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Stretched strings in tachyon condensation models

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Abstract: In this note, we consider the two derivative truncation of boundary string field theory for the unstable D9-brane in type-IIA string theory. We construct multiples of the stable codimension 1 solitons that correspond to stacks of D8-branes. We find the fluctuation modes that correspond to open strings stretching between the branes, and find that their masses are consistent with the string tension. We show that these modes are localized halfway between the branes and that their width is independent of the brane separation.

Keywords: Brane Dynamics in Gauge Theories, Tachyon Condensation.

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Simple field theory models have proven to be quite useful in giving qualitative and surprisingly good quantitative descriptions for tachyon condensation[1] in open string field theory [2]–[10]. These models turn out to be two derivative truncations of boundary string field theory [11]–[16]. In particular, the model described in [6] is the two derivative truncation of boundary superstring field theory [14]. The lagrangian for this model is given by

$$-\mathcal{T}_9 \left(\frac{1}{2} \partial_\mu T \partial_\mu T + 1 \right) e^{-T^2/4}, \tag{1}$$

where $\mathcal{T}_9$ is the tension of the the unstable D9-brane. The potential has an unstable vacuum at $T = 0$, where the tachyon mass squared is $-1/2$ (we have set $\alpha' = 1$). It also has stable vacuua at $T = \pm\infty$, where the mass of the tachyon becomes infinite. There is a stable codimension one soliton solution with a tension $\mathcal{T}_8 = 2\sqrt{2\pi}\mathcal{T}_9$, a result reasonably close to the actual value of $\mathcal{T}_8 = \sqrt{2}\pi\mathcal{T}_9$. It was also shown in [6] that the fluctuation modes about this soliton have integer mass squared level spacing. Thus, this model nicely conforms with the Sen conjectures [17].

Of course, open string field theory should contain an infinite number of fields. In [7] a prescription was given for including higher level fields into the models. For instance, it was argued that including gauge fields with the Lagrangian

$$-\frac{1}{4} F_{\mu\nu} F^{\mu\nu} e^{-T^2/4} \tag{2}$$

led to fluctuation modes on the codimension one-brane with the correct mass squared spectra. Moreover, it was shown that a zero momentum mode for the gauge field was gauged away on the D8-brane but was compensated by the zero mode coming from the tachyon field. Some preliminary evidence was also given for how to include higher level fields in the model. In any event, the surprising conclusion is reached that dropping the higher derivative terms does not seem to drastically effect the physics, at least at tree-level. Dropping the higher derivatives changes the tensions of the branes and fattens them out, but it does not seem to effect the spectrum.

The "open string" modes that we find in these models are localized around the brane, that is, the center of the soliton. One interesting question is what do the stretched strings look like in these models. By stretched string we mean a mode that corresponds to a string stretching between two D-branes separated by some distance. The naive guess is that the mode is more or less evenly distributed between the branes. This turns out not to be the case. Instead we will show that the mode *is* localized halfway between the branes, with the localization width independent of the brane separation. This last fact seems to indicate that the modes will be localized, not stretched, even when higher derivative terms are taken into account.

Nevertheless, the modes get an extra contribution to their mass-squared which is precisely the amount expected for a string stretched between two D-branes.

In order to describe multiple brane solutions, the tachyon field needs to be generalized to an N by N tensor field [14] and [18]–[22]. The tachyon potential is then [14]

$$V(T) = \mathrm{Tr}(e^{-T^2/4}). \tag{3}$$

[1] For early pioneering work in the field of tachyon condensation see [1].

Generalizing the kinetic term requires some care because of ordering ambiguities. One consideration is to find a solvable spectrum. We will see that

$$\frac{1}{2}\,\mathrm{Tr}(\partial_\mu T e^{-T^2/8}\partial_\mu T e^{-T^2/8}) \tag{4}$$

serves this purpose.

Non-trivial codimension 1 solutions to the equations of motion are the diagonal matrices

$$T = \mathrm{diag}(T_1, T_2, \ldots, T_N)\,, \tag{5}$$

where

$$T_i = \pm\sqrt{2}(x - x_i) \tag{6}$$

and where x_i is the location of the i^{th} brane. The $+$ $(-)$ sign corresponds to a D8-brane (anti-brane).

Let us begin by assuming that $N = 2$, our results being easily generalized to higher N. In this case, we can rewrite T as

$$T = T_0\sigma_0 + \Delta T\sigma_3\,, \tag{7}$$

where

$$\sigma_0 = \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix}, \qquad \sigma_3 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \tag{8}$$

We now look for fluctuations about these solutions. The diagonal fluctuations are exactly as in [6, 7]. However, the off-diagonal fluctuations are non trivial if $\Delta T \neq 0$. In order to find these, let us write the off-diagonal fluctuation piece as

$$T_{fl} = \begin{pmatrix} 0 & T_+ \\ T_- & 0 \end{pmatrix}. \tag{9}$$

We now make use of the fact that

$$\mathrm{Tr}\left(\begin{pmatrix} 0 & A \\ B & 0 \end{pmatrix}(\sigma_3)^n \begin{pmatrix} 0 & A \\ B & 0 \end{pmatrix}(\sigma_3)^m\right) = AB\,[(-1)^n + (-1)^m]\,, \tag{10}$$

which among other things implies that

$$\begin{aligned} &\mathrm{Tr}\left[\begin{pmatrix} 0 & A \\ B & 0 \end{pmatrix} e^{-T_0^2/8-\Delta T^2/8-T_0\Delta T\sigma_3/4} \begin{pmatrix} 0 & A \\ B & 0 \end{pmatrix} e^{-T_0^2/8-\Delta T^2/8-T_0\Delta T\sigma_3/4}\right] = \\ &\quad = 2AB\exp\left(-\frac{T_0^2}{4} - \frac{\Delta T^2}{4}\right). \end{aligned} \tag{11}$$

To quadratic order in the off-diagonal fluctuations, the lagrangian is

$$-\left(\left(\partial_\mu T_+ \partial_\mu T_- + \frac{1}{4}\left[(\partial_\mu T_0)^2 - (\partial_\mu \Delta T)^2 - 2\right] T_+ T_-\right) e^{-T_0^2/4-\Delta T^2/4}\,, \tag{12}$$

where we have used the equations of motion for T_0 and ΔT,

$$\begin{aligned} \partial^2 T_0 - \frac{1}{4}T_0\left[(\partial T_0)^2 + (\partial\Delta T)^2\right] - \frac{1}{2}\Delta T\partial_\mu T_0\partial_\mu\Delta T + \frac{1}{2}T_0 &= 0\,, \\ \partial^2\Delta T - \frac{1}{4}\Delta T\left[(\partial T_0)^2 + (\partial\Delta T)^2\right] - \frac{1}{2}T_0\partial_\mu T_0\partial_\mu\Delta T + \frac{1}{2}\Delta T &= 0\,. \end{aligned} \tag{13}$$

Making further use of the equations of motion and integrating by parts, we find that the expression in (12) can be rewritten as

$$-\partial_\mu\left(\frac{T_+}{\Delta T}\right)\partial_\mu\left(\frac{T_-}{\Delta T}\right)e^{-T_0^2/4-\Delta T^2/4}(\Delta T)^2\,. \tag{14}$$

The U(N) symmetry of the tachyon field is actually a gauge symmetry, so accordingly, the ordinary derivatives should be replaced by covariant derivatives. We will assume that the kinetic term has the form

$$-\frac{1}{2}\,\mathrm{Tr}([\partial_\mu - iA_\mu, T]e^{-T^2/8}[\partial_\mu - iA_\mu, T]e^{-T^2/8})\,. \tag{15}$$

We then expect *all* of the off-diagonal fluctuations of T to be eaten by the gauge fields. Again, let us consider the case that $N = 2$ and let us find the terms quadratic in A_μ and the tachyon fluctuations. Using the properties in (10) and (11), we find the quadratic off-diagonal piece to be

$$-4\left(A_{+\mu}+\frac{i}{2}\partial_\mu\left(\frac{T_+}{\Delta T}\right)\right)\left(A_{-\mu}-\frac{i}{2}\partial_\mu\left(\frac{T_-}{\Delta T}\right)\right)e^{-T_0^2/4-\Delta T^2/4}(\Delta T)^2\,. \tag{16}$$

To show the importance of ordering in the tachyon kinetic terms, let us suppose that the appropriate term in the lagrangian were

$$\frac{1}{2}\,\mathrm{Tr}\left([\partial_\mu - iA_\mu, T]^2 e^{-T^2/4}]\right)\,. \tag{17}$$

In this case, one would find that the off-diagonal fluctuations have the form

$$4\left(A_{+\mu}+\frac{i}{2}\partial_\mu\left(\frac{T_+}{\Delta T}\right)\right)\left(A_{-\mu}-\frac{i}{2}\partial_\mu\left(\frac{T_-}{\Delta T}\right)\right)e^{-T_0^2/4-\Delta T^2/4}\cosh\left(\frac{T_0\Delta T}{4}\right)\Delta T^2\,, \tag{18}$$

and so would clearly have a different spectrum than that coming from (16).

Now consider the kinetic term for the gauge fields. Again, there will be ordering ambiguities. There is also some ambiguity as to how even the diagonal action should be chosen. For the D9-brane, one expects the quadratic piece of the lagrangian to come from expanding the Born-Infeld lagrangian

$$\sqrt{\det(\eta_{\mu\nu}-2\pi F_{\mu\nu})}\left(\frac{1}{2}G^{\mu\nu}\partial_\mu T e^{-T^2/8}\partial_\nu T e^{-T^2/8}+e^{-T^2/4}\right)\,, \tag{19}$$

where $G^{\mu\nu}$ is the open string metric which is the symmetric part of [23]

$$(\eta_{\mu\nu}-2\pi F_{\mu\nu})^{-1}\,. \tag{20}$$

The problem with this formulation is that the D8 solution for T will have no quadratic piece for the gauge fields; the contribution from the kinetic term cancels off the contribution from the potential. These difficulties are probably due to our two derivative limitation as well as the assumption in (19) that the field strengths are constant.

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We will instead assume that the quadratic gauge lagrangian has the form

$$-\frac{(2\pi)^2}{4}\,\mathrm{Tr}\left(\frac{1}{2}F_{\mu\nu}\partial_\gamma T e^{-T^2/8}\partial_\gamma T F_{\mu\nu}e^{-T^2/8}+F_{\mu\nu}e^{-T^2/8}F_{\mu\nu}e^{-T^2/8}\right). \tag{21}$$

This form of the lagrangian is consistent with reduction from the unstable D9 to the stable D8-branes. The ordering in (21) leads to a particularly simple form for the off-diagonal gauge fluctuations,

$$\begin{aligned}(2\pi)^2\left(\partial_\mu A_{+\nu}\partial_\nu A_{-\nu}-\partial_\mu A_{+\nu}\partial_\nu A_{-\mu}\right)e^{-T_0^2/4-\Delta T^2/4}\times\\ \times\left(\frac{1}{2}\partial_\mu T_0\partial_\mu T_0+\frac{1}{2}\partial_\mu\Delta T\partial_\mu\Delta T+1\right).\end{aligned} \tag{22}$$

Defining $A'_{+\mu}$ as

$$A'_{+\mu}=A_{+\mu}+\frac{i}{2}\partial_\mu\left(\frac{T_+}{\Delta T}\right), \tag{23}$$

the complete quadratic lagrangian for the off-diagonal gauge fluctuations is

$$\begin{aligned}-\Bigg[(2\pi)^2\left(\partial_\mu A'_{+\nu}\partial_\mu A'_{-\nu}-\partial_\mu A'_{+\nu}\partial_\nu A'_{-\mu}\right)\left(\frac{1}{2}(\partial T_0)^2+\frac{1}{2}(\partial\Delta T)^2+1\right)+\\ +4A'_{+\mu}A'_{-\mu}\Delta T^2\Bigg]e^{-T_0^2/4-\Delta T^2/4}.\end{aligned} \tag{24}$$

We can now use (24) to compute the spectrum of $A'_{\pm\mu}$. Let us take the classical solutions in (6), and assume that we have two D8-branes, one at $x=x_0/2$ and the other at $x=-x_0/2$. Thus,

$$\begin{aligned}T_1&=\sqrt{2}(x-x_0/2)\,,\\ T_2&=\sqrt{2}(x+x_0/2)\,,\end{aligned} \tag{25}$$

and so

$$\begin{aligned}T_0&=\sqrt{2}x\,,\\ \Delta T&=-\frac{\sqrt{2}}{2}x_0\,.\end{aligned} \tag{26}$$

The equations of motion for T_+ lead to the equation

$$\partial_\mu(A'_{+\mu}e^{-x^2/2})=0\,. \tag{27}$$

Substituting this into the equation of motion for $A'_{+\nu}$, we have

$$-\partial^2A'_{+\nu}+x\partial_xA'_{+\nu}+A'_{+,x}\delta_{\nu x}+\frac{x_0^2}{(2\pi)^2}A'_{+\nu}=0\,. \tag{28}$$

Letting $A_{+\nu}=B_\nu e^{-x^2/4}$ we find the equation

$$-\partial^2B_\nu+\left[\frac{1}{4}x^2-\frac{1}{2}+\delta_{\nu x}+\frac{x_0^2}{(2\pi)^2}\right]B_\nu=0\,. \tag{29}$$

Hence the mass spectra for the modes is

$$m^2 = n + \frac{x_0^2}{(2\pi)^2}\,, \qquad n \geq 0 \tag{30}$$

for polarizations along the brane ($\nu \neq x$) and

$$m^2 = n + 1 + \frac{x_0^2}{(2\pi)^2}\,, \qquad n \geq 0 \tag{31}$$

for the transverse polarization ($\nu = x$). Since the string tension is $1/(2\pi)$, we obtain the desired contribution to the mass coming from a string stretching between the two-branes separated by a distance x_0. Note that the massive modes of T_+ and T_- have been eaten by the B_x modes while the massless modes of T_+ and T_- have been Higgsed by the massless gauge mode.

Let us examine a few facts that might seem surprising about these stretched modes. First of all, they are not really stretched. If one were to consider a classical wave corresponding to one of these modes, they would find the energy density localized at $x = 0$. This is half-way between the branes, which contrasts with the diagonal modes which are localized on the branes. One also finds that the width of the modes is independent of the brane separation.

However, there is no suppression of the interactions between the charged and the diagonal modes, even at large D-brane separation. For instance, consider the full non-abelian lagrangian coming from (21). Writing the gauge field as

$$A_\mu = \begin{pmatrix} A_{3\mu} & A_{+\mu} \\ A_{-\mu} & -A_{3\mu} \end{pmatrix}, \tag{32}$$

the cubic interaction term coming from (21) is

$$2i(2\pi)^2 A_{3\nu}\left[A_{+\mu}\partial_{[\nu}A_{-\mu]} - A_{-\mu}\partial_{[\nu}A_{+\mu]}\right] e^{-T_0^2/4-\Delta T^2/4}. \tag{33}$$

The kinetic term for $A_{3\mu}$ comes with a gaussian factor localizing the mode around the branes. But the normalized modes for $A_{3\mu}$ have no such suppression. For instance, the zero mode has $A_{3\mu}$ constant in the x direction. Hence the 3 point interaction is not suppressed even though the charged modes and the neutral modes are localized at different positions in x.

It is straightforward to generalize these results to multiple branes. In this case, the off-diagonal fluctuations corresponding to strings attached to branes i and j are

$$\left[2(2\pi)^2\left(\partial_\mu A'_{ij,\nu}\partial_\nu A'_{ji,\nu} - \partial_\mu A'_{ij,\nu}\partial_\nu A'_{ji,\mu}\right) + A'_{ij,\nu}A'_{-\nu}(T_i - T_j)^2\right] e^{-T_i^2/8-T_j^2/8}\,. \tag{34}$$

This will lead to the anticipated spectrum for the stretched strings.

One can also consider the case where the transverse direction is compactified on a circle such that

$$x \equiv x + 2\pi R\,. \tag{35}$$

The analysis will be similar to that in [24] for matrix models on the circle. Now the gauge group must be extended to U(∞). A stable D8-brane on the circle is given by the infinite diagonal matrix

$$T = \mathrm{diag}(\ldots, T_n, T_{n+1}, \ldots)\,, \qquad T_n = \sqrt{2}(x - 2\pi n R)\,. \tag{36}$$

Hence, the solution is invariant under the shift $x \to x + 2\pi R$ up to a gauge transformation. An open string wound n times around the circle is then given by a superposition of the $A_{i,i+n}$ modes.

Note that the center of mass U(1) is not present in (34). The usual picture for an open string stretched between two-branes is that of a flux tube between opposite charges of the diagonal U(1). That is not the picture that emerges here. Instead, everything operates as an adjoint Higgsing of the SU(N) group. Of course, this is expected for the low energy effective gauge theory on the world volume, but it is nonetheless surprising to see it happen for the stringy modes as well.

This also raises questions about how the closed strings will appear [25]–[32]. One could certainly look for closed strings in a D-brane vacuum. For example, in the compact case, there should be closed strings that wrap around the compact direction. It seems from the analysis here that these are more likely to come from the SU(∞) sector and not the COM U(1), since the decay of a wound open string into an unwound open string and a closed string with zero momentum in the transverse direction still needs to preserve the winding. The wound open string modes are localized, either on the brane or diametrically opposite from it. One would expect a closed string to be delocalized on the circle. It would be interesting to see how this could occur, if indeed it can occur when only considering a finite number of the open string fields. One might need the full machinery of cubic string field theory [33] or vacuum string field theory [34] to see this.

It is also possible that the localization is gauge dependent. For instance, in boundary string field theory, the width of a D-brane has zero size [12], but in cubic string field theory, the width is non zero [33]. These theories must be related by a gauge transformation, so the width of the D-brane must be a gauge artifact. Likewise, the widths of fluctuation modes might also be gauge artifacts and that a gauge exists where the mode is evenly distributed between the branes.

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