Description of Intersecting Branes via Tachyon Condensation

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Abstract

We construct a model describing BPS brane-systems using low energy effective theory of brane-antibrane system. Both parallel branes and intersecting branes can be treated by this model. After tachyon condensation, the dynamics of fluctuations around such brane-systems is supersymmetric if the degrees of freedom are restricted on the branes. The form of the tachyon potential and the application of this model to the black hole physics are discussed.

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1. Introduction

Low energy dynamics of the combined D-brane systems is important, for example, for black hole physics [1]. It is easy to describe such system which contains only one kind of D-branes. The effective theory of the system is the dimensional reduction of the ten-dimensional super Yang-Mills theory [2]. However, construction of effective theory is not so straightforward when the system consists of a combination of D-branes with different dimensions. In most cases, one pays attention to a specific part of the worldvolume and consider the effective theory living only on that region [1][3]. For example, the effective theory of the D1-D5 system is a two-dimensional superconformal field theory living on the D1-branes. Although this is enough to discuss the physics of the black holes, it is preferable to have another way of describing the system as a whole.

The difficulty in such description may be how to treat the brane-intersections. When one writes down the effective theory of the massless modes on the whole brane-system, the theory contains a delta-function-like interaction at the intersection. It would not be an easy task to deal with such theory, although this description has some applications [4].

In this paper we propose another way of describing the D-brane system which includes intersecting branes. Our idea is to apply the tachyon condensation [7]. In this situation, pair annihilations of D-branes can be discussed and in some cases the lower dimensional branes can remain as a result. The worldvolume of the resulting brane is determined by the form of the tachyon field and, in particular, it need not be a smooth manifold. Therefore the theory describing the above phenomenon would be able to handle the low energy dynamics of the intersecting branes.

This paper is organized as follows. We review the story of the tachyon condensation briefly in section 2. The construction of the model is discussed in section 3. This model can describe both parallel branes (section 4) and intersecting branes (section 5). Section 6 is devoted to discussions. In Appendix, we summarize the properties of the vortex solutions used in this paper.

2. Review of the tachyon condensation

In this section, we will review some properties of Dp-˘Dp system. While both branes preserve half of the supersymmetries, this system breaks all of the supersymmetries when they exist together.

There exists a tachyonic mode coming from the string stretched between Dp-brane and ˘Dp-brane, which survives the GSO projection [5]. This indicates the instability of the system. Assuming that the tachyon potential has a minimum, the tachyon would condense and this system decays to some stable vacuum.

Let $T$ denote the tachyon field whose potential $V(T)$ has a minimum at $T = T_0$. When $T = T_0$ everywhere, there remains nothing: pair annihilation of the D-branes occurs. This has been checked by examining that $V(T_0)$ would exactly cancel the tension of the D-branes, using string field theory [6]. On the other hand, when $T$ has some nontrivial
form, there remains a lower dimensional brane which is stable but not necessarily BPS \[7\]. In this case, the tension of the originally unstable brane-antibrane system is cancelled in almost all regions except at those places where the tachyon field behaves nontrivially. The energy therefore remains nonzero only on that regions, and this is regarded as a lower dimensional brane. It is argued that K-theory determines what kind of branes can appear after tachyon condensation \[8\][9].

For example, consider the D1-\(\overline{D}1\) system extended along \(x^1\) direction \[7\]. Suppose that \(T\) supports a kink,

\[
T(x) = \begin{cases} 
+T_0 & (x^1 \to +\infty) \\
-T_0 & (x^1 \to -\infty) 
\end{cases}

T(x^1 = 0) = 0,
\]

where we have assumed that \(V(-T) = V(T)\) and \(T = -T_0\) is also a minimum of \(V(T)\). Then the remaining energy is localized around \(x^1 = 0\), which can be regarded as D0-brane. In ordinary Type IIB theory this is of course unstable, but this is stable in Type I theory \[7\].

The brane annihilation summarized above is also described in terms of gauge theory on the D-brane worldvolume \[9\]. The worldvolume theory would be a gauge theory coupled to the tachyon. When the tachyon acquires VEV, the gauge symmetry is broken down by the Higgs mechanism. Thus the gauge fields acquire masses and are decoupled from the low energy effective theory. This indicates the disappearance of the D-branes. The case of nontrivial \(T\) can be described similarly. If \(T = 0\) in some region, the gauge symmetry is recovered there, and this corresponds to the resulting lower dimensional brane.

In the gauge theory language, the classification of the remaining branes can be done in terms of homotopy groups of the vacuum manifold \[9\].

3. The model

We will consider in this section the model of D9-\(\overline{D}9\) system in Type IIB theory. The model of Dp-\(\overline{D}p\) system can be obtained by dimensional reduction.

There is a \(U(1)\) gauge multiplet on each brane and also a tachyon field \(T\) which couples to both gauge fields. One linear combination of these gauge fields (so called diagonal \(U(1)\)) will decouple from the dynamics (it is argued that this gauge field becomes a fundamental string \[10\], and its generalizations are discussed \[11\]). Then the resulting model is ten-dimensional \(U(1)\) gauge theory coupled to the tachyon field,

\[
S = \int d^{10}x \left\{-\frac{1}{4} F_{\mu\nu}F^{\mu\nu} + \frac{i}{2} \bar{\psi} \Gamma^\mu \partial_\mu \psi - |D_\mu T|^2 - V(T) \right\}
\]

where \(D_\mu T = \partial_\mu T + iA_\mu T\). \(\Gamma^\mu\) are the ten-dimensional gamma matrices and \(\psi\) is a ten-dimensional Majorana-Weyl spinor. Although there exists a charged massless fermion coming from the string stretched between branes, it is not important for the following discussions and we will ignore it. This action would be a first few terms in the expansion.
of the more complicated DBI-like action. The spacetime action of non-BPS brane is
argued in [12]. It is determined by requiring that it reduces to the ordinary DBI action if
the tachyon and half of the massless fermions which are absent on the BPS brane are set
to zero. This seems to imply that in the Dp-¯Dp case the action is also supersymmetric if
the tachyon field T is set to zero. The action of this system is discussed in [13].

The action (1) is obviously not supersymmetric. According to the discussion in string
theory, however, classical solutions should exist which preserve some of the supersymme-
tries. The first two terms are invariant under the ordinary transformations.

\[ \delta A_\mu = \frac{i}{2} \epsilon \Gamma_\mu \psi \]
\[ \delta \psi = - \frac{1}{4} F_{\mu \nu} \Gamma^{\mu \nu} \epsilon \] (2)

The tachyon should be invariant: \( \delta T = 0 \). Then,

\[ \delta S = \int d^{10}x \frac{i}{2} \bar{\psi} \Gamma_\mu (i T^\dagger \Gamma_\mu T - i D^\mu T^\dagger T). \] (4)

This vanishes when \( D_\mu T = 0 \) which leads to trivial solutions.

Now we will modify the transformations of the fermions (3) as follows,

\[ \delta \psi = - \frac{1}{4} F_{\mu \nu} \Gamma^{\mu \nu} \epsilon + f(|T|^2) \epsilon, \] (5)

where \( f(x) \) is a function which may have spinor indices. Then the variation of the action is

\[ \delta S = \int d^{10}x \frac{i}{2} \bar{\psi} \Gamma_\mu (-i T^\dagger \Gamma_\mu T + i D^\mu T^\dagger T + 2 \partial_\mu f(|T|^2)) \epsilon. \] (6)

We will show below that for a suitable choice of \( f(x) \), there exist classical solutions which
preserve some of supersymmetries.

4. Parallel D-branes

We will consider for example D4-¯D4 system which decays to D2-branes. Suppose that
D4-brane and D4-brane extend along the (01234) directions, and the resulting D2-branes
extend along (034) directions. The action of the model is eq.(1) dimensionally reduced to
five dimensions. For the solutions corresponding to such branes to exist, the appropriate
choice of \( f(x) \) is

\[ f(|T|^2) = - \frac{1}{2} s(|T|^2 - \zeta) \Gamma^1 \Gamma^2 \quad (\zeta > 0). \] (7)

The variation of the action is

\[ \delta S = - \int d^5x \left[ \frac{i}{2} \bar{\psi} T^\dagger \left( \Gamma^1 (i D_1 T + s D_2 T) + \Gamma^2 (i D_2 T - s D_1 T) \\
+ D_3 T (i \Gamma^i + \Gamma^{12}) + i A_k T (i \Gamma^k + \Gamma^{k12}) \right) \epsilon + (h.c.) \right], \] (8)
where \( i = 0, 3, 4 \) and \( k = 5, 6, \ldots, 9 \). In this case \( A_k \) is the neutral scalars. Thus the conditions for \( \delta S = 0 \) are

\[
\begin{align*}
D_i T - i s D_2 T &= 0 \\
D_i T &= 0 \\
A_k T &= 0,
\end{align*}
\]

We assume translational invariance along \( i \)-th direction, thus eq.(10) sets \( A_i = 0 \), and also \( A_k = 0 \) by eq.(11). \( \delta \psi = 0 \) means

\[
F_{12} + s(|T|^2 - \zeta) = 0.
\]

We will investigate the stability of this BPS solutions. Their energy is rewritten as follows.

\[
E = \int d^4 x \left\{ \frac{1}{2} \left( F_{12} + s(|T|^2 - \zeta) \right)^2 + |D_1 T - i s D_2 T|^2 + s \zeta F_{12} + V(T) - \frac{1}{2}(|T|^2 - \zeta)^2 \right\}
\]

If the tachyon potential takes the form

\[
V(T) = \frac{1}{2}(|T|^2 - \zeta)^2,
\]

then eq.(13) becomes

\[
E = \int d^4 x \left\{ \frac{1}{2} \left( F_{12} + s(|T|^2 - \zeta) \right)^2 + |D_1 T - i s D_2 T|^2 + s \zeta F_{12} \right\}.
\]

Therefore the BPS solutions are stable topologically.

Now we have the following BPS equations whose solutions preserve 16 supercharges.

\[
\begin{align*}
D_i T - i s D_2 T &= 0 \\
F_{12} + s(|T|^2 - \zeta) &= 0
\end{align*}
\]

This is the equations for the Nielsen-Olesen vortex [14]. The properties of this solutions are well-known. Some of these are collected in the Appendix. The solutions are labeled by an integer \( n \) (quantized magnetic flux), and is determined by specifying \( n \) distinct points in 1-2 plane, at which \( T = 0 \). As explained in section 2, zero loci of the tachyon field \( T \) correspond to the D-brane worldvolume. Thus the \( n \)-vortex solution describes \( n \) parallel D2-branes.

The worldvolume theory on the D2-branes should be a supersymmetric one. For this, eqs.(1)(10)(11) has to be satisfied even when the fluctuations are included. Let \( a_\mu, \varphi \) denote the fluctuations of \( A_\mu, \psi \) around the vortex solution respectively, then eqs.(1)(10)(11) means

\[
a_\mu \neq 0 \quad \text{only at} \quad T = 0
\]

From the transformations (2), \( \varphi \) is also restricted to the region specified by \( T = 0 \). Then the resulting model is supersymmetric and the physical degrees of freedom exist only at the cores of the vortices, as is expected from the D-brane interpretation.
It is interesting that we can construct \( n \) D-brane system from \( U(1) \) gauge theory. Moreover system with different number of D-branes merely corresponds to taking another classical solution in the same model. In the ordinary approach to the D-brane worldvolume theory, the number of D-branes must be fixed to construct the theory. It might be possible to relate this feature of our model to the second quantization of D-branes.

5. Intersecting branes

We will show in this section that our model can also describe intersecting branes in a similar way which was discussed in the previous section. We will consider the D4-D4 system which extends along \((01234)\) directions. The action of the model is again eq.\( (1) \) dimensionally reduced to five dimensions, and the tachyon potential and supersymmetry transformations are given as eqs.\( (14) \( \mathcal{F}(7) \). Suppose that this system decays to the intersecting D2-D2' system, in which D2-branes extend along \((034)\) directions and D2'-branes along \((012)\) directions. The conditions for the remaining supersymmetries are

\[
\Gamma^1 \Gamma^2 \epsilon = s' \Gamma^3 \Gamma^4 \epsilon \\
(s' = \pm 1).
\]

The BPS conditions are then

\[
\begin{align*}
D_1 T - isD_2 T &= 0 \\
D_3 T - i ss' D_4 T &= 0 \\
F_{12} + s' F_{34} + s(|T|^2 - \zeta) &= 0 \\
F_{13} - s' F_{24} &= 0 \\
F_{14} + s' F_{23} &= 0,
\end{align*}
\]

where we have assumed \( A_k = 0 \) \((k = 0, 5, 6, \ldots, 9)\) and that the solutions are static. It is easy to check that the solutions of these equations are stable topologically.

At first sight, these conditions are overdetermined. However the last two equations \( (23)(24) \) are in fact redundant. From eqs.\( (20)(21) \), \( T \) and \( A_i \) \((i = 1, 2, 3, 4)\) are written as follows.

\[
\begin{align*}
T &= \rho^{\frac{1}{2}} e^{i \omega} \\
A_1 &= \frac{1}{2} s \partial_2 \log \rho - \partial_1 \omega \\
A_2 &= -\frac{1}{2} s \partial_1 \log \rho - \partial_2 \omega \\
A_3 &= \frac{1}{2} ss' \partial_4 \log \rho - \partial_3 \omega \\
A_4 &= -\frac{1}{2} ss' \partial_3 \log \rho - \partial_4 \omega,
\end{align*}
\]
These solve eqs. (23)–(24) automatically. The only nontrivial equation (22) is then

$$-\frac{1}{2} \sum_{i=1}^{4} (\partial_i)^2 \log \rho + \rho - \zeta = 0. \quad (30)$$

$\omega$ is determined by requiring the regularity of $A_i$ and the single-valuedness of $T$.

The solutions we would like to find are symmetric under the rotations in 1-2 and 3-4 planes (D2(D2')-branes coincide). Now we take the polar coordinates for these planes,

$$(x_1, x_2) \rightarrow (r_1, \theta_1), \quad (x_3, x_4) \rightarrow (r_2, \theta_2),$$

and assume that $\rho$ is independent of both $\theta_1$ and $\theta_2$. Then (30) becomes

$$\left[ \frac{\partial^2}{\partial r_1^2} + \frac{1}{r_1} \frac{\partial}{\partial r_1} + \frac{1}{r_2} \frac{\partial}{\partial r_2} \right] \log \rho = 2(\rho - \zeta). \quad (31)$$

The region around $r_1 = r_2 = 0$, $\rho$ behaves as

$$\rho \sim r_1^{2n_1} r_2^{2n_2}. \quad (32)$$

For the regularity and the single-valuedness of the solution, $n_1$ and $n_2$ are positive integers.

It is convenient to make the following change of variables: $r_1 = e^{t_1}, r_2 = e^{t_2}$. Then

$$\left[ e^{-2t_1} \frac{\partial^2}{\partial t_1^2} + e^{-2t_2} \frac{\partial^2}{\partial t_2^2} \right] \log \rho = 2(\rho - \zeta). \quad (33)$$

When $t_2 \to \infty$, the second term of LHS is negligible and $\rho$ becomes $t_2$ (thus $r_2$) independent. This means that in this region $\rho$ approaches the vortex solution with magnetic flux $n_1$ discussed in the previous section. The same is true with the behavior of $\rho$ when $t_1 \to \infty$ ($r_1 \to \infty$) and its flux is $n_2$. By the continuity of $\omega$, $n_1$ and $n_2$ take the same values as the ones in eq. (32).

The global behavior of the solutions can also be discussed. Now take $u = \log \frac{\rho}{\zeta}$ and define $D_+$ to be the region where $u > 0$. $D_+$ is a finite region in $\mathbb{R}^2$. Then using eq. (33),

$$0 = \int_{D_+} dt_1 dt_2 \left[ u (-e^{2t_2} \partial_{t_2}^2 u - e^{2t_1} \partial_{t_1}^2 + 2\zeta e^{2t_1+2t_2}(e^u - 1)) \right]$$

$$= \int_{D_+} dt_1 dt_2 \left[ e^{2t_2} (\partial_1 u)^2 + e^{2t_1} (\partial_2 u)^2 + 2\zeta e^{2t_1+2t_2} u (e^u - 1) \right], \quad (34)$$

where we have used that $u = 0$ at the boundary of $D_+$. The integrand is strictly positive, and therefore, eq. (34) means $u \leq 0$ (i.e. $\rho \leq \zeta$) everywhere.

The above analyses imply that $\rho = 0$ at $r_1 = 0$ and/or $r_2 = 0$, and $\rho \sim \zeta$ away from the cores of the vortices. $\omega$ is determined to be $n_1\theta_1 + n_2\theta_2$, and this gives

$$\frac{1}{2\pi} \int dx_1 dx_2 F_{12} = n_1 \quad (35)$$

$$\frac{1}{2\pi} \int dx_3 dx_4 F_{34} = n_2. \quad (36)$$
Therefore this solution corresponds to \( n_1 \) coincident D2-branes and \( n_2 \) coincident D2'-branes.

The dynamics of the fluctuations around this solution is supersymmetric if they are restricted at the cores of the vortices, as in the case of parallel D-branes. This means that we can, in principle, construct the effective theory of D-brane system whose worldvolume is not a smooth manifold.

The intersecting D4-D4'-D4” system can also be described, starting from D6-\( \bar{\text{D}}6 \) system. The BPS conditions for this case are

\[
\begin{align*}
D_1 T - i s D_2 T &= 0 \\
D_3 T - i s' s' D_4 T &= 0 \\
D_5 T - i s s'' D_6 T &= 0 \\
F_{12} + s' F_{34} + s'' F_{56} + s(|T|^2 - \zeta) &= 0 \\
F_{13} - s' F_{24} &= 0 \\
F_{14} + s' F_{23} &= 0 \\
F_{15} - s'' F_{26} &= 0 \\
F_{16} + s'' F_{25} &= 0 \\
s' F_{35} - s'' F_{46} &= 0 \\
s' F_{36} + s'' F_{45} &= 0 \\
(s', s'' = \pm 1).
\end{align*}
\]

As in the D2-D2’ case, the last six equations are redundant and the only nontrivial equation is rewritten as follows,

\[
-\frac{1}{2} \sum_{k=1}^{6} (\partial_k)^2 \log \rho + \rho - \zeta = 0,
\]

where \( T = \rho^{1/2} e^{i\omega} \). The analysis of the solutions can be done similarly as for the previous case.

6. Discussions

We have discussed the discription of the BPS brane-system via the tachyon condensation. The model (\[\]) can describe both parallel branes and intersecting branes. Moreover the system with different number of D-branes can be treated in the same model. The dynamics of the fluctuations which localize on the branes is supersymmetric, and this model can, in principle, provide a way of describing the effective theory of the brane-system whose worldvolume is not a smooth manifold.

For the BPS D-branes to exist, we have seen that the specific form of the tachyon potential is needed. The existence of BPS D-branes would also require similar restriction
on the tachyon potential in the case of more complicated DBI-like action [12]. This might provide some information of the profile of the tachyon potential.

D4-D4'-D4'' system is considered in the end of the previous section. To this system, one can add D0-branes without breaking any supersymmetries. The bound states of such a system was discussed in [13] and it is conjectured that this has four bosonic states and four fermionic states (although the system considered in [13] is the D-branes wrapped on some cycle in a Calabi-Yau manifold). If the model (1) is generalized to the non-Abelian gauge thoery, there would exist a solution which contains the D0-branes in addition to the D4-branes. Therefore, the above conjecture could be checked by quantizing the collective coordinates of such a solution.

As shown in section 3, $n$ parallel D-branes are described in terms of $U(1)$ gauge theory. The generic vortex solution corresponds to the separated D-branes and the gauge symmetry on each worldvolume is $U(1)$. It should be expected that when two or more vortices coincide, the gauge symmetry is enhanced. This will certainly be achieved in the non-Abelian model. This gauge enhancement might occur already in the $U(1)$ model, however.

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Appendix : Vortex solutions

The Nielsen-Olesen vortex solution is the solution of the following equations [14],

\[ D_1 T - iD_2 T = 0 \]  
\[ F_{12} + |T|^2 - \zeta = 0, \]

where \( D_k T = \partial_k T + iA_k T \) \((k = 1, 2)\). \( A_k \) and \( T \) can be written as

\[ T = \rho^\frac{1}{2} e^{i\omega} \]  
\[ A_k = \frac{1}{2} \varepsilon_{kl} \partial_l \log \rho - \partial_k \omega. \]

\( \rho \) is determined by

\[ -\frac{1}{2} \left( \partial_1^2 + \partial_2^2 \right) \log \rho + \rho - \zeta = 0. \]

We now consider the radially symmetric solutions. Taking the polar coordinate \((r, \theta)\) and assuming that \( \rho \) is independent of \( \theta \), (41) becomes

\[ -\frac{1}{2} \left( \frac{d^2}{dr^2} + \frac{1}{r} \frac{d}{dr} \right) \log \rho + \rho - \zeta = 0. \]

\( \rho \) is a monotonic function and behaves as

\[ \rho \sim \left\{ \begin{array}{ll} r^n & (r \to 0) \\ \zeta & (r \to \infty) \end{array} \right\}, \]

where \( n \) is a positive integer. This behavior makes \( A_k \) singular at \( r = 0 \). To eliminate this singularity,

\[ \omega = \frac{n}{2} \theta. \]

For the single-valuedness of \( T \), \( n \) must be an even integer.

General solutions are also known [14]. The \( n \)-vortex solution is determined by specifying the positions of \( n \) points in the plane, at which \( T = 0 \) (cores of the vortices).
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