Brane with transverse rotation and background fields: boundary state and tachyon condensation

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Abstract The boundary state corresponding to the $D_p$-brane with a transverse rotation in the presence of the Kalb–Ramond and tachyon background fields and a $U(1)$ internal field will be constructed. We shall investigate effects of the open string tachyon condensation on this brane via its boundary state. We demonstrate that the background fields and transverse rotation cannot protect the brane against the collapse. Our calculations are in the context of the bosonic string theory.

1 Introduction

Some significant and important steps have been made to introduce the D-branes as essential objects in the string theory [1–4]. One of the main problems of the D-branes is their stability. The fate of an unstable D-brane can be investigated via the dynamics of the open string tachyon, i.e., the tachyon condensation process [5–12]. An unstable D-brane usually decays to another lower dimensional unstable D-brane as an intermediate state [13–16]. This intercurrent state eventually collapses to the closed string vacuum or decays to a lower dimensional stable configuration. There are various trustworthy approaches for studying these concepts, e.g.,: string field theory [17–20], the first quantized string theory [5–12, 21, 22], the renormalization group flow method [23–25], and the boundary string field theory [14, 17–20, 26, 27].

On the other hand, we have the boundary state formalism for describing the D-branes [28–67]. A boundary state prominently encodes all properties of its corresponding D-brane, and is a source for emitting all closed string states. Thus, this adequate state can be used to study the time evolution of the brane during the tachyon condensation process [46–51]. Note that the rolling tachyon has a boundary state description which is valid during the finite time. Therefore, after elapsing this time the energy of the system will be completely dissipated into the bulk [47, 48].

Among the various D-branes the dynamical-dressed branes motivated us to examine their behaviors under the tachyon condensation experience. This stimulation is due to the background fields and dynamics of such branes. Thus, in this paper we shall consider a single $D_p$-brane with a transverse rotation, which has been dressed with the Kalb–Ramond field, a $U(1)$ gauge potential and an open string tachyon field. The boundary state, associated with this $D_p$-brane, enables us to study the response of it in conflicting with the tachyon condensation phenomenon. We shall observe that the rotation of the brane and its field-dressing do not induce a resistance to protect it against the collapse. That is, the dimensional reduction of the brane will drastically occur.

This paper is organized as follows. In Sect. 2, the boundary state, corresponding to a rotating $D_p$-brane with the foregoing background fields, will be constructed. In Sect. 3, evolution of this $D_p$-brane under the condensation of the open string tachyon will be investigated. Section 4 is devoted to the conclusions.

2 The boundary state corresponding to our dynamical-dressed $D_p$-brane

We begin with the following closed string action

$$S = -\frac{1}{4\pi\alpha'} \int_{\Sigma} d^2 \sigma \left( \sqrt{-h} h^{ab} g_{\mu\nu} \partial_a X^\mu \partial_b X^\nu + \epsilon^{ab} B_{\mu\nu} \partial_a X^\mu \partial_b X^\nu \right) + \frac{1}{2\pi\alpha'} \int_{\partial\Sigma} d\sigma \left( A_a \partial_b X^a + T^2(X^a) \right),$$

(2.1)

where $\Sigma$ is the worldsheet of a closed string which is emitted by a static $D_p$-brane, and $\partial\Sigma$ is the boundary of it. The coordinates $\{x^a| a = 0, 1, \ldots, p\}$ specify the directions which are along the worldvolume of this brane, and the set
In fact, the rotation of the perpendicular coordinates to a brane worldvolume, which describes a spinning brane, deforms the metric of the background spacetime. However, imposing some other motions to the brane does not change the metric. For example, in the flat spacetime see the D-branes with transverse velocities [2, 35, 62–65], the rotated D-branes [63], the D-branes with tangential rotations [66, 67], and so on. In fact, in these examples the first approximation of the background metric has been manifestly applied. Our D-brane will rotate in a transverse plane to itself, thus, for a small angular velocity we have a quasi-static D-brane. Hence, similar to the foregoing examples, at least for such small rotations we can apply the first approximation of the metric. This elaborates that the equations of the boundary state, corresponding to the rotating brane with the transverse rotation, and also the equation of motion of an emitted closed string from the brane will be reliably written in the initial flat spacetime.

Here we compare effects of the different dynamics of the branes on the background metric. In a spinning brane for each rotating perpendicular plane to the brane there exists one rotational parameter. These adequate variables induce a deformation to the spacetime metric. For a boosted brane the velocity components are introduced, which do not deform the metric. In our system the only rotational parameter is the angular velocity. By employing the foregoing first approximation of the metric, accompanied by the quasi-static rotation of the brane, one can see that the flat metric under the transformations (2.3) remains invariant.

Now we impose a transverse rotation to the brane. Let $x^{i_0}$ be the horizontal axis and $x^{a_0}$ (with $a_0 \neq 0$) be the vertical one. At the time $t = 0$ the direction $x^{a_0}$ is along the brane, and the direction $x^{i_0}$ is perpendicular to it. The brane is rotating, e.g. counterclockwise, with the constant angular velocity \( \omega \). The axis of the rotation is one of the normal directions to the plane $x^{i_0}x^{a_0}$. The coordinate system \( \{x^{i\mu}\} \) is stuck to the brane such that at each moment the planes $x^{i_0}x^{a_0}$ and $x^{i_0}x^{\tilde{a}0}$ have common origin and they are coincident. Thus, we receive the following coordinate transformations

$$
\begin{align*}
    x^{i_0} &= x^{i_0} \cos(\omega t) + x^{a_0} \sin(\omega t), \\
    x^{a_0} &= -x^{i_0} \sin(\omega t) + x^{a_0} \cos(\omega t), \\
    x^{\tilde{a}} &= x^{\tilde{a}}, \\
    x^{\tilde{i}} &= x^{\tilde{i}},
\end{align*}
$$

(2.3)

where the new indices $\tilde{a}$ and $\tilde{i}$ belong to the sets

$$
\tilde{a} \in \{0, 1, \ldots, p\} - \{a_0\}, \\
\tilde{i} \in \{p + 1, \ldots, d - 1\} - \{i_0\}.
$$

The equation of motion, extracted from the action (2.1) for the flat spacetime and worldsheet, clearly takes the form \( (\partial_t^2 - \partial_\tau^2)X^\mu(\sigma, \tau) = 0 \). By applying the transformations (2.3) and the quasi-static approximation for the brane dynamics we receive \( (\partial_t^2 - \partial_\tau^2)X^{\tilde{\mu}}(\sigma, \tau) = 0 \). Besides, the quasi-static
rotating Dp-brane possesses the following boundary state equations

\[
\begin{align*}
\partial_\tau X^\alpha + \frac{\mathcal{F}^\alpha}{p} \partial_\sigma X^\beta \\
+ \frac{\mathcal{F}^\alpha}{a_0} \cos(\omega t) \left(- \sin(\omega t) \partial_\tau X^i + \cos(\omega t) \partial_\sigma X^{a_0} \right) \\
- 2\pi i \alpha' U^\alpha_{\beta} \bar{X}^\beta X^i \\
+ 4\pi i \alpha' U^a_{\alpha} \cos(\omega t) \left(- \alpha'_m \sin(\omega t) \right) \\
+ \mathcal{X}^{a_0} \cos(\omega t) \right]_{\tau=0} |B(t)\rangle = 0, \\
\left[ X^i \cos(\omega t) + X^{a_0} \sin(\omega t) \right]_{\tau=0} |B(t)\rangle = 0, \\
\left( \mathcal{X}^i - \mathcal{Y}^i \right)_{\tau=0} |B(t)\rangle = 0.
\end{align*}
\]

for the zero-mode part, and we have

\[
\begin{align*}
\left[ \alpha'_m + \tilde{\alpha}'_{-m} - \frac{2\pi \alpha'}{m} U^\alpha_{\beta} \right] (\alpha'^m_m - \tilde{\alpha}'_{-m}^m) \\
+ \left( \frac{\mathcal{F}^\alpha}{a_0} - \frac{2\pi \alpha'}{m} U^a_{a_0} \right) \cos(\omega t) [ (\alpha'^i + \tilde{\alpha}'_{-i}) \sin(\omega t) ] \\
- (\alpha'^m_m - \tilde{\alpha}'_{-m}^m) \cos(\omega t) ] |B(t)\rangle^{(osc)} = 0, \\
\left[ (\alpha'^m_m + \tilde{\alpha}'_{-m}^m) \cos(\omega t) - (\alpha'^i + \tilde{\alpha}'_{-i}) \sin(\omega t) \right] \\
- \frac{2\pi \alpha'}{m} U^a_{a_0} \cos^2(\omega t) \left[ (\alpha'^m_m - \tilde{\alpha}'_{-m}^m) \sin(\omega t) \right] \\
- (\alpha'^m_m - \tilde{\alpha}'_{-m}^m) \cos(\omega t) ] |B(t)\rangle^{(osc)} = 0, \\
\left[ (\alpha'^m_m - \tilde{\alpha}'_{-m}^m) \cos(\omega t) + (\alpha'^i + \tilde{\alpha}'_{-i}) \sin(\omega t) \right] |B(t)\rangle^{(osc)} = 0, \\
\left( \alpha'^m_m - \tilde{\alpha}'_{-m}^m \right) |B(t)\rangle^{(osc)} = 0.
\end{align*}
\]

for the oscillating part, with \( m \in \mathbb{Z} - \{0\} \). Note that we decomposed the boundary state to the zero-mode portion and the oscillating part, i.e., \( |B(t)\rangle = |B(t)\rangle^{(0)} \otimes |B(t)\rangle^{(osc)} \).

In fact, solving Eqs. (2.5) and (2.6) is very difficult. For simplification we impose the restriction \( U_{a_0a_0} = 0 \), or equivalently \( U_{a_0a_0} = U_{a_0a_0} = 0 \). Therefore, the solution of the zero-mode part of the boundary state is given by

\[
|B(t)\rangle^{(0)} = \frac{1}{\sqrt{\det U}} \int_0^\infty \exp \left[ - \frac{1}{4\pi} \sum \xi(t) t \right] \left( p^\xi \right)^2 \\
- \frac{1}{2\pi} \sum \left( \xi(t) \right) \left( \prod \left| p^\xi \right| dp^\xi \right) \\
\times \left( \mathcal{X}^i \cos(\omega t) + \mathcal{X}^{a_0} \sin(\omega t) \right) \\
\times \prod \left( \delta(x^i - y^i) |p^0 = 0\rangle \otimes |p^{a_0} = 0\rangle \right)_{\tau=0} |B(t)\rangle^{(0)} = 0, \\
\left( \mathcal{X}^i - \mathcal{Y}^i \right)_{\tau=0} |B(t)\rangle^{(0)} = 0.
\]

where, according to the condition \( U_{a_0a_0} = 0 \), the \( p \times p \) symmetric matrix \( U \) is defined by eliminating the \( a_0 \)th column and \( a_0 \)th row of the \((p+1) \times (p+1)\) tachyon matrix \( U \). From the disk partition function we deduce the prefactor \( 1/\sqrt{\det U} \) [52, 53]. The exponential part of \( |B(t)\rangle^{(0)} \), which is absent for the conventional boundary states, clearly is an effect of the tachyon field. We observe that the zero-mode part of the boundary state is independent of the total field strength and the parameter \( \alpha' \). This is due to the fact that we considered a non-compact brane. The compact case extremely contains these factors [36–38].
For solving Eq. (2.6) we define the new oscillators
\[
A_m^\alpha = a_m^\alpha \cos(\omega t) + \tilde{a}_m^\alpha \sin(\omega t),
\]
\[
\tilde{A}_m^\alpha = \tilde{a}_m^\alpha \cos(\omega t) + a_m^\alpha \sin(\omega t),
\]
\[
B_m^\alpha = a_m^\alpha \cos(\omega t) + \tilde{a}_m^\alpha \sin(\omega t),
\]
\[
\tilde{B}_m^\alpha = \tilde{a}_m^\alpha \cos(\omega t) - a_m^\alpha \sin(\omega t).
\]

These oscillators possess the following nonzero commutators
\[
[A_m, A_n] = [\tilde{A}_m, \tilde{A}_n] = [B_m, B_n] = [\tilde{B}_m, \tilde{B}_n] = m \delta_{m+n,0},
\]
and all other commutators among them vanish.

By multiplying the coherent state method, and after some
heavy calculations, the oscillating part of the boundary state
finds the feature
\[
|B(t)|^{(osc)} = \frac{T_p}{g_s} \prod_{n=1}^{\infty} \left[ \det \left( 1 - F'(t) + \frac{2\pi \alpha'}{n} \tilde{U} \right) \right]^{-1} \times \exp \left\{ -\sum_{m=1}^{\infty} \frac{1}{m} \left[ \alpha^\alpha_m Q_{(m)}^{\alpha\beta} \tilde{a}_m^\beta - \alpha^\dagger_m \tilde{a}^\dagger_m \right. \right.
\]
\[
- A_m \tilde{A}_m
\]
\[
+ \left. \left( 1 + 2 \left( M_m^{-1} \right)^{\alpha\beta}_{\bar{\alpha}\bar{\beta}} \tilde{F}^\bar{\alpha}\bar{\beta} a_0 \tilde{F}^\alpha\beta a_0 \cos^2(\omega t) \right) B_m \tilde{B}_m
\]
\[
+ 2 \tilde{F}^\bar{\beta}_\alpha \cos(\omega t) \left( \left( M_m^{-1} \right)^{\alpha\beta}_{\bar{\alpha}\bar{\beta}} \tilde{a}_m^\alpha \tilde{B}_m \right)
\]
\[
- \left. \left( M_m^{-1} \right)^{\alpha\beta}_{\bar{\alpha}\bar{\beta}} B_m \tilde{a}_m^\alpha \right) \right\} |0\rangle,
\]

where by putting the \(\alpha_0\)th column and \(\alpha_0\)th row of the tachyon
matrix \(\tilde{U}\) to zero the \((p+1) \times (p+1)\) matrix \(\tilde{U}\) is obtained.

By multiplying the \(\alpha_0\)th column and \(\alpha_0\)th row of the field
strength matrix \(F\) with \(\cos(\omega t)\) we receive the matrix \(F'(t)\).
The other matrices are defined by
\[
Q_{(m)}^{\alpha\bar{\alpha}} = \left( M_m^{-1} N_m \right)^{\alpha\bar{\alpha}},
\]
\[
M_m^{\alpha\bar{\beta}} = -\tilde{q}_\beta^\beta \tilde{F}^\beta_\alpha + \frac{2\pi \alpha'}{m} U_\beta^\beta - \tilde{F}^\beta_\alpha a_0 \tilde{F}^\alpha\beta \cos^2(\omega t),
\]
\[
N_m^{\alpha\bar{\beta}} = \tilde{q}_\alpha^\bar{\alpha} \tilde{F}^\alpha_\beta + \frac{2\pi \alpha'}{m} U_\beta^\beta - \tilde{F}^\alpha_\beta a_0 \tilde{F}^\alpha\beta \cos^2(\omega t).
\]

As we see these matrices depend on the mode number “m”
which is induced by the tachyon matrix. The normalization
factor, i.e., the infinite product in the first line of Eq. (2.10),
is originated by the disk partition function. The state (2.10)
specifies that \(A_m\) and \(\tilde{A}_m\) are Dirichlet oscillators. Similarly,
in the case \(\tilde{F}^\alpha_\beta a_0 = 0\) the variables \(B_m\) and \(\tilde{B}_m\) became Neumann
oscillators.

In Eq. (2.6) one can express the right-moving annihilation
oscillators in terms of the left-moving creation oscillators.
This obviously eventuates to the boundary state (2.10).
However, in these equations it is possible to express the
left-moving annihilation oscillators in terms of the right-moving
creation oscillators. In this case, applying the coherent state
method leads to another form for the boundary state of the
oscillating part. Equality of these boundary states elaborates
the following conditions
\[
M_m M_m^{T} = N_m N_m^{T},
\]
\[
2 \left( M_m^{-1} \right)^{\alpha\beta}_{\bar{\alpha}\bar{\beta}} \tilde{F}^\beta_\alpha a_0 = -\tilde{F}^\alpha_\beta a_0 \left( Q_m + 1 \right) \tilde{F}^\alpha_\beta a_0,
\]
\[
2 \left( N_m^{-1} \right)^{\alpha\beta}_{\bar{\alpha}\bar{\beta}} \tilde{F}^\beta_\alpha a_0 = -\tilde{F}^\alpha_\beta a_0 \left( Q_m + 1 \right) \tilde{F}^\alpha_\beta a_0,
\]
(2.12)

In fact, by substituting the explicit forms of the matrices from
Eqs. (2.11) and (2.13) into Eq. (2.12) we see that the first, the
two and the third equations of (2.12) are trivial identities.
That is, they do not impose any relation among the parameters
of our setup. For the odd values of the brane dimension “p”
the fourth equation is an identity, and for the even values of
“p” it only gives rise to the condition \(\det (F_{\alpha\beta}) = 0\).

Note that the total boundary state includes a part which
comes from the conformal ghosts. This portion manifestly is
independent of the background fields and the brane rotation.
Thus, it obviously is null under the tachyon condensation
process. Hence, we shall not consider it.

3 Effect of the tachyon condensation on our Dp-brane

According to the Sen’s papers [5–12], in the presence of the
open string tachyonic field our knowledge about the vacua of
the string theories, the fate of the D-branes, their instability,
and so on, was improved. During the process of the tachyon
condensation the brane drastically collapses, and finally we
receive a collection of the closed strings. These imply that
decadence of unstable objects is very important phenomenon.
For example, these objects specify an approach to achieve the
background independent formulation of string theory.

Since the boundary state is a source for emitting all quantum
states of closed string, and accurately describes all properties
of the corresponding brane, and comprises a specific
normalization factor, it is a favorable and convenient tool for
finding the treatment and behavior of a single D-brane under the experience of the tachyon condensation. Hence, in this section we shall use this adequate formalism.

For imposing the condensation on the tachyon field, some of the matrix elements of the tachyon matrix should be infinite. For this purpose let the system tend to the infrared fixed point via the limit \( U_{pp} \to \infty \). This defines the tachyon condensation along the \( x^\rho \)-direction, where we assume \( x^{a_0} \neq x^\rho \). Now we should take the limit of the total boundary state to acquire the behavior of our dynamical-dressed D\( p \)-brane under the tachyon condensation process.

At first we obtain the behavior of the zero-mode part of the boundary state, i.e., Eq. \( (2.7) \). Under the limit \( U_{pp} \to \infty \) its prefactor transforms to

\[
\frac{1}{\sqrt{U_{pp} \det U}}.
\]

where the \( (p - 1) \times (p - 1) \) symmetric matrix \( \tilde{U} \) is defined by eliminating the last and \( a_0 \)th columns and also the last and \( a_0 \)th rows of the tachyon matrix \( U \). At the IR fixed point limit we have

\[
\lim_{U_{pp} \to \infty} \tilde{U}^{-1} = \left( \begin{array}{cc} \tilde{U}^{-1} & 0 \\ 0 & 1 \end{array} \right).
\]

Adding all these together we receive the limit

\[
|B(t)|^{(0)} = \frac{2\pi}{\sqrt{U_{pp} \det U}} \int_{-\infty}^{\infty} \exp \left[-\frac{1}{4\pi} \sum_{a} \left( \tilde{U}^{-1} \right)_{aa} \left( p^a \right)^2 \right]
\]

\[
-\frac{1}{2\pi} \sum_{a \neq b} \left( \tilde{U}^{-1} \right)_{ab} p^a p^b \left( \prod_{a} |p^a| d p^a \right)
\]

\[
\times \delta \left[ \chi^0 \cos(\omega t) + x^{a_0} \sin(\omega t) \right]
\]

\[
\times \prod_{i} \delta (x^i - y^i) |p^i| = 0 \]

\[
\otimes |p^{a_0} = 0 \rangle \otimes |p^{a_0} = 0 \rangle,
\]

where the profitable relation \( 2\pi \sqrt{\chi^0} T_p = T_{p-1} \) was used. For the forms of the \( p \times p \) matrices \( \tilde{F}' \) and \( \tilde{U} \), eliminate the last rows and last columns of the \( (p + 1) \times (p + 1) \) matrices \( \tilde{F}' \) and \( \tilde{U} \), respectively.

For calculating the limit of \( M_m^{-1} \) we use

\[
\lim_{U_{pp} \to \infty} \det M_m = \frac{2\pi \alpha'}{m} U_{pp} \det M_m^{(p-1)}.
\]

where by eliminating the last row and last column of \( M_m \) the \( (p - 1) \times (p - 1) \) matrix \( M_m^{(p-1)} \) is acquired. Since the last row and last column of \( M_m^{-1} \) contain the factor \( 1/U_{pp} \), we receive

\[
\lim_{U_{pp} \to \infty} M_m^{-1} = \left( \begin{array}{cc} (M_m^{(p-1)})^{-1} & 0 \\ 0 & 1 \end{array} \right).
\]

Beside, the limit of \( Q_m = M_m^{-1} N_m \) finds the feature

\[
\lim_{U_{pp} \to \infty} Q_m = \left( \begin{array}{cc} (M_m^{(p-1)})^{-1} N_m^{(p-1)} & 0 \\ 0 & 1 \end{array} \right).
\]

Note that the limit of the matrix \( Q_m \) is not the product of the limits of \( M_m^{-1} \) and \( N_m \). After performing the product \( M_m^{-1} N_m \) we have taken the limit of \( Q_m \). The structures of the matrices \( M_m^{(p-1)}, N_m^{(p-1)} \) and \( Q_m^{(p-1)} \) are similar to the matrices of Eq. \( (2.11) \) in which \( \tilde{a} \) and \( \tilde{b} \) must be replaced with the indices \( a, b \in \{ 0, 1, \ldots, p \} - \{ a_0, p \} \).

Adding all these together, the effect of the tachyon condensation on the oscillating part of the boundary state is given by

\[
|B(t)|^{(osc)} = \frac{T_p \sqrt{U_{pp}}}{g_s} \prod_{n=1}^{\infty} \left[ \det \left( 1 - \tilde{F}' + \frac{2\pi \alpha'}{n} \tilde{U} \right) \right]^{-1}
\]

\[
\times \exp \left[ -\sum_{m=1}^{\infty} \left( \alpha^m \left( Q_m^{(p-1)} \right)_{\tilde{a}b, m} \right) \right]
\]

\[
-\tilde{a}_{m} \tilde{a}_m \tilde{a}_{m} + \tilde{a}_{m} \tilde{a}_m \tilde{a}_{m} - A_{m} \tilde{A}_{m}
\]

\[
+ 2 \tilde{F}_{ab} \cos(\omega t) \left[ \left( M_m^{(p-1)} \right)^{-1} \right]_{ab} \tilde{a}_{m} \tilde{B}_{m}
\]

\[
\left( \left( M_m^{(p-1)} \right)^{-1} \right)_{ab} \tilde{B}_{m}
\]

\[
+ [1 + 2 \left( M_m^{(p-1)} \right)^{-1}] \tilde{F}_{ab} \tilde{a}_m \tilde{a}_m \cos(\omega t) \left( B_{m} \tilde{B}_{m} \right) \right] \otimes |0\rangle.
\]

As expected, the sign of the operator \( \alpha_{n-m} \tilde{a}_{n-m} \) has changed, i.e., the previous Neumann direction \( x^\rho \) has been transformed to a Dirichlet direction. By comparing this equation with Eq. \( (2.10) \) we observe that, up to the factor \( \sqrt{U_{pp}} \), Eq. \( (3.7) \) manifestly describes the oscillating part of the boundary state which is corresponding to the D\( (p - 1) \)-brane.

For the total boundary state at the IR fixed point the extra factors \( 1/\sqrt{U_{pp}} \) and \( \sqrt{U_{pp}} \) of Eqs. \( (3.2) \) and \( (3.7) \) exactly
cancel each other. Similar cancellation between the zero-mode portion and the oscillating part also occurs in the D-\(\bar{D}\) systems [20,26,27,68]. However, according to the product of the states (3.2) and (3.7) we have proved that, during the lapse. That is, the unstable D\(p\)-brane lost its \(x^p\)-direction and conveniently reduced to a D\((p-1)\)-brane. The resulted brane is rotating inside the \(x^b x^{p_0}\)-plane with the same frequency “\(\omega\)”. The delta functions of Eq. (3.2) prominently clarify that this D\((p-1)\)-brane has been localized at the position \(x^p = 0\), \(x^i = y^i\), and its configuration at the times \(t \in \{\frac{2\pi n}{\omega} | n \in \mathbb{Z}\}\) is along the directions \([x^1, x^2, \ldots, x^{p-1}]\).

4 Conclusions

In the framework of the bosonic string theory we constructed a profitable boundary state, associated with a dynamical D\(p\)-brane with a transverse rotation, in the presence of the antisymmetric tensor field \(B_{\mu\nu}\), a \(U(1)\) internal gauge potential and a tachyonic field of the open string spectrum. Though we imposed a uniform rotation to the brane but the time dependence of the corresponding boundary state is very intricate. Besides, the rotational dynamics induced the deformed versions of the tachyon matrix and total field strength to the boundary state.

We investigated the effects of the tachyon condensation on the foregoing D\(p\)-brane through its boundary state. We demonstrated that at the infrared fixed point the background fields, accompanied by the transverse rotation of the brane, cannot prevent the unstable brane against the collapse. Therefore, the tachyon condensation was eventually terminated by the dimensional reduction of the brane. The resulted D\((p-1)\)-brane possessed the same angular frequency as the previous one. Presence of the remaining tachyon field implies that the subsequent brane also is an unstable object, and at the IR fixed point will be collapsed.

Data Availability Statement This manuscript has no associated data or the data will not be deposited. [Authors’ comment:This paper is a theoretical work. No experimental data were used].

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References

1. J. Polchinski, String Theory, vol. I (Cambridge University Press, Cambridge, 1998)
2. J. Polchinski, String Theory, vol. II (Cambridge University Press, Cambridge, 1998)
3. C.V. Johnson, D-Branes (Cambridge University Press, Cambridge, 2003)
4. J. Polchinski, Phys. Rev. Lett. 75, 4724 (1995)
5. A. Sen, Int. J. Mod. Phys. A 14, 4061 (1999)
6. A. Sen, Int. J. Mod. Phys. A 20, 5513 (2005)
7. A. Sen, JHEP 9808, 010 (1998)
8. A. Sen, JHEP 9808, 012 (1998)
9. A. Sen, JHEP 9912, 027 (1999)
10. A. Sen, JHEP 9809, 023 (1998)
11. A. Sen, JHEP 9812, 021 (1998)
12. A. Sen, JHEP 9910, 008 (1999)
13. K. Hashimoto, P.M. Ho, J.E. Wang, Mod. Phys. Lett. A 20, 79 (2005)
14. D. Kutasov, M. Marino, G. Moore, JHEP 0010, 045 (2000)
15. T. Lee, Phys. Rev. D 64, 106004 (2001)
16. T. Lee, Phys. Lett. B 520, 385 (2001)
17. E. Witten, Phys. Rev. D 47, 3405 (1993)
18. E. Witten, Phys. Rev. D 46, 5467 (1992)
19. E. Witten, Nucl. Phys. B 268, 253 (1986)
20. P. Kraus, F. Larsen, Phys. Rev. D 63, 106004 (2001)
21. E. Witten, JHEP 12, 019 (1998)
22. O. Bergman, M.R. Gaberdiel, Phys. Lett. B 441, 133 (1998)
23. J.A. Harvey, D. Kutasov, E.J. Martinec, arXiv:hep-th/0003101
24. S. Dasgupta, T. Dasgupta, arXiv:hep-th/0010247
25. A.A. Gerasimov, S.L. Shatashvili, JHEP 10, 034 (2000)
26. S.L. Shatashvili, Phys. Lett. B 311, 83 (1993)
27. T. Takayanagi, S. Terashima, T. Uesugi, JHEP 03, 019 (2001)
28. M.B. Green, M. Gutperle, Nucl. Phys. B 476, 484 (1996)
29. C. Schmidhuber, Nucl. Phys. B 467, 146 (1996)
30. M.B. Green, P. Wai, Nucl. Phys. B 431, 131 (1994)
31. M. Bilbo, P. Di Vecchia, M. Frau, A. Lerda, I. Pesando, R. Russo, S. Sciuto, Nucl. Phys. B 526, 199 (1998)
32. P. Di Vecchia, M. Frau, I. Pesando, S. Sciuto, A. Lerda, R. Russo, Nucl. Phys. B 507, 259 (1997)
33. F. Hassanain, R. Iengo, C. Nunez, Nucl. Phys. B 497, 205 (1997)
34. O. Bergman, M. Gaberdiel, G. Lifschytz, Nucl. Phys. B 509, 194 (1998)
35. M. Bilbo, P. Di Vecchia, D. Cangemi, Phys. Lett. B 400, 63 (1997)
36. H. Arfaei, D. Kamani, Phys. Lett. B 452, 54 (1999)
37. H. Arfaei, D. Kamani, Nucl. Phys. B 561, 57 (1999)
38. H. Arfaei, D. Kamani, Phys. Lett. B 475, 39–45 (2000)
39. D. Kamani, Phys. Lett. B 487, 187–191 (2000)
40. D. Kamani, Europhys. Lett. 57, 672–676 (2002)
41. D. Kamani, Nucl. Phys. B 601, 149–168 (2001)
42. D. Kamani, Mod. Phys. Lett. A 15, 1655–1664 (2000)
43. H. Arfaei, D. Kamani, Nucl. Phys. B 907, 109702 (2014)
44. M. Saidy-Sarjoubi, D. Kamani, Nucl. Phys. B 922, 280 (2017)
45. M. Naka, T. Takayanagi, T. Uesugi, JHEP 0006, 007 (2000)
46. T. Okuda, S. Sugimoto, Nucl. Phys. B 647, 101 (2002)
47. F. Larsen, A. Naqvi, S. Terashima, JHEP 0302, 039 (2003)
48. A. Sen, JHEP 0204, 048 (2002)
49. A. Sen, JHEP 0207, 065 (2002)
50. D. Kamani, Ann. Phys. 354, 394–400 (2015)
52. C.G. Callan, C. Lovelace, C.R. Nappi, S.A. Yost, Nucl. Phys. B 308, 221–284 (1988)
53. E.S. Fradkin, A.A. Tseytlin, Phys. Lett. B 163, 123 (1985)
54. J. Dai, R.G. Leigh, J. Polchinski, Mod. Phys. Lett. A 4, 2073 (1989)
55. R.G. Leigh, Mod. Phys. Lett. A 4, 2767 (1989)
56. S.J. Rey, S. Sugimoto, Phys. Rev. D 68, 026003 (2003)
57. M. Frau, I. Pesando, S. Sciuto, A. Lerda, R. Russo, Phys. Lett. B 400, 52 (1997)
58. M. Li, Nucl. Phys. B 460, 351 (1996)
59. S. Gukov, I.R. Klebanov, A.M. Polyakov, Phys. Lett. B 423, 64 (1998)
60. T. Kitao, N. Ohta, J.G. Zhou, Phys. Lett. B 428, 68 (1998)
61. C. Bachas, Phys. Lett. B 374, 37 (1996)
62. C.G. Callan, I.R. Klebanov, Nucl. Phys. B 465, 473 (1996)
63. P. Di Vecchia, A. Liccardo, D branes in string theory, II, YITP Proceedings Series No. 4 (Kyoto, 1999)
64. F. Hussain, R. Iengo, C. Nunez, C.A. Scrucca, Nucl. Phys. B 517, 92 (1998)
65. F. Hussain, R. Iengo, C. Nunez, C.A. Scrucca, Phys. Lett. B 409, 101 (1997)
66. F. Safarzadeh-Maleki, D. Kamani, Phys. Rev. D 89, 026006 (2014)
67. E. Maghsoodi, D. Kamani, Nucl. Phys. B 942, 381 (2019)
68. S.P. de Alwis, Phys. Lett. B 505, 215 (2001)