of the tachyon, initial exponential growth of the scale factor switches over to a power law growth (see Fig. 4). Eventually, as the rolling tachyon picks up speed, growth of $a(t)$ as a function of t slows down and settles at $t^{2/3}$ indicating a transition to a tachyon (and other) matter dominated era. It will be interesting to investigate tachyon driven quintessence models in detail. We hope to come back to it in future.

Note added

Since an earlier version of our Letter, a number of authors have studied various aspects related to tachyon driven inflation. Most notably Ref. [17], (see also [18]), argued on general grounds that inflation driven by tachyon cannot produce enough expansion.

When we re-examined our results, it was found that incompatible conventions ($2\pi\ell_s^2 = 1$ and $\kappa = 1$) were inadvertently used. We correct this unfortunate error in the present version. The conclusions are modified in the light of the new result.

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