

Deforming tachyon kinks and tachyon potentials

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ABSTRACT: In this paper we investigate deformation of tachyon potentials and tachyon kink solutions. We consider the deformation of a DBI type action with gauge and tachyon fields living on D1-brane and D3-brane world-volume. We deform tachyon potentials to get other consistent tachyon potentials by using properly a deformation function depending on the gauge field components. Resolutions of singular tachyon kinks via deformation and applications of deformed tachyon potentials to scalar cosmology scenario are discussed.

KEYWORDS: Tachyon Condensation, Solitons Monopoles and Instantons, D-branes.

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1. Introduction

It has been conceived in the literature [1, 2] that the D-branes in string theory can be viewed as realizations of tachyon kinks [3–10]. Higher dimensional branes can be related to lower dimensional ones via descent relations [11]. Unstable Dp -branes may allow for stable or unstable $D(p-1)$ -brane solutions on their world-volume. This can be understood, as for example, a picture where stable type IIB $D(p-1)$ -branes with “right dimension” (odd dimension) may live on a unstable type IIB Dp -brane with “wrong dimension” (even dimension). Under this consideration, in order to establish the equivalence of tachyon kinks and $D(p-1)$ -branes, it is necessary to find tachyon kink solutions with finite energy per unit $(p-1)$ -volume [11]. Thus, it is interesting to investigate the properties of tachyon kink solutions in some tachyon models. In particular, it would be interesting to find a way of ‘resolving’ a singular tachyon kink solution whose energy is infinite via deformation of tachyon potentials. As we shall show later, we can introduce a deformation function associated with electromagnetic fields that can play this role. Thus, a singular tachyon kink of a theory describing a Dp -brane can be smoothed out by properly deforming the tachyon potential, living on the Dp -brane world-volume, into a new potential of a deformed theory describing a “deformed Dp -brane”. We show that while a deformed and smooth tachyon kink can confine electromagnetic fields on its world-volume, the singular tachyon kink of the original (non-deformed) theory cannot. Specially, for a D3-brane the electromagnetic field components can be localized everywhere on its world-volume. However, as we deform the tachyon potential, the world-volume of the deformed D3-brane may not localize all the components of the gauge field. This is because the smooth tachyon kink that we identify with a D2-brane living on the world-volume of the deformed D3-brane, now can trap

components of the gauge field. For instance, for some choice of parameters one might have a D2-brane confining a magnetic component or an electric component of the gauge field. This would characterize an alternative to compactification of a deformed D3-brane into a deformed D2-brane (the smooth tachyon kink) as a magnetically or an electrically charged D2-brane. A similar thing happens as one wraps strings living on the world-volume of a Dp-brane carrying a gauge field. A string will become magnetically or electrically charged depending on the direction of the Dp-brane world-volume manifold the string is wrapped.

In this paper we investigate deformed tachyon potentials and deformed tachyon kinks, following the lines of developments on tachyon kinks put forward in [3, 5]. Our main goal here is to address the aforementioned issue just by deforming “string theory motivated” tachyon potentials to get other consistent tachyon potentials by using properly the functional form of the gauge field components. We also discuss possible applications of deformed tachyon potentials to scalar cosmology scenario. The paper is organized as follows. In section 2 we consider the energy-momentum tensor and equations of motion for a DBI type action with gauge and tachyon fields living in a two-dimensional world-volume of a D1-brane. In section 3 we develop a general formalism to deform tachyon potentials by making use of a deformation function. The section 4 is devoted to a generalization of the DBI action by adding a Chern-Simons like term, which makes necessary to perform deformation of tachyon kinks/potentials with non-constant gauge field components. In section 5 we look for a tachyon kink solution and its deformed counterpart. The section 6 generalizes the treatment of section 3 for higher dimensional branes, specially D3-branes. In section 7 we make our final comments, considering some issues in tachyon cosmology.

2. Preliminaries

Let us consider a DBI type action describing tachyons and gauge fields on the world-volume of a non-BPS Dp-brane [12, 13]

$$S = -T_p \int d^{p+1}x V(T) \sqrt{-\det(\eta_{\mu\nu} + \partial_\mu T \partial_\nu T + F_{\mu\nu})}. \tag{2.1}$$

For electric and tachyon static fields living on the world-volume of a D1-brane (or D-string) we have the Lagrangian

$$\mathcal{L} = -T_1 V(T) \sqrt{1 + T'^2 - F_{01}^2}, \tag{2.2}$$

where $V(T)$ is the tachyon potential, $T \equiv T(x)$ and $E \equiv E_1(x) = F_{01}$. The equations of motion are

$$\left(\frac{V(T)T'}{\sqrt{1 + T'^2 - E^2}} \right)' - V'(T) \sqrt{1 + T'^2 - E^2} = 0, \tag{2.3}$$

$$\Pi' = \left(T_1 \frac{V(T)E}{\sqrt{1 + T'^2 - E^2}} \right)' = 0, \tag{2.4}$$

where the prime means derivative with respect to the argument of the function. The conjugate momentum Π is constant. The non-vanishing energy-momentum tensor components are

$$T_{00} = T_1 \frac{V(T)(1 + T'^2)}{\sqrt{1 + T'^2 - E^2}}, \tag{2.5}$$

$$T_{11} = -T_1 \frac{V(T)}{\sqrt{1 + T'^2 - E^2}}. \tag{2.6}$$

Notice that from equations (2.4) and (2.6) and from the conservation of the energy-momentum tensor $\partial_\mu T^{\mu\nu} = 0$ we find the following constraint

$$-T'_{11} = \left(\frac{\Pi}{E}\right)' = 0. \tag{2.7}$$

This constraint together with (2.4) forces the electric field E to be constant, unless we include some “external source” into equation (2.4). Indeed, this is achieved via Chern-Simons like terms, that we shall consider later, since we plan to focus on non-constant electric and magnetic fields.

3. Deformed tachyon potentials

In this section we address the issue of transforming the DBI type action for tachyons and gauge fields into another DBI type action with only a deformed tachyon field. In the “deformed theory” the gauge fields are encoded into a “deformation function”. We can transform the Lagrangian (2.2) into a deformed one as follows. Firstly we rewrite (2.2) as

$$\begin{aligned} \mathcal{L} &= -T_1 V(T) \sqrt{1 - E^2} \sqrt{1 + \frac{T'^2}{1 - E^2}} \\ &= -T_1 V(\tilde{T}) \sqrt{1 - E^2} \sqrt{1 + \tilde{T}'^2}. \end{aligned} \tag{3.1}$$

Let us now consider the transformation

$$\mathcal{L} \rightarrow \tilde{\mathcal{L}}, \quad V(\tilde{T}) \sqrt{1 - E^2} \rightarrow \tilde{V}(\tilde{T}) = \frac{V(T)}{\lambda \sqrt{1 - E^2}}. \tag{3.2}$$

This transformation leads to a new theory given by the Lagrangian

$$\tilde{\mathcal{L}} = -T_1 \tilde{V}(\tilde{T}) \sqrt{1 + \tilde{T}'^2}. \tag{3.3}$$

We regard this theory as “deformed theory”, with the deformed tachyon field

$$\tilde{T} = \pm \int \frac{dT}{\sqrt{1 - E^2}}, \tag{3.4}$$

and deformed tachyon potential $\tilde{V}(\tilde{T})$. The transformation (3.2) is justified by the fact we are requiring that both theories (3.1) and (3.3) maintain their energy-momentum components T_{11} and \tilde{T}_{11} conserved. The pressures T_{11} and \tilde{T}_{11} are constants related to each other via real parameter $\lambda > 0$. We can check this explicitly, i.e.,

$$\tilde{T}_{11} = -T_1 \frac{\tilde{V}(\tilde{T})}{\sqrt{1 + \tilde{T}'^2}} = -T_1 \frac{V(T)}{\lambda \sqrt{1 - E^2}} \frac{1}{\sqrt{1 + \frac{T'^2}{1 - E^2}}} = -T_1 \frac{V(T)}{\lambda \sqrt{1 - E^2 + T'^2}} = \frac{T_{11}}{\lambda}. \quad (3.5)$$

In this work we are interested in resolving tachyon solutions with nonvanishing pressure T_{11} , which show divergent energy density — the pressureless case is treated by Sen in ref. [14]. Thus the DBI type action with electric and tachyon fields can be transformed to another with only a deformed tachyon field. We summarize the transformations above as follows:

$$d\tilde{T} = \pm \frac{dT}{\sqrt{1 - E^2}} = \pm \frac{\lambda dT}{[f'(\tilde{T})]^2}, \quad (3.6)$$

$$\tilde{V}(\tilde{T}) = \frac{V(T)}{\lambda \sqrt{1 - E^2}} = \frac{V(T)}{[f'(\tilde{T})]^2}, \quad (3.7)$$

where $\tilde{T} = f^{-1}(T)$. Here we have followed the idea of deformed defects put forward in a former investigation [15] — see also [16, 17]. We regard the function $f(\tilde{T})$ as the deformation function defined as

$$f(\tilde{T}) = \int \lambda^{1/2} (1 - E^2)^{1/4} d\tilde{T}. \quad (3.8)$$

If E is constant this means nothing but $f(\tilde{T}) = \lambda^{1/2} (1 - E^2)^{1/4} \tilde{T}$. However, as we can see below, there are interesting deformations if we consider non-constant gauge fields.

An important feature we want to emphasize here is that we are deforming a singular kink with infinite energy to a smooth kink with finite energy. Note that this can be achieved because the energy densities of the deformed and non-deformed theories are given by

$$\tilde{T}_{00} = -\tilde{T}_{11} (1 + \tilde{T}'^2) = -\frac{T_{11}}{\lambda} \left(1 + \frac{T'^2}{f'(\tilde{T})^2} \right), \quad T_{00} = -T_{11} (1 + T'^2), \quad (3.9)$$

where we have used the equations (2.5), (2.6), (3.5) and the fact that the deformation function defines the map $\tilde{T} = f^{-1}(T)$. Supposing that \tilde{T} is a smooth tachyon kink, such that behaves asymptotically as $\tilde{T}'(|x| \rightarrow \tilde{x}_{\text{vac}}) = 0$, its vacuum energy density $\tilde{T}_{00} = -\tilde{T}_{11}$ is clearly finite. On the other hand, if T is a singular tachyon kink behaving as $T'(|x| \rightarrow \tilde{x}_{\text{vac}}) \rightarrow \infty$, with nonvanishing pressure, then $T_{00} \rightarrow \infty$. These statements can only be consistent with (3.9) if

$$\lim_{|x| \rightarrow \tilde{x}_{\text{vac}}} \frac{T'^2}{f'(\tilde{T})^2} = 0. \quad (3.10)$$

This is always true for the simple case $f'(\tilde{T}) \propto (T')^\gamma$, with $1 < \gamma < 2$. For $\tilde{T}'(|x| \rightarrow \tilde{x}_{\text{vac}}) \propto -|x| + \tilde{x}_{\text{vac}}$, we find the simple deformation function $f(x) \propto x (-|x| + \tilde{x}_{\text{vac}})^{\frac{2-\gamma}{1-\gamma}}$.

4. Brane within branes and Chern-Simons like couplings

There exists a remarkable feature in D-branes systems, that is the possibility of constructing D-branes with lower or higher dimension from others with a fixed dimension (see refs. [18] for comprehensive reviews, and references therein). This is possible by adding Chern-Simons like terms [19] to the action (2.1), which are couplings of the Dp -brane to background Ramond-Ramond (R-R) fields, e.g.,

$$S = -T_p \int d^{p+1}x V(T) \sqrt{-\det(\eta_{\mu\nu} + \partial_\mu T \partial_\nu T + F_{\mu\nu})} + T_p \int C_{(p-1)} \wedge F, \quad (4.1)$$

where $F = dA_{(1)}$ is the Abelian Born-Infeld 2-form field strength on the Dp -brane and $C_{(p-1)}$ is a R-R $(p-1)$ -form. The R-R potential C_{p-1} acts as a source of a “dissolved” $D(p-2)$ -brane living in the world-volume of a Dp -brane. The Chern-Simons like part of the action above is indeed a BF model [20, 21] which is topological, i.e., a metric independent action $\int_M B_q \wedge dA_{n-q-1}$ on a n -dimensional manifold M . Here, the manifold M is our $(p+1)$ -dimensional Dp -brane world-volume. Since the energy-momentum tensor is given by $(1/\sqrt{-g}) \delta S / \delta g^{\mu\nu}$, where $g_{\mu\nu} = \eta_{\mu\nu} + \partial_\mu T \partial_\nu T$, and the topological BF term is metric independent, the components (2.5) and (2.6) remain unchanged. Note that the equation of motion of the tachyon field does not change either, because this term does not depend on the tachyon field explicitly. For a D1-brane, we simply write down the Lagrangian

$$\mathcal{L} = -T_1 V(T) \sqrt{1 + T'^2 - F_{01}^2} + T_1 C_0(x) F_{01}. \quad (4.2)$$

Now $C_{(0)}$ is a scalar potential from the R-R sector of a type IIB string theory [19, 18]. The equation of motion of the electric field now reads

$$\Pi' = \left(T_1 \frac{V(T)E}{\sqrt{1 + T'^2 - E^2}} - T_1 C_0(x) \right)' = 0. \quad (4.3)$$

The non-constant background R-R scalar field $C_0(x)$ acts as a source of a $D(-1)$ -brane (a Dirichlet instanton in type IIB string theory). Since now we have $\Pi = -T_{11}E - T_1 C_0(x)$, the equation (4.3) does not force the electric field to be constant anymore; it suffices that $T_{11}E(x) = -T_1 C_0(x)$ to hold (4.3). This equation relates the electric field on the D1-brane with dissolved $D(-1)$ -branes living on the D1-brane, that agrees with earlier discussion about dissolved branes.

5. Deformed Tachyon Kinks

We are now ready to look for deformed tachyon kinks for non-constant electric field. This extends the analysis developed in [3, 5]. Let us first rewrite the equation (2.6) as

$$\varepsilon = \frac{1}{2} T'^2 + U(E, T). \quad (5.1)$$

This is similar to the total energy of a particle of mass equal to unity, with $\varepsilon = -1/2$ and potential energy

$$U(E, T) = -\frac{1}{2} \left[\frac{T_1 V(T)}{-T_{11}} \right]^2 - \frac{1}{2} E^2. \quad (5.2)$$

From the total energy (5.1) we can reduce the problem of finding tachyon kinks solutions to the problem of solving a first-order differential equation consistent with the tachyon equation (2.3). The solutions are found via equation

$$\int_0^T \frac{dT}{\sqrt{\left[\frac{T_1 V(T)}{-T_{11}}\right]^2 - (1 - E^2)}} = \pm x. \tag{5.3}$$

Now using the transformations (3.6) and (3.7) we can easily rewrite (5.3) as

$$\int_0^{\tilde{T}} \frac{d\tilde{T}}{\sqrt{\left[\frac{\lambda T_1 \tilde{V}(\tilde{T})}{-T_{11}}\right]^2 - 1}} = \pm x. \tag{5.4}$$

Notice that this is equivalent to work with the equation (3.5) for the \tilde{T}_{11} ($= T_{11}/\lambda$). By using the equation (5.4), up to difficulties with integrability, we should be able to find deformed tachyon kinks \tilde{T} from deformed tachyon potentials \tilde{V} .

Let us now consider some examples. First we consider the following tachyon potential in the DBI type action given in (4.2)

$$V(T) = \text{sech}(T/T_0). \tag{5.5}$$

The deformed potential is

$$\tilde{V}(\tilde{T}) = \frac{\text{sech}(T/T_0)}{[f'(\tilde{T})]^2} = \frac{\text{sech}(T/T_0)}{\sqrt{1 - E^2}}. \tag{5.6}$$

If the electric field E is constant and $\lambda=1$ as considered in [3, 5], the equation (5.4) turns out to be

$$\int_0^T \frac{dT}{\sqrt{\frac{E^2 T_1^2}{\Pi^2(1-E^2)} \text{sech}^2(T/T_0) - 1}} = \pm \sqrt{1 - E^2} x, \tag{5.7}$$

where we have used $T_{11} = \Pi/E$. Working out the equation (5.7) gives

$$T(x) = T_0 \text{arcsinh} \left[\sqrt{\frac{E^2 T_1^2}{\Pi^2(1 - E^2)} - 1} \sin \left(\frac{\sqrt{1 - E^2}}{T_0} x \right) \right], \tag{5.8}$$

which is the result found in [3, 5].

Now consider the electric field E is not constant and can be related to a deformation function according to (3.8). For a more general analysis let us assume $V(T) = \text{sech}^q(T/T_0)$. We choose, among the ones satisfying the criterion (3.10), the following deformation function:

$$T = f(\tilde{T}) = T_0 \text{arctanh}(\tilde{T}/T_0). \tag{5.9}$$

As we will see below, this function allows us to construct analytically a periodic tachyon kink that is consistent with a tachyon kink wrapped on a circle. Substituting this into equation (5.6) we find the deformed tachyon potential

$$\tilde{V}(\tilde{T}) = \frac{\text{sech}^q(T/T_0)}{[f'(\tilde{T})]^2} = \left(\sqrt{1 - \tanh^2(T/T_0)} \right)^q \left(1 - \frac{\tilde{T}^2}{T_0^2} \right)^2 = \left(1 - \frac{\tilde{T}^2}{T_0^2} \right)^{\frac{q}{2}+2}, \quad (5.10)$$

which has global minima at $\tilde{T} = \pm T_0$ and global maximum at $\tilde{T} = 0$, for $\frac{q}{2} + 2 > 0$. Notice that this potential is defined only in the interval $(-T_0, T_0)$. Substituting this potential into (5.4) we find

$$\int_0^{\tilde{T}} \frac{d\tilde{T}}{\sqrt{\frac{\lambda^2 T_1^2}{T_1^2} \left(1 - \frac{\tilde{T}^2}{T_0^2} \right)^{q+4} - 1}} = \pm x. \quad (5.11)$$

The general solution involves polynomial terms and elliptical functions, and is not necessary for our discussions here. The case $q = 1$ which recover the potential (5.5) is still hard to work with. Let us turn, however, to the simplest case, i.e., $q = -3$ to get the simple nontrivial solution

$$\tilde{T}(x) = T_0 \sqrt{1 - \frac{T_{11}^2}{\lambda^2 T_1^2} \frac{\tan(\omega x)}{\sqrt{\tan(\omega x)^2 + 1}}}, \quad (5.12)$$

where $\omega = \sqrt{\lambda^2 T_1^2 / T_0^2 T_{11}^2}$. This is a periodic tachyon kink whose period is π/ω — see figure 1, that is in accord with a tachyon kink on a circle of radius $R = 1/2\omega$ [11]. This tachyon kink is stable because the vacuum manifold of the deformed potential \tilde{V} consists of a pair of points $\pm T_0$, and the solution connects different vacua. This solution can be recognized as a stable D0-brane living on the D1-brane (or D-string) world-volume [11]. Since the D1-brane is unstable, it may decay leaving behind a stable D0-brane. This relation between the D1-brane and the D0-brane is an example of ‘descent relations’ that exist among D-branes, as first pointed out by Sen [11].

Using the deformation function (5.9) one can also obtain the tachyon kink T of the original theory (4.2) given by

$$T(x) = T_0 \text{arctanh} \left(\sqrt{1 - \frac{T_{11}^2}{\lambda^2 T_1^2} \frac{\tan(\omega x)}{\sqrt{\tan(\omega x)^2 + 1}}} \right). \quad (5.13)$$

This tachyon kink becomes singular for $T_{11}^2/\lambda^2 T_1^2 = 1/\omega^2 T_0^2 = 4R^2/T_0^2$ sufficiently small, i.e., $R \ll T_0/2$ — see figure 2. It is singular in the sense that its energy diverges, whereas the deformed tachyon kink is smooth and has finite energy. A singular tachyon kink wrapped on a circle S^1 with radius $R \ll T_0/2$ of a $D(p+1)$ -brane world-volume manifold $S^1 \times \mathbb{M}^{p+1}$ becomes “resolved” in the deformed theory with the same compactification radius. Thus we can think of the deformation process as a way of “smoothing out” singular tachyon kink solutions through gauge fields. This is similar to “brane resolution”, where singularity of branes can be resolved by turning on fluxes [22].

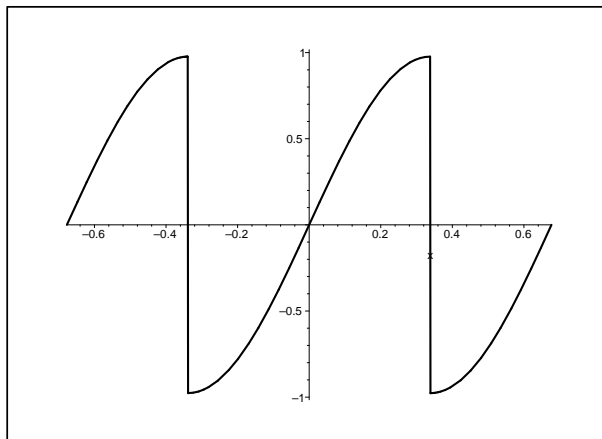


Figure 1: The deformed tachyon kink behavior in the interval $-\pi/\omega \leq x \leq \pi/\omega$, with $T_0 = 1$, $T_1 = 0.1$, $T_{11} = 10$ and $\lambda = 464$.

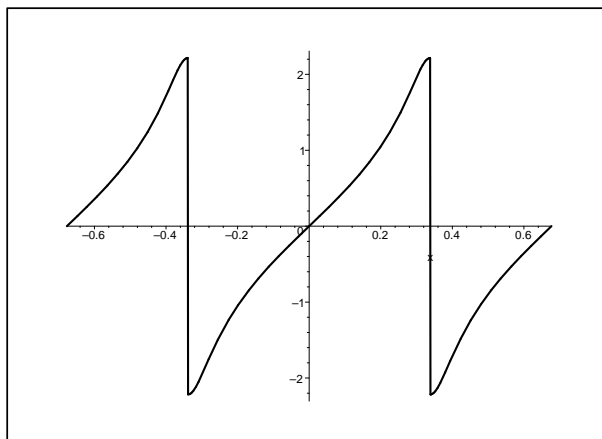


Figure 2: The singular tachyon kink behavior in the interval $-\pi/\omega \leq x \leq \pi/\omega$, with $T_0 = 1$, $T_1 = 0.1$, $T_{11} = 10$ and $\lambda = 464$.

The electric field can be calculated via eq. (3.8). Applying the deformation function $f(\tilde{T})$ defined as in (5.9) gives

$$E = \sqrt{1 - \lambda^{-2} \left(1 - \frac{\tilde{T}^2}{T_0^2}\right)^{-4}}. \tag{5.14}$$

Substituting the deformed tachyon kink solution (5.12) into (5.14), the electric field $E(x)$ develops the behavior depicted in figure 3. The electric field is confined to the new lower dimensional D0-brane — note it falls off fast along the transverse coordinate x . Both kink solution (5.12) and electric field function (5.14) impose constraints on the parameter λ , i.e.,

$$\frac{T_{11}}{T_1} \leq \lambda \leq \left(\frac{T_{11}}{T_1}\right)^{4/3}, \quad T_{11} > T_1. \tag{5.15}$$

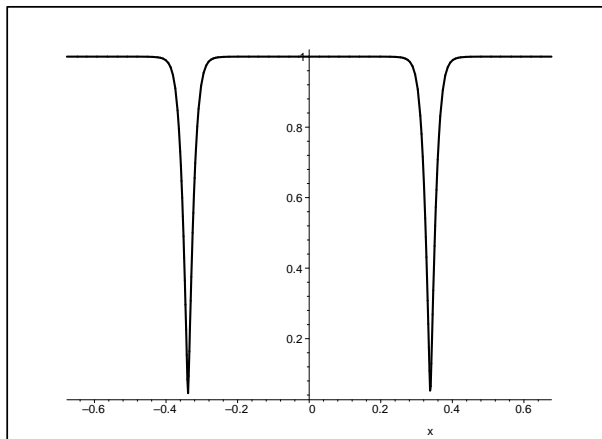


Figure 3: The electric field on the deformed tachyon kink, in the interval $-\pi/\omega \leq x \leq \pi/\omega$, for $T_0=1, T_1=0.1, T_{11}=10$ and $\lambda = 464$.

The electric field $E(x)$ we found above is localized either on the original D3-brane or on the D2-brane that appears as a deformed tachyon kink living on the D3-brane of the deformed theory. Note this is consistent with the fact that in the original D3-brane the tachyon kink is singular and cannot localize (or expel) the electric field. As a consequence, the electric field is present everywhere on the D3-brane world-volume and appears explicitly on the DBI action. On the other hand, the D3-brane world-volume of the the deformed theory has no explicit dependence on the electric field. This field now reappears as a localized field on the deformed tachyon kink (the D2-brane) and cannot be present everywhere on the *deformed* D3-brane world-volume anymore.

6. Extension to higher dimensional D-branes

Our previous analysis can be extended to Dp -branes with $p + 1$ dimensional world-volume. The Lagrangian for DBI and Chern-Simons like terms now reads

$$\mathcal{L} = -T_p V(T) \sqrt{-X} + \frac{1}{2(p-1)!} \epsilon^{\mu_1 \dots \mu_{p+1}} C_{\mu_1 \dots \mu_{p-1}} F_{\mu_p \mu_{p+1}}, \tag{6.1}$$

where in terms of components we write

$$-X \equiv 1 + (\vec{\nabla}T)^2 + (\vec{\nabla}T \cdot \vec{B})^2 - (\vec{\nabla}T \times \vec{E})^2 + \vec{B}^2 - \vec{E}^2 - (\vec{E} \cdot \vec{B})^2. \tag{6.2}$$

As a concrete example, let us focus on static field configurations on a D3-brane. The field equation for the tachyon living on the world-volume of a D3-brane is

$$\vec{\nabla} \cdot \left[\frac{V(T)}{\sqrt{-X}} \left((1 - \vec{E}^2) \vec{\nabla}T + (\vec{E} \cdot \vec{\nabla}T) \vec{E} + (\vec{B} \cdot \vec{\nabla}T) \vec{B} \right) \right] - V'(T) \sqrt{-X} = 0. \tag{6.3}$$

The conjugate momenta for the gauge fields, Π_i , can be written as

$$\Pi_1 = \frac{T_3 V(T)}{\sqrt{-X}} \left[(1 + (\vec{\nabla}T)^2) E_1 + (\vec{E} \cdot \vec{B}) B_1 - (\vec{E} \cdot \vec{\nabla}T)(\vec{\nabla}T)_1 \right] + \Pi_{1CS}, \quad (6.4)$$

$$\Pi_2 = \frac{T_3 V(T)}{\sqrt{-X}} \left[(\vec{E} \cdot \vec{B}) E_2 - (\vec{B} \cdot \vec{\nabla}T)(\vec{\nabla}T)_2 - B_2 \right] + \Pi_{2CS}, \quad (6.5)$$

$$\Pi_3 = \frac{T_3 V(T)}{\sqrt{-X}} \left[(\vec{E} \cdot \vec{B}) E_3 - (\vec{B} \cdot \vec{\nabla}T)(\vec{\nabla}T)_3 - B_3 \right] + \Pi_{3CS}, \quad (6.6)$$

$$\Pi_4 = \frac{T_3 V(T)}{\sqrt{-X}} \left[(1 + (\vec{\nabla}T)^2) E_3 + (\vec{E} \cdot \vec{B}) B_3 - (\vec{E} \cdot \vec{\nabla}T)(\vec{\nabla}T)_3 \right] + \Pi_{4CS}, \quad (6.7)$$

$$\Pi_5 = \frac{T_3 V(T)}{\sqrt{-X}} \left[(\vec{E} \cdot \vec{B}) E_1 - (\vec{B} \cdot \vec{\nabla}T)(\vec{\nabla}T)_1 - B_1 \right] + \Pi_{5CS}, \quad (6.8)$$

$$\Pi_6 = \frac{T_3 V(T)}{\sqrt{-X}} \left[(1 + (\vec{\nabla}T)^2) E_2 + (\vec{E} \cdot \vec{B}) B_2 - (\vec{E} \cdot \vec{\nabla}T)(\vec{\nabla}T)_2 \right] + \Pi_{6CS}, \quad (6.9)$$

where Π_{iCS} are contributions from the Chern-Simons like term. The components of the energy-momentum tensor turn out to be

$$T_{00} \equiv \rho = \frac{T_3 V(T)}{\sqrt{-X}} \left[1 + (\vec{\nabla}T)^2 + \vec{B}^2 + (\vec{B} \cdot \vec{\nabla}T)^2 \right], \quad (6.10)$$

$$T_{0i} \equiv P_i = \frac{T_3 V(T)}{\sqrt{-X}} \left[(\vec{B} \cdot \vec{\nabla}T)(\vec{E} \times \vec{\nabla}T)_i + (\vec{E} \times \vec{B})_i \right], \quad (6.11)$$

and

$$T_{ij} = \frac{T_3 V(T)}{\sqrt{-X}} \left[(\vec{\nabla}T)_i (\vec{\nabla}T)_j - E_i E_j - B_i B_j + (\vec{E} \times \vec{\nabla}T)_i (\vec{E} \times \vec{\nabla}T)_j - \left(1 - \vec{E}^2 + (\vec{\nabla}T)^2 \right) \delta_{ij} \right]. \quad (6.12)$$

Let us look for tachyon kink configurations, where we specialize ourselves on tachyon and gauge fields depending only on one coordinate, say x . The integrability of the system simplifies for such configurations. We can now write the following:

$$-X \equiv 1 + T'^2 + T'^2 (B_1^2 - E_2^2 - E_3^2) + \vec{B}^2 - \vec{E}^2 - (\vec{E} \cdot \vec{B})^2, \quad (6.13)$$

and the equations of motion for tachyon and gauge fields as

$$\left[\frac{V(T)}{\sqrt{-X}} (1 - \vec{E}^2 + E_1^2 + B_1^2) T' \right]' - V'(T) \sqrt{-X} = 0, \quad (6.14)$$

$$\Pi'_1 = \left\{ \frac{T_3 V(T)}{\sqrt{-X}} \left[E_1 + (\vec{E} \cdot \vec{B}) B_1 \right] + C_{32}(x) \right\}' = 0, \quad (6.15)$$

$$\Pi'_2 = \left\{ \frac{T_3 V(T)}{\sqrt{-X}} \left[(\vec{E} \cdot \vec{B}) E_2 - B_2 \right] + C_{20}(x) \right\}' = 0, \quad (6.16)$$

$$\Pi'_3 = \left\{ \frac{T_3 V(T)}{\sqrt{-X}} \left[(\vec{E} \cdot \vec{B}) E_3 - B_3 \right] + C_{03}(x) \right\}' = 0. \quad (6.17)$$

The system has energy density ρ and carries the linear momentum density P_i given by

$$T_{00} \equiv \rho = \frac{T_3 V(T)}{\sqrt{-X}} \left[1 + \vec{B}^2 + (1 + B_1^2) T'^2 \right], \quad (6.18)$$

$$T_{0i} \equiv P_i = \frac{T_3 V(T)}{\sqrt{-X}} \left[B_1 T' (\vec{E} \times \vec{\nabla} T)_i + (\vec{E} \times \vec{B})_i \right]. \quad (6.19)$$

The other non-vanishing components of the energy-momentum tensor are

$$T_{ij} = \frac{T_3 V(T)}{\sqrt{-X}} \left[T'^2 \delta_{i1} \delta_{j1} - E_i E_j - B_i B_j + (\vec{E} \times \vec{\nabla} T)_i (\vec{E} \times \vec{\nabla} T)_j - \left(1 - \vec{E}^2 + T'^2 \right) \delta_{ij} \right], \quad (6.20)$$

where $(\vec{E} \times \vec{\nabla} T)_i = T'(E_3 \delta_{i2} - E_2 \delta_{i3})$. The conservation of energy-momentum tensor, i.e., $\partial^\mu T_{\mu\nu} = 0$, implies the four conserved quantities $T_{01} = T_{10}$, T_{11} , $T_{12} = T_{21}$, and $T_{13} = T_{31}$. As in the one dimensional example above, now we shall make use of the quantity T_{11} , as well as the quantities T_{01} , T_{12} , T_{13} given explicitly by

$$T_{01} = -T_3 \frac{V(T)}{\sqrt{-X}} (B_2 E_3 - B_3 E_2), \quad (6.21)$$

$$T_{11} = -T_3 \frac{V(T)}{\sqrt{-X}} (1 + B_1^2 - E_2^2 - E_3^2), \quad (6.22)$$

$$T_{12} = -T_3 \frac{V(T)}{\sqrt{-X}} (E_1 E_2 + B_1 B_2), \quad (6.23)$$

$$T_{13} = -T_3 \frac{V(T)}{\sqrt{-X}} (E_1 E_3 + B_1 B_3). \quad (6.24)$$

These are constants of motion that will be useful for integrability of the system, as we shall see below. In addition to these four independent conserved quantities, there are other three constants coming from the Faraday's law, i.e., $\partial^\mu F_{\mu\nu}^* = 0$. The constants are the gauge field components E_2 , E_3 and B_1 , such that we are left with only $E_1(x)$, $B_2(x)$ and $B_3(x)$ as possible ‘‘dynamical’’ components. As a consequence of such constant components, from (6.22) we note that $V(T)/\sqrt{-X}$ is also constant.

The generalization of the discussion of section 3 to higher dimensions is straightforward. Specially in three spatial dimensions, to include new field components, we consider the following. Let us first rewrite (6.13) as

$$-X = \beta + \alpha T'^2, \quad (6.25)$$

where $\alpha = 1 + B_1^2 - E_2^2 - E_3^2$, and $\beta = 1 + \vec{B}^2 - \vec{E}^2 - (\vec{E} \cdot \vec{B})^2$. Thus, for a D3-brane we can write (6.1) as

$$\begin{aligned} \mathcal{L} &= -T_3 V(T) \sqrt{\beta} \sqrt{1 + \frac{\alpha T'^2}{\beta}} + \mathcal{L}_{\text{CS}} \\ &= -T_3 V(\tilde{T}) \sqrt{\beta} \sqrt{1 + \tilde{T}'^2} + \mathcal{L}_{\text{CS}}, \end{aligned} \quad (6.26)$$

where \mathcal{L}_{CS} is the Chern-Simons like Lagrangian. Following earlier steps, we can find a new related theory by making the deformation

$$\mathcal{L} \rightarrow \tilde{\mathcal{L}}, \quad V(\tilde{T})\sqrt{\beta} \rightarrow \tilde{V}(\tilde{T}) = \frac{\alpha V(T)}{\lambda\sqrt{\beta}}, \quad (6.27)$$

where the deformed theory is given by

$$\tilde{\mathcal{L}} = -T_3 \tilde{V}(\tilde{T})\sqrt{1 + \tilde{T}'^2} + \mathcal{L}_{\text{CS}}. \quad (6.28)$$

Note that the Chern-Simons like term that couples to the original D3-brane, is not affected by the deformation. However, this may not be true for other coupling terms, such as $dT \wedge F$, involving tachyon field, also used in the literature [13, 11]. The deformed tachyon field is defined as

$$\tilde{T} = \pm \int \sqrt{\frac{\alpha}{\beta}} dT, \quad (6.29)$$

and the deformed tachyon potential is $\tilde{V}(\tilde{T})$. Again, as previously considered, by performing the transformation (or deformation) (6.27) both theories (6.26) and (6.28) maintain their energy-momentum components T_{11} and \tilde{T}_{11} conserved. The pressures T_{11} and \tilde{T}_{11} are constants related to each other via real parameter λ as we can check explicitly:

$$\tilde{T}_{11} = -T_3 \frac{\tilde{V}(\tilde{T})}{\sqrt{1 + \tilde{T}'^2}} = -T_3 \frac{\alpha V(T)}{\lambda\sqrt{\beta}} \frac{1}{\sqrt{1 + \frac{\alpha T'^2}{\beta}}} = -T_3 \frac{\alpha V(T)}{\lambda\sqrt{\beta + \alpha T'^2}} = \frac{T_{11}}{\lambda}, \quad (6.30)$$

with T_{11} given in (6.22). The transformations (3.6)-(3.7) generalize and are now given by

$$d\tilde{T} = \pm \sqrt{\frac{\alpha}{\beta}} dT = \pm \frac{\lambda\alpha^{-1/2}dT}{[f'(\tilde{T})]^2}, \quad (6.31)$$

$$\tilde{V}(\tilde{T}) = \frac{\alpha V(T)}{\lambda\sqrt{\beta}} = \frac{V(T)}{[f'(\tilde{T})]^2}, \quad (6.32)$$

where $\tilde{T} = f^{-1}(T)$, and the deformation function $f(\tilde{T})$ is defined as

$$f(\tilde{T}) = \int \lambda^{1/2} \frac{\beta^{1/4}}{\alpha^{1/2}} d\tilde{T}. \quad (6.33)$$

Before going into calculations some comments are in order. Note that since the deformation function captures all the information related to the gauge fields, the previous calculations concerning deformed tachyon kinks continues valid here. The field components contributes to a same deformation. The novelty is that now we can deal with the effects of such components individually. The whole deformation is related to the field components as

$$[f'(\tilde{T})]^4 \equiv \lambda^2 \frac{\beta}{\alpha^2} = \lambda^2 \frac{1 + \vec{B}^2 - \vec{E}^2 - (\vec{E} \cdot \vec{B})^2}{(1 + B_1^2 - E_2^2 - E_3^2)^2}. \quad (6.34)$$

Note that for $\vec{B} = 0$, $E_1 = E$ and $E_2 = E_3 = 0$ this equation recovers the equation (5.14). The problem with all field components involve six variables $(E_1, E_2, E_3, B_1, B_2, B_3)$. So

in principle we need six equations to solve the problem. However, the equation (6.34) together with the independent conserved quantities we found before are sufficient to solve the system. Using the fact that B_1 , E_2 and E_3 are constants, a simple nontrivial and consistent example is given by setting $E_2 = 0$, $B_2 = 0$ and $B_1 = E_3$, that leaves us with the following relevant equations for $E_1(x)$ and $B_3(x)$

$$[f'(\tilde{T})]^4 = \lambda^2[1 + B_3^2 - E_1^2 - (E_1E_3 + E_3B_3)^2], \quad (6.35)$$

$$T_{13} = T_{11}(E_1E_3 + E_3B_3). \quad (6.36)$$

These equations can be solved for the dynamical field components

$$E_1(x) = \frac{1}{2} \frac{[f'(\tilde{T})]^2 T_{11}^2 E_3^2 + \lambda^2 (T_{11}^2 E_3^2 + T_{13}^2 - E_3^2 T_{13}^2)}{\lambda^2 T_{13} T_{11} E_3}, \quad (6.37)$$

$$B_3(x) = \frac{1}{2} \frac{[f'(\tilde{T})]^2 T_{11}^2 E_3^2 + \lambda^2 (-T_{11}^2 E_3^2 + T_{13}^2 + E_3^2 T_{13}^2)}{\lambda^2 T_{13} T_{11} E_3}, \quad (6.38)$$

where λ , T_{11} , T_{13} , and E_3 are non-vanishing constants. Thus, up to these constants, the problem is fully determined as long as we know the deformation function $f(\tilde{T})$. Let us now use the deformation function (5.9) and the deformed tachyon kink \tilde{T} given in (5.12), to find the field components — note that $[f'(\tilde{T})]^2 = 1/[1 - (\tilde{T}^2/T_0^2)]^2$. The magnitude of the field components, i.e., $|E_1(x)|$ and $|B_3(x)|$ can be peaked around $x = 0$ according to the sign of their concavity

$$d_x^2 |E_1(x)|_{x=0} = \text{sgn} [T_{11}^2 E_3^2 (1 - \lambda^2) + \lambda^2 T_{13}^2 (E_3^2 - 1)] \left(1 - \frac{T_{11}^2}{\lambda^2 T_3^2}\right) \zeta, \quad (6.39)$$

$$d_x^2 |B_3(x)|_{x=0} = \text{sgn} [T_{11}^2 E_3^2 (1 - \lambda^2) + \lambda^2 T_{13}^2 (E_3^2 + 1)] \left(1 - \frac{T_{11}^2}{\lambda^2 T_3^2}\right) \zeta, \quad (6.40)$$

where $\zeta = \frac{2|T_{11}||E_3|\omega^2}{\lambda^2 |T_{13}|} > 0$, and T_3 is the 3-brane tension, i.e., the 3d analog of T_1 given in section 5. Thus, replacing $T_1 \rightarrow T_3$ the constraint on λ given in (5.15) implies that $\frac{T_{11}^2}{\lambda^2 T_3^2} < 1$ (and $T_{11} > T_3$, $\lambda^2 > 1$), such that the signs of the concavities (6.39)-(6.40) are simply determined by the “sgn” prefactor. For example, we can easily see that for $E_3 = 1$ and $T_{13} = T_{11}$ the concavities are $d_x^2 |E_1(x)|_{x=0} < 0$ and $d_x^2 |B_3(x)|_{x=0} > 0$. Thus, for this choice of parameters the only field component that is localized around $x = 0$ is $E_1(x)$, whereas the component $B_3(x)$ is not. They are depicted in figure 4 and figure 5, respectively. In the example shown in figures 4 and 5, the electric field component $E_1(x)$ is confined to the deformed tachyon kink (a D2-brane), while the magnetic field component is not. $B_3(x)$ is “peaked” just outside the D2-brane. As another example, one can also check that for $E_3 = 1$ and $T_{13} = nT_{11}$, both components $B_3(x)$ and $E_1(x)$ get localized on the D2-brane, as long as $n < (1/\lambda)[(\lambda^2 - 1)/2]^{1/2}$. Finally, for $T_{13} = T_{11}$ and $E_3^2 \sim \lambda^2$ we find that both components $B_3(x)$ and $E_1(x)$ are expelled from the D2-brane. In summary, the fields \vec{E} and \vec{B} are not present everywhere on the deformed D3-brane, because the presence of the deformed D2-brane inside, always confines (or expels) electromagnetic components. This is only possible in the deformed theory provided its D2-brane has finite tension. This is assured if the deformation function used satisfies criterion (3.10). As the deformed

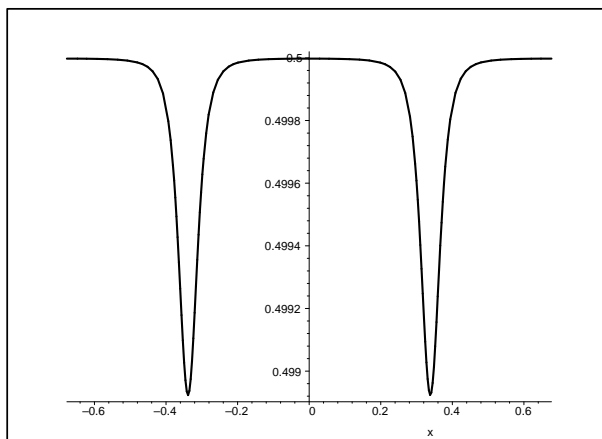


Figure 4: The behavior of the electric field component $E_1(x)$ around the deformed tachyon kink (at $x \approx 0$), in the interval $-\pi/\omega \leq x \leq \pi/\omega$, for $E_3=1$, $T_0=1$, $T_3=0.1$, $T_{13}=T_{11}=10$, and $\lambda=464$.

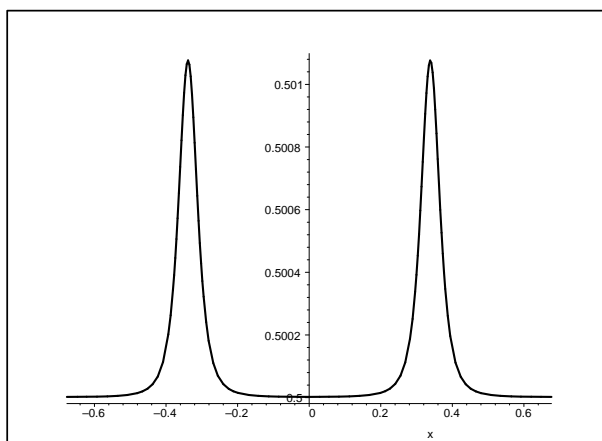


Figure 5: The behavior of the magnetic field component $B_3(x)$ around the deformed tachyon kink (at $x \approx 0$), in the interval $-\pi/\omega \leq x \leq \pi/\omega$, for $E_3=1$, $T_0=1$, $T_3=0.1$, $T_{13}=T_{11}=10$, and $\lambda=464$.

D3-brane decay it may leave behind stable D2-branes with electromagnetic components localized on them.

7. Conclusions

In this paper we addressed the issue of deforming tachyon potentials and tachyon kinks by using a deformation function given in terms of gauge field components. We have found that a singular tachyon kink can be resolved via such deformation. The resolved (non-singular) tachyon kinks can represent Dp -branes in string theory [1, 2, 11].

Concerning tachyon cosmology that makes use of time-dependent homogeneous tachyons $T(t)$ [23], some brief comments are in order. In realistic models of scalar cosmology, it is expected that the small fluctuations of the scalar field during the reheating are stable. This is because there is the possibility of the fast growth of inhomogeneous scalar field fluctuations after inflation. In ref. [24] it was pointed out that tachyon cosmology, for some string theory motivated tachyon potentials, indeed does develop such instability. The

instability of the inhomogeneous tachyon fluctuations $\delta T(t, \vec{x}) = \int d^3k T_k(t) e^{i\vec{k}\cdot\vec{x}}$ around $T(t)$, governed by the following equation

$$\frac{\ddot{T}_k}{1 - \dot{T}^2} + \frac{2\dot{T}\ddot{T}}{(1 - \dot{T}^2)^2}\dot{T}_k + [k^2 + (\log V)_{,TT}]T_k = 0, \tag{7.1}$$

becomes non-linear very quickly [24]. However, as has been shown in [25], this has to do with the ‘concavity’ $(\log V)_{,TT}$ of the tachyon potential, which means that for $(\log V)_{,TT}$ positive (negative) one finds stable (unstable) fluctuations. In ref. [25] it was considered a scalar potential which is essentially a ‘tachyon potential’ with reversed concavity. For example, a tachyon potential $V(T) = e^{-(T/T_0)^2}$ with negative concavity $-2/T_0^2$ is replaced to another $\tilde{V}(\tilde{T}) = e^{(\tilde{T}/\tilde{T}_0)^2}$ with positive concavity $2/\tilde{T}_0^2$, i.e., a massive inflaton potential. Given that we have investigated deformation of tachyon potentials in the previous sections, we can make use of this tool in order to replace a cosmological scenario to another. Thus, a tachyon potential motivated by string theory could be deformed into another by a deformation function that is given in terms of gauge field components associated with background R-R fields.

In fact, in this paper, we have investigated an example in which the concavity of potentials indeed changes through deformation. This is the case $p = -3$, where the potential $V(T) = 1/\text{sech}^3(T/T_0)$ with positive concavity $(3/T_0^2)\text{sech}^2(T/T_0)$ is deformed into the potential $\tilde{V}(\tilde{T}) = \left(1 - \frac{\tilde{T}^2}{T_0^2}\right)^{\frac{1}{2}}$ with negative concavity $-(\tilde{T}^2 + T_0^2)/(-T_0^2 + \tilde{T}^2)^2$.

In a certain sense the deformed theory is good for tachyon kinks (inhomogeneous tachyons), but not good for tachyon cosmology (homogeneous tachyon). The smooth tachyon kink is generated by the deformed tachyon potential \tilde{V} which has negative concavity, whereas the tachyon cosmology is stable for the “non-deformed” theory with tachyon potential V which has positive concavity.

In summary, in this paper we have found an example where the deformation of a tachyon theory via gauge field components connects two different limits. On one hand, in the deformed theory we can find smooth tachyon kinks representing lower dimensional D-branes, but without stable cosmology, while on the other hand we have stable cosmology without smooth lower dimensional D-branes in the original (non-deformed) theory. In particular, note that the same deformation that makes the smooth tachyon kinks (or deformed D2-branes) to localize gauge field components, in some way also prevents stable cosmology on the deformed D3-brane. This seems to be consistent with the fact that for potentials with negative concavity (the deformed potential), the instability of the tachyon cosmology after inflation takes place and inhomogeneous tachyon fluctuations grow very fast. In such scenario the instability may start populating the universe with extended objects such as deformed tachyon kinks or D2-branes, which are beyond the perturbative spectrum. This issue and its higher dimensional extensions are to be further explored elsewhere.

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