We study quantum corrections at the one loop level in open superstring tachyon condensation using the boundary string field theory (BSFT) method. We find that the tachyon field action has the same form as at the disc level, but with a renormalized effective coupling $\lambda' = g_s^{\text{ren}} e^{-T^2/4}$ ($g_s^{\text{ren}}$ is the renormalized dimensionless closed string coupling, $e^{-T^2/4}$ is the tachyonic field expectation value) and with an effective string tension. This result is in agreement with that based on general analysis of loop effects.
1 Introduction

Spontaneous symmetry breaking in dual models leading to removal of tachyonic vacuum instability has been studied some years ago [1-4]. Recently, the problem of tachyon condensation on unstable D-branes and on D-\(\bar{D}\) system has been the subject of active investigation. Two distinct methods have been used to study this problem. The first, earlier method, uses string field theory. The main approach here has been through Witten’s cubic string field theory [5] and remarkable results have been obtained using this approach [6, 7]. It is shown that as a result of tachyon condensation the unstable D-brane decays into the perturbative vacuum of closed strings or to lower dimensional D-brane. It has been further conjectured that the tachyon potential on a D-\(\bar{D}\) system in type II string theory has a universal form [8] independent of the boundary conformal field theory describing the D-brane, and that the open string field theory at the tachyonic vacuum must contain closed string states [9]. In Witten’s string field theory approach, tachyon condensation phenomena in general involves giving expectation values to an infinite number of components of the string field, and as a consequence one has to resort to level truncation [5].

A second approach, which appears simpler, involves using the background independent string field theory (BSFT) of Witten [10, 11] and of Shatashvili [12]. In this version of the open string field theory, the classical configuration space is the space of two dimensional world sheet theories which are conformal in the bulk (of disc topology) but contain arbitrary boundary interaction terms. Since the tachyon condensation is an off-shell process, the BSFT approach is well suited to studying this problem. A large number of papers have been written using this approach [13-21]. The BSFT method has been extended to the case of superstrings in [17, 23-27]. The tree level tachyonic action has been shown to be given by

\[
S_{\text{tree}} = T_9 \int d^{10}X e^{-\frac{4}{T^2}}[(2 \ln 2)\partial_\mu T \partial^\mu T + 1],
\]

where \(T_9\) is the tension of the non-BPS D9-brane.

The question of quantum corrections to tachyonic action deserves a detailed study. Partial studies of one loop effects in bosonic case are contained in [28, 29]. Earlier studies...
on the annulus effective action are based on the sigma model approach to string theory \[30\], and also \[21, 31\], where they calculate the Dirac-Born-Infeld (DBI) action for the D-brane. We use the BSFT approach to calculate the tachyonic field action in the superstring theory at one loop level. The partition function \(Z\) of the BSFT takes the form \[30\]

\[
Z = \sum_{\chi} g_s^{-\chi} \int d\mu(\lambda) \int e^{-I_{\text{bulk}}-I_{\text{bndy}}} [dx^\mu] [d\psi] [d\tilde{\psi}],
\]

(2)

where \(\lambda_1, \ldots, \lambda_N\) denote Teichmüller parameters of the two dimensional surface with holes \((\chi = 1, N = 0; \chi = 0, N = 1; \chi \leq 1, N = -3\chi)\). We are interested in calculating the annulus partition function where \(\chi = 0\) and \(N = 1\). The tree level partition function has already been evaluated; see (1). \(g_s\) denotes the unrenormalized dimensionless closed string coupling. The measure \(\mu\) in (2) includes the contribution of the ghost determinant.

The bulk action \(I_{\text{bulk}}\) is given by the standard NSR action

\[
I_{\text{bulk}} = \frac{1}{4\pi} \int_{\Sigma} d^2z \left( \partial X^\mu \partial X_\mu + \psi^\mu \partial \psi_\mu + \tilde{\psi}^\mu \partial \tilde{\psi}_\mu \right),
\]

(3)

where \(\psi\) and \(\tilde{\psi}\) are the holomorphic and antiholomorphic fermionic fields. \(\Sigma\) is a two-dimensional world sheet with the coordinates \(X^\mu\). (Here we choose \(\alpha' = 2\) for the closed string.)

The supersymmetric boundary action \(I_{\text{bndy}}\) is given by \[23\] as

\[
I_{\text{bndy}} = \frac{1}{8\pi} \int_{\partial \Sigma} d\theta \left[ (T(X))^2 + (\psi^\mu \partial_\mu T) \frac{1}{\partial \theta} (\psi^\mu \partial_\mu T) + (\tilde{\psi}^\mu \partial_\mu T) \frac{1}{\partial \theta} (\tilde{\psi}^\mu \partial_\mu T) \right],
\]

(4)

where \(T(X)\) is a tachyon profile chosen to be of the form

\[
T(X) = a + u_\mu X^\mu,
\]

(5)

and

\[
\frac{1}{\partial \theta} (\psi^\mu \partial_\mu T) = \frac{1}{2} \int d\theta' \epsilon(\theta - \theta')(\psi^\mu \partial_\mu T)(\theta'),
\]

(6)

where \(\epsilon(x) = +1\) for \(x > 0\) and \(\epsilon(x) = -1\) for \(x < 0\).
It was first conjectured in [23] and later verified in [24, 25] that, using Batalin-Vilkovisky (BV) formalism [32] of background independent superstring field theory, the spacetime superstring field action is the partition function defined in (2) on the disc. Since the derivation of this result using the BV formalism appears to go through for other geometries (the annulus here), we assume that the string field action is given by the partition function defined in (2) on the world sheet with the topology of an annulus. We justify \textit{aposteriori} this assumption.

The plan of the paper is as follows. In section 2 we review the results for the disc. The bulk fermionic Green function in the NS sector is derived and the first two non-trivial terms, i.e. the kinetic and potential energy terms, in the tachyonic action are derived. In section 3, we first derive the required bosonic and fermionic Green functions on the annulus. The bosonic Green function has been derived in [28, 29, 34] previously. Using these the one loop tachyonic field action is derived. Section 4 has conclusions.

2 Tachyon Action at the Tree Level

In this section we choose the world sheet $\Sigma$ to be a disc with a rotationally invariant flat metric

\[
 ds^2 = d\sigma_1^2 + d\sigma_2^2, \tag{7}
\]

the complex variable $z = \sigma_1 + i\sigma_2$, with $|z| \leq 1$.

In (5), the constant $a$ can be shifted away by the integration constant, so we take boundary tachyon profile of the form

\[
 T(X) = u_\mu X^\mu. \tag{8}
\]

The boundary term can then be written as

\[
 I' = \frac{y}{8\pi} \int_0^{2\pi} d\theta \left( X^2 + \psi \frac{1}{\partial_\theta} \bar{\psi} + \bar{\psi} \frac{1}{\partial_\theta} \psi \right), \tag{9}
\]

where $y \equiv u^2$. 
The boundary conditions are derived by varying the action $I = I_{\text{bulk}} + I_{\text{bndy}}$. The bosonic Green function satisfies,

$$\partial_z \bar{\partial}_z G(z, w) = -2\pi \delta^{(2)}(z - w),$$

with the boundary condition

$$z \partial_z G_B + \bar{z} \partial_{\bar{z}} G_B + u G_B \bigg|_{\rho = 1} = 0.$$  \hspace{1cm} (11)

The boundary conditions for the fermionic Green functions are

$$(1 + iy \frac{1}{\partial_{\theta}}) G_F \bigg|_{\rho = 1} = (1 - iy \frac{1}{\partial_{\theta}}) \tilde{G}_F \bigg|_{\rho = 1},$$

where $G(z, w)$ and $\tilde{G}(\bar{z}, \bar{w})$ are the holomorphic and antiholomorphic Green function.

The fermionic Green functions satisfy the equations:

$$\partial G_F(z, w) = -i \sqrt{zw} \delta^{(2)}(z - w),$$
$$\partial \tilde{G}_F(\bar{z}, \bar{w}) = +i \sqrt{\bar{z} \bar{w}} \delta^{(2)}(\bar{z} - \bar{w})$$

with the boundary conditions (12). Note that the factors $\sqrt{zw}$ and $\sqrt{\bar{z} \bar{w}}$ before the $\delta$ function are inserted to make the Green function dimensionless.

To solve (13) with the boundary conditions (12) on the disc, we start with the ansatz,

$$i G_F(z, w) = \frac{\sqrt{zw}}{z - w} + \sum_r a_r (z \bar{w})^r,$$
$$i \tilde{G}_F(\bar{z}, \bar{w}) = -\frac{\sqrt{\bar{z} \bar{w}}}{\bar{z} - \bar{w}} - \sum_r b_r (\bar{z} \bar{w})^r.$$  \hspace{1cm} (14)

It is easy to verify that these are solutions of (13). $a_r$’s and $b_r$’s are coefficients to be determined by the boundary conditions (12).

Inserting this ansatz into the boundary conditions (12), expanding it by series and following the same procedure as in [28], we get

$$i G_F(z, w) = \frac{\sqrt{zw}}{z - w} - \frac{\sqrt{\bar{z} \bar{w}}}{1 - \bar{z} \bar{w}} + \sum_{r \geq 1} \frac{2y}{r + y} (z \bar{w})^r,$$

$$i \tilde{G}_F(\bar{z}, \bar{w}) = -\frac{\sqrt{\bar{z} \bar{w}}}{\bar{z} - \bar{w}} + \frac{\sqrt{z w}}{1 - z w} - \sum_{r \geq 1} \frac{2y}{r + y} (\bar{z} \bar{w})^r.$$  \hspace{1cm} (15)
where the sums are taken over all positive half integers.

When $y = 0$, we obtain the familiar result

$$
\langle \psi(\theta)\psi(\theta') \rangle |_{y=0} + \langle \tilde{\psi}(\theta)\tilde{\psi}(\theta') \rangle |_{y=0}
= GF(e^{i\theta}, e^{i\theta'}) |_{y=0} + \tilde{G}_F(e^{-i\theta}, e^{-i\theta'}) |_{y=0}
= \frac{2}{\sin \frac{\theta - \theta'}{2}}.
$$

(16)

To evaluate the partition function, we need the explicit form of the propagators of $X$ and $\psi(\tilde{\psi})$ at boundary points $z = e^{i\theta}$ and $w = e^{i\theta'}$. That of $X$ was computed in [11],

$$
\langle X(\theta)X(\theta') \rangle \equiv G_B(\theta - \theta') = 2 \sum_{k \in \mathbb{Z}} \frac{1}{|k| + y} e^{ik(\theta - \theta')}.
$$

(17)

The propagator for the fermions on the boundary in the NS sector is

$$
\langle \psi(\theta)\psi(\theta') \rangle + \langle \tilde{\psi}(\theta)\tilde{\psi}(\theta') \rangle \equiv GF(e^{i\theta}, e^{i\theta'}) + \tilde{G}_F(e^{-i\theta}, e^{-i\theta'})
= 2i \sum_{r \in \mathbb{Z} + \frac{1}{2}} \frac{r}{|r| + y} e^{i\theta r}(\theta - \theta').
$$

(18)

which is the expected result as in [23].

Now we can calculate the partition function on the disc. From the definition (2), the partition function on the disc is ($\chi = 1, N = 0$)

$$
Z = \frac{1}{g_s} \int e^{-I_{\text{bulk}} - I_{\text{bndy}}} [dx^\mu][d\psi][d\tilde{\psi}],
$$

(19)

by using the explicit form of $I_{\text{bulk}}$ and $I_{\text{bndy}}$ in (3) and (4), we have

$$
\frac{d}{dy} \ln Z = -\frac{1}{8\pi g_s} \int_0^{2\pi} d\theta \langle X^2 + \psi \frac{1}{\partial \theta} \psi + \tilde{\psi} \frac{1}{\partial \theta} \tilde{\psi} \rangle.
$$

(20)

The right hand side of (20) at the boundary is defined in [23] as

$$
\langle X^2 + \psi \frac{1}{\partial \theta} \psi + \tilde{\psi} \frac{1}{\partial \theta} \tilde{\psi} \rangle \equiv \lim_{\epsilon \to 0} \langle X(\theta)X(\theta + \epsilon) + \psi(\theta) \frac{1}{\partial \theta} \psi(\theta + \epsilon) + \tilde{\psi}(\theta) \frac{1}{\partial \theta} \tilde{\psi}(\theta + \epsilon) \rangle.
$$

(21)
Define
\[ G'_F(\epsilon; y) \equiv \langle \psi(\theta) \frac{1}{\partial \theta} \psi(\theta + \epsilon) \rangle, \]
\[ \tilde{G}'_F(\epsilon; y) \equiv \langle \tilde{\psi}(\theta) \frac{1}{\partial \theta} \tilde{\psi}(\theta + \epsilon) \rangle. \] (22)

We obtain
\[ G'_F(\epsilon; y) + \tilde{G}'_F(\epsilon; y) = -2 \sum_{r \in \mathbb{Z}^+} \frac{1}{|r| + y} e^{ir\epsilon} = G_B(\epsilon; y) - 2G_B(\frac{\epsilon}{2}; 2y). \] (23)

This gives
\[ \langle X^2 + \psi \frac{1}{\partial \theta} \psi + \tilde{\psi} \frac{1}{\partial \theta} \tilde{\psi} \rangle = -8 \ln 2 + 2f(y) - 2f(2y), \] (24)
where
\[ f(y) = \frac{2}{y} - 4y \sum_{k=1}^{\infty} \frac{1}{k(k + y)}. \] (25)

Integrating over \( y \), we find that the partition function is given by
\[ Z(y) = Z'4^y \frac{Z_B(y)^2}{Z_B(2y)}. \] (26)
where \( Z' \) is an integration constant and \( Z_B \) is the partition function for the bosonic case \([11]\) given by,
\[ Z_B(y) = \sqrt{ye^{\gamma y} \Gamma(y)}. \] (27)

According to the conjecture we mentioned in the introduction, for the superstring theory, the BSFT action is simply this partition function. Writing down the action in terms of tachyonic field \( T = uX \), we find the action of the tachyonic field as
\[ S_{\text{tree}} = \frac{T_9}{g_s} \int d^{10}X e^{-\frac{1}{4} T^2} (2 \ln 2 \partial_\mu T \partial^\mu T + 1) \] (28)
where \( T_9 \) is the tension of the non-BPS D9-brane. The integration constant \( Z' \) can then be fixed by comparing to the standard tension of a non-BPS D9-brane [17, 33].

6
3 Tachyon Action at the One Loop level

In this section, we choose the world sheet $\Sigma$ to be an annulus with a rotationally invariant flat metric

$$ds^2 = d\sigma_1^2 + d\sigma_2^2, \quad (29)$$

the complex variable $z = \sigma_1 + i\sigma_2$, with $a \leq |z| \leq b$.

We first take the constant tachyon profile $T(X) = a$, Then the direct calculation gives

$$S_{1-\text{loop}} = V_0 e^{-\frac{1}{2}T^2}. \quad (30)$$

Thus the tachyon potential of the one-loop level is

$$V(T) \sim e^{-\frac{1}{2}T^2}. \quad (31)$$

Next, we take the tachyon profile as in (8),

$$T(X) = \begin{cases} u_a X & \rho = a \\ u_b X & \rho = b \end{cases}, \quad (32)$$

and the boundary term can be written as

$$I' = \frac{y_b}{8\pi} \int_0^{2\pi} d\theta \left( X^2 + \frac{1}{\partial_b} \psi + \bar{\psi} \frac{1}{\partial_b} \bar{\psi} \right)_{\rho=b} + \frac{y_a}{8\pi} \int_0^{2\pi} d\theta \left( X^2 + \frac{1}{\partial_b} \psi + \bar{\psi} \frac{1}{\partial_b} \bar{\psi} \right)_{\rho=a}, \quad (33)$$

where $y_a \equiv u_a^2$ and $y_b \equiv u_b^2$.

The boundary conditions for the Green function of bosonic field are [28, 29]

$$(z\partial + \bar{z}\partial - y_a)G_B(z, w)|_{\rho=a} = 0,$$

$$(z\partial + \bar{z}\partial + y_b)G_B(z, w)|_{\rho=b} = 0. \quad (34)$$

To solve for $G_B(z, w)$, we start with the ansatz,

$$G_B(z, w) = -\ln|z - w|^2 + C_1 \ln|z| |w|^2 + C_2 (\ln|z|^2 + \ln|w|^2) + C_3$$

$$+ \sum_{-\infty}^{\infty} a_k [(z\bar{w})^k + (\bar{z}w)^k] + \sum_{-\infty}^{\infty} b_k \left[ (\frac{z}{w})^k + (\frac{\bar{z}}{\bar{w}})^k \right] \quad (35)$$
As before, one can easily verify that these are solutions of (10) on an annulus. $C_1$, $C_2$, $C_3$, $a_r$’s and $b_r$’s are coefficients to be determined by the boundary conditions (34).

Inserting this ansatz into the boundary conditions (34), expanding it by series and following the procedure as in [28], we get

\[
G_B(z, w) = -\ln(|z - w|^2/b^2) - [ya \ln(|z|^2/b^2) + 2][yb \ln(|w|^2/b^2) - 2 - 2(yb + ya) \ln(|w|^2/b^2) + 2ya \ln(a^2/b^2)
\]

\[
- \sum_{n=1}^{\infty} \ln \left[ \left( \frac{a}{b} \right)^{2n} \frac{z \bar{w}}{a^2 \bar{b}^2} \left| \frac{1 - \left( \frac{a}{b} \right)^{2n} \frac{b^2}{z \bar{w}} \cdot 1 - \left( \frac{a}{b} \right)^{2n} \frac{b^2}{w \bar{z}} \cdot 1 - \left( \frac{a}{b} \right)^{2n} \frac{w}{z} \right|^2 \right]
\]

\[
-2 \sum_{k=1}^{\infty} \frac{[yb(k + ya)b^{2k} + ya(k - yb)a^{2k}]a^{2k}}{k(b^{2k} - a^{2k})[\ln(k + ya)(k + yb)b^{2k} - (k - yb)(k - yb)a^{2k}]}
\]

\[
-2 \sum_{k=1}^{\infty} \frac{[ya(k + yb)b^{2k} - yb(k - ya)a^{2k}]a^{2k}}{k(b^{2k} - a^{2k})[\ln(k + ya)(k + yb)b^{2k} - (k - yb)(k - yb)a^{2k}]}
\]

\[
-2 \sum_{k=1}^{\infty} \frac{[yb + ya]k^2a^{2k}b^{2k}}{k(b^{2k} - a^{2k})[\ln(k + ya)(k + yb)b^{2k} - (k - yb)(k - yb)a^{2k}]}
\]

\[
-2 \sum_{k=1}^{\infty} \frac{[yb + ya]k^2a^{2k}b^{2k}}{k(b^{2k} - a^{2k})[\ln(k + ya)(k + yb)b^{2k} - (k - yb)(k - yb)a^{2k}]}
\]

\[
(1 - iy_a \frac{1}{\partial_\theta}) G_F \bigg|_{\rho = a} = (1 + iy_a \frac{1}{\partial_\theta}) \tilde{G}_F \bigg|_{\rho = a},
\]

\[
(1 + iy_b \frac{1}{\partial_\theta}) G_F \bigg|_{\rho = b} = (1 - iy_b \frac{1}{\partial_\theta}) \tilde{G}_F \bigg|_{\rho = b}.
\]

(38)
To solve the eq. (13) with the boundary conditions (38) on the annulus, we start with the ansatz,

$$iG_F(z, w) = \frac{\sqrt{zw}}{z - w} + \sum_{r \in Z + \frac{1}{4}} a_r (zw)^r + \sum_{r \in Z + \frac{1}{4}} a'_r \left( \frac{z}{w} \right)^r,$$

$$i\tilde{G}_F(\bar{z}, \bar{w}) = -\frac{\sqrt{\bar{z}\bar{w}}}{\bar{z} - \bar{w}} - \sum_{r \in Z + \frac{1}{4}} b_r (\bar{z}\bar{w})^r - \sum_{r \in Z + \frac{1}{4}} b'_r \left( \frac{\bar{z}}{\bar{w}} \right)^r. \quad (39)$$

After a straightforward, albeit lengthy, calculation we get the following results for $G_F(z, w)$ and $\tilde{G}_F(\bar{z}, \bar{w})$

$$iG_F(z, w) = \frac{\sqrt{zw}}{z - w} - \sum_{n=1}^{\infty} \sum_{r \geq \frac{1}{4}} \left[ y_b(r + y_a)b^{2r} + y_a(r - y_b)a^{2r} \right] a^{2r} \left( \frac{z}{w} \right)^r$$

$$-2 \sum_{r \geq \frac{1}{4}} \left[ y_a(r + y_b)b^{2r} + y_b(r - y_a)a^{2r} \right] a^{2r} \left( \frac{b^2}{z\bar{w}} \right)^r$$

$$+2 \sum_{r \geq \frac{1}{4}} \left[ y_b(r + y_a)b^{2r} + y_a(r - y_b)a^{2r} \right] a^{2r} \left( \frac{b^2}{z\bar{w}} \right)^r$$

$$-2 \sum_{r \geq \frac{1}{4}} \left[ y_a(r + y_b)b^{2r} + y_b(r - y_a)a^{2r} \right] a^{2r} \left( \frac{z}{w} \right)^r$$

$$-2 \sum_{r \geq \frac{1}{4}} \left[ y_b(r + y_a)b^{2r} + y_a(r - y_b)a^{2r} \right] a^{2r} \left( \frac{w}{z} \right)^r$$

and

$$i\tilde{G}_F(z, w) = -\frac{\sqrt{zw}}{z - w} + \sum_{n=1}^{\infty} \sum_{r \geq \frac{1}{4}} \left[ y_b(r + y_a)b^{2r} + y_a(r - y_b)a^{2r} \right] a^{2r} \left( \frac{b^2}{z\bar{w}} \right)^r$$

$$-2 \sum_{r \geq \frac{1}{4}} \left[ y_b(r + y_a)b^{2r} + y_a(r - y_b)a^{2r} \right] a^{2r} \left( \frac{z}{w} \right)^r$$

$$-2 \sum_{r \geq \frac{1}{4}} \left[ y_b(r + y_a)b^{2r} + y_a(r - y_b)a^{2r} \right] a^{2r} \left( \frac{w}{z} \right)^r.$$
+ \sum_{n=1}^{\infty} \frac{\sqrt{\left(\frac{a}{b}\right)^{2n} \frac{z}{w}}}{1 - \left(\frac{a}{b}\right)^{2n} \frac{z}{w}} - \sum_{n=1}^{\infty} \frac{\sqrt{\left(\frac{a}{b}\right)^{2n} \frac{b^2}{wz}}}{1 - \left(\frac{a}{b}\right)^{2n} \frac{b^2}{wz}}

-2 \sum_{r \geq \frac{1}{2}} \frac{[y_b(r + y_a)b^{2r} + y_a(r - y_b)\frac{a^2}{b^{2r}}]}{(b^{2r} - a^{2r})[(r + y_a)(r + y_b)b^{2r} - (r - y_a)(r - y_b)a^{2r}]} \left(\frac{\bar{w}}{a^2}\right)^r

+ 2 \sum_{r \geq \frac{1}{2}} \frac{[y_a(r + y_b)b^{2r} + y_b(r - y_a)\frac{a^2}{b^{2r}}]}{(b^{2r} - a^{2r})[(r + y_a)(r + y_b)b^{2r} - (r - y_a)(r - y_b)a^{2r}]} \left(\frac{\bar{w}}{b}\right)^r

-2 \sum_{r \geq \frac{1}{2}} \frac{\frac{(y_b + y_a)ar^{2r}/2}{b^{2r} - a^{2r}}[(r + y_a)(r + y_b)b^{2r} - (r - y_a)(r - y_b)a^{2r}]}{(\bar{w})^r}

+ 2 \sum_{r \geq \frac{1}{2}} \frac{\frac{(y_b + y_a)ar^{2r}/2}{b^{2r} - a^{2r}}[(r + y_a)(r + y_b)b^{2r} - (r - y_a)(r - y_b)a^{2r}]}{(\bar{w})^r} \cdot (41)

To evaluate the partition function, we need to calculate the propagators on the boundary, for both the bosonic and the fermionic fields. We first work at the \(\rho = b\) boundary. Let \(z = be^{i\theta}\) and \(w = be^{i(\theta + \epsilon)}\), we have

\[
\lim_{\epsilon \to 0} G_B(y_a, y_b; \epsilon; a, b) \\
\equiv \lim_{\epsilon \to 0} G_B(be^{i\theta}, be^{i(\theta + \epsilon)})
\]

\[
= -2 \ln(1 - e^{i\epsilon}) - 2 \ln(1 - e^{-i\epsilon}) - 8 \sum_{n=1}^{\infty} \ln \left[1 - \left(\frac{a}{b}\right)^{2n}\right] + \frac{4 - 2y_a \ln \frac{a^2}{b^2}}{2y_b + 2y_a - y_a y_b \ln \frac{a^2}{b^2}}
\]

\[
\quad -4 \sum_{k=1}^{\infty} \frac{[y_b(k + y_a)b^{2k} + y_a(k - y_b)a^{2k}b^{2k}]}{k(b^{2k} - a^{2k})[(k + y_a)(k + y_b)b^{2k} - (k - y_a)(k - y_b)a^{2k}]} - 4 \sum_{k=1}^{\infty} \frac{[y_a(k + y_b)b^{2k} + y_b(k - y_a)a^{2k}b^{2k}]}{k(b^{2k} - a^{2k})[(k + y_a)(k + y_b)b^{2k} - (k - y_a)(k - y_b)a^{2k}]} - 8 \sum_{k=1}^{\infty} \frac{(y_b + y_a)ka^{2k}b^{2k}}{k(b^{2k} - a^{2k})[(k + y_a)(k + y_b)b^{2k} - (k - y_a)(k - y_b)a^{2k}]} \equiv -2 \ln(1 - e^{i\epsilon}) - 2 \ln(1 - e^{-i\epsilon}) + F(y_a, y_b; a, b), \quad (42)
\]

where \(F(y_a, y_b; a, b) \equiv F(y_a, y_b; \epsilon = 0; a, b)\) are the terms which are not singular when we take the limit \(\epsilon = 0\).
As in the case of disc topology, we define
\[ G'_F(\epsilon, y_a, y_b) \equiv G_F'(be^{i\theta}, be^{i(\theta + \epsilon)}) = \langle \psi(\theta) \frac{1}{\partial_{\theta}} \psi(\theta + \epsilon) \rangle, \]
\[ \tilde{G}'_F(\epsilon, y_a, y_b) \equiv \tilde{G}_F'(be^{-i\theta}, be^{-i(\theta + \epsilon)}) = \langle \tilde{\psi}(\theta) \frac{1}{\partial_{\theta}} \tilde{\psi}(\theta + \epsilon) \rangle. \] (43)

We expand \( G'_F(\epsilon, y_a, y_b) \) and \( \tilde{G}'_F(\epsilon, y_a, y_b) \), then compare them to the expansion of \( G_B(y_a, y_b; \epsilon; a, b) \). By a straightforward calculation it can be shown that
\[ G'_F(\epsilon, y_a, y_b) + \tilde{G}'_F(\epsilon, y_a, y_b) = G_B(y_a, y_b; \epsilon; a, b) - 2G_B(2y_a, 2y_b; \frac{\epsilon}{2}; \sqrt{a}, \sqrt{b}). \] (44)

Now, we are ready to calculate the right hand side of (20) at the boundary \( \rho = b \). From eqn. (21), we have
\[ \langle X^2(\theta) + \psi(\theta) \frac{1}{\partial_{\theta}} \psi(\theta) + \tilde{\psi}(\theta) \frac{1}{\partial_{\theta}} \tilde{\psi}(\theta) \rangle|_{\rho=b} = \lim_{\epsilon \to 0} [G_B(y_a, y_b; \epsilon; a, b) + G'_F(\epsilon, y_a, y_b) + \tilde{G}'_F(\epsilon, y_a, y_b)] - 2\lim_{\epsilon \to 0} [G_B(2y_a, 2y_b; \epsilon; \sqrt{a}, \sqrt{b})] = -8 \ln 2 + 2F(y_a, y_b; a, b) - 2F(2y_a, 2y_b; \sqrt{a}, \sqrt{b}). \] (45)

From its definition (2), the partition function on an annulus (\( \chi = 0 \), and \( N = 1 \)) is
\[ Z = \int d\lambda \mu(\lambda) \int e^{-I_{\text{bulk}} - I_{\text{bndy}}} \langle dx^\mu \rangle \langle d\psi \rangle \langle d\tilde{\psi} \rangle \equiv \int d\lambda Z(a/b), \] (46)
where \( \lambda = \ln(a/b) \) is the Teichmüller parameter on the annulus. \( \mu(\lambda) \) includes the contribution of the ghost determinant. In superstring theory, the ghost contribution of the bosonic fields exactly cancels out that of the fermionic fields at the one loop level. So in our case (annulus), \( \mu(\lambda) = 1 \).

Differentiating with respect to the parameter \( y_b \), we have
\[ \frac{\partial}{\partial y_b} \ln Z(a/b) = -\frac{1}{8\pi} \int_0^{2\pi} d\theta \langle X^2(\theta) + \psi(\theta) \frac{1}{\partial_{\theta}} \psi(\theta) + \tilde{\psi}(\theta) \frac{1}{\partial_{\theta}} \tilde{\psi}(\theta) \rangle|_{\rho=b} = (2\ln 2) - \frac{1}{2} F(y_a, y_b; a, b) + \frac{1}{2} F(2y_a, 2y_b; \sqrt{a}, \sqrt{b}). \] (47)
Integrating over $y_b$, up to an integration constant, we get
\[
\ln Z(a/b) = (2 \ln 2) y_b - \frac{1}{2} \ln \left[ y_a + y_b - \frac{y_a y_b}{2} \ln \left( \frac{a^2}{b^2} \right) \right]
+ \sum_{k=1}^{\infty} \left\{ \ln \left[ \left( 1 + \frac{2y_a}{k} \right) \left( 1 + \frac{2y_b}{k} \right) - \left( 1 - \frac{2y_a}{k} \right) \left( 1 - \frac{2y_b}{k} \right) \left( \frac{a}{b} \right)^k \right] - 2 \ln \left[ \left( 1 + \frac{y_a}{k} \right) \left( 1 + \frac{y_b}{k} \right) - \left( 1 - \frac{y_a}{k} \right) \left( 1 - \frac{y_b}{k} \right) \left( \frac{a}{b} \right)^{2k} \right] \right\} + f(y_a),
\]
where $f(y_a)$ is an arbitrary function of $y_a$.

Similarly, we can compute the propagator at the boundary $\rho = a$. Repeating the above procedure, we obtain
\[
\ln Z(a/b) = (2 \ln 2) y_a - \frac{1}{2} \ln \left[ y_a + y_b - \frac{y_a y_b}{2} \ln \left( \frac{a^2}{b^2} \right) \right]
+ \sum_{k=1}^{\infty} \left\{ \ln \left[ \left( 1 + \frac{2y_a}{k} \right) \left( 1 + \frac{2y_b}{k} \right) - \left( 1 - \frac{2y_a}{k} \right) \left( 1 - \frac{2y_b}{k} \right) \left( \frac{a}{b} \right)^k \right] - 2 \ln \left[ \left( 1 + \frac{y_a}{k} \right) \left( 1 + \frac{y_b}{k} \right) - \left( 1 - \frac{y_a}{k} \right) \left( 1 - \frac{y_b}{k} \right) \left( \frac{a}{b} \right)^{2k} \right] \right\} + f(y_b),
\]
where $f(y_b)$ is an arbitrary function of $y_b$.

Comparing the results (48) and (49), we can fix the arbitrary functions $f(y_a)$ and $f(y_b)$. The final expression is
\[
\ln Z(a/b) = (2 \ln 2) y_b + (2 \ln 2) y_a - \frac{1}{2} \ln \left[ y_a + y_b - \frac{y_a y_b}{2} \ln \left( \frac{a^2}{b^2} \right) \right]
+ \sum_{k=1}^{\infty} \left\{ \ln \left[ \left( 1 + \frac{2y_a}{k} \right) \left( 1 + \frac{2y_b}{k} \right) - \left( 1 - \frac{2y_a}{k} \right) \left( 1 - \frac{2y_b}{k} \right) \left( \frac{a}{b} \right)^k \right] - 2 \ln \left[ \left( 1 + \frac{y_a}{k} \right) \left( 1 + \frac{y_b}{k} \right) - \left( 1 - \frac{y_a}{k} \right) \left( 1 - \frac{y_b}{k} \right) \left( \frac{a}{b} \right)^{2k} \right] \right\}. \tag{50}
\]

The partition function, then, can be obtained as
\[
Z(a/b) = Z_B^4 y_b + y_a \frac{Z_B^2(y_a, y_b; a, b)}{Z_B(2y_a, 2y_b; \sqrt{a}, \sqrt{b})}.
\tag{51}
\]
where $Z'$ is the integration constant which we can choose to be the same as the disc case for convenience, and

$$Z_B(y_a, y_b; a, b) = \left[ y_a + y_b - \frac{y_a y_b}{2} \ln \left( \frac{a^2}{b^2} \right) \right]^{-1/2} \cdot \prod_{k=1}^{\infty} \left[ (1 + \frac{y_a}{k})(1 + \frac{y_b}{k}) - (1 - \frac{y_a}{k})(1 - \frac{y_b}{k}) \left( \frac{a}{b} \right)^{2k} \right]^{-1}$$

(52)

is the bosonic partition function on the annulus.

It is convenient to take the outer radius $b = 1$ in the following. Integrating over the modulus $d\lambda = da/a$, we obtain

$$Z = \int d\lambda Z(a)$$

$$= \sqrt{2} Z' 4^{y_a+y_b} \int_0^1 \frac{da}{a} \left[ y_a + y_b - \frac{y_a y_b}{2} \ln a^2 \right]^{-1/2} \cdot \prod_{k=1}^{\infty} \left[ (1 + \frac{2y_a}{k})(1 + \frac{2y_b}{k}) - (1 - \frac{2y_a}{k})(1 - \frac{2y_b}{k})a^k \right]^{2}.$$  (53)

In the one loop level, there are two types of corrections coming from two different configurations:

Case 1. The two ends of the open strings end on two different D-branes on which they may have different tachyonic couplings $y_a$ and $y_b$. In this case, as above, $Z$ is a function of both $y_a$ and $y_b$. To write the background independent action in term of the tachyonic field on one of the two D-branes, say $\rho = b$ boundary, we take $y_a = 0$ and $y_b = y$, then

$$Z = Z' 4^y \int_0^1 \frac{da}{a} \sqrt{\frac{2}{y}} \cdot \prod_{k=1}^{\infty} \left[ \frac{(1 + \frac{2y}{k}) - (1 - \frac{2y}{k})a^{2k}}{(1 + \frac{y}{k}) - (1 - \frac{y}{k})a^{2k}} \right]^{2}.$$  (54)

Expanding (54) in powers of $y$, the expression of $Z$ takes the form

$$Z = Z' \int_0^1 \frac{da}{a} \prod_{k=1}^{\infty} \left( \frac{1 - a^{2k-1}}{1 - a^{2k}} \right) \cdot \left( \sqrt{\frac{2}{y}} + \left[ 2\sqrt{2} \ln 2 - 4\sqrt{2} \sum_{n=1}^{\infty} \ln(1 - a^{2n-1}) \right] \sqrt{y} + \cdots \right).$$  (55)

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Writing the first two terms (corresponding to the kinetic term and the tachyon potential) of the partition function in terms of the tachyonic field, we obtain the tachyon action

\[
S_{1-\text{loop}} = T_9 \int_0^1 \frac{da}{a} \prod_{k=1}^{\infty} \left( \frac{1 - a^{2k-1}}{1 - a^{2k}} \right) 
\]

\[
\int d^{10} X e^{-\frac{4}{T}T^2} \left( 2 \ln 2 - 4 \sum_{n=0}^{\infty} \ln(1 - a^{2n-1}) \right) \partial_{\mu} T \partial^{\mu} T + 1 \right). \quad (56)
\]

Let us consider the integral over the modulus \( a \). The integral

\[
\int_0^1 \frac{da}{a} \prod_{k=1}^{\infty} \left( \frac{1 - a^{2k-1}}{1 - a^{2k}} \right)
\]

is divergent at \( a = 0 \). Introducing a cut-off parameter \( \Lambda \), we may write this integral as

\[
\int_0^1 \frac{da}{a} \prod_{k=1}^{\infty} \left( \frac{1 - a^{2k-1}}{1 - a^{2k}} \right) = \Lambda + \Lambda_{\text{finite}}, \quad (57)
\]

where \( \Lambda_{\text{finite}} \) is the finite part of the integral. The cut off \( \Lambda \) will be absorbed by renormalization.

In the coefficient of the kinetic energy term, the integrand

\[
-4 \int_0^1 \frac{da}{a} \prod_{k=1}^{\infty} \left( \frac{1 - a^{2k-1}}{1 - a^{2k}} \right) \sum_{n=1}^{\infty} \ln(1 - a^{2n-1}) \quad (58)
\]

is finite, so we can ignore this term comparing to the first term of the kinetic energy term. Thus the one-loop action can be written as

\[
S_{1-\text{loop}} = T_9 (\Lambda + \Lambda_{\text{finite}}) \int d^{10} X e^{-\frac{4}{T}T^2} \left( (2 \ln 2) \partial_{\mu} T \partial^{\mu} T + 1 \right). \quad (59)
\]

**Case 2.** The two ends of open strings end on different D-branes, but with the same coupling \( y_a = y_b = y \). In this case, the partition function \( Z(y_a, y_b) |_{y_a = y_b = y} \) may be used, since the ends are independent. The partition function is given by

\[
Z = Z' 4^y \int_0^1 \frac{da}{a} \sqrt{\frac{1}{y - (y^2/4) \ln a^2}} \prod_{k=1}^{\infty} \frac{(1 + \frac{2y}{k})^2 - (1 - \frac{2y}{k})^2 a^k}{(1 + \frac{2y}{k})^2 - (1 - \frac{2y}{k})^2 a^{2k}} \quad (60)
\]

\[\text{Note the factor before the integration is 4^y, not 4^{2y} as it might appear to do so from (52). This is because when we set } y_a = y_b = y, \text{ we do not need to add the extra terms } f(y_a) \text{ and } f(y_b) \text{ in (46) and (47).} \]
Again expanding in $y$, we obtain

$$ Z = Z' \int_0^1 \frac{da}{a} \frac{1}{\sqrt{1-(y/2) \ln a}} \cdot \prod_{k=1}^{\infty} \left( \frac{1-a^{2k-1}}{1-a^{2k}} \right) 
\cdot \left( \frac{1}{\sqrt{y}} + \left[ 2 \ln 2 - 8 \sum_{n=1}^{\infty} \ln(1-a^{2n-1}) \right] \sqrt{y} + \cdots \right). \quad (61) $$

Note that we keep the square root form of the zero mode without expanding it in $y$, this is because that it is not proper to expand it for small $y$ at the integral end $a = 0$.

The integral in (61) has a divergence at $a = 0$. (Similarly, here we can ignore the second term, which is finite, of the kinetic energy term.) This is handled by replaces the lower limit $a = 0$ by $a = \delta$ and lets $\delta \to 0$ at the end. Consider the integral

$$ I = \lim_{\delta \to 0} \int_{\delta}^{1} \frac{da}{a} \frac{1}{\sqrt{1-(y/2) \ln(a)}} \prod_{k=1}^{\infty} \left( \frac{1-a^{2k-1}}{1-a^{2k}} \right) 
\cdot \left( 1 + 2 \ln 2 - 8 \sum_{n=1}^{\infty} \ln(1-a^{2n-1}) \right) \sqrt{y} + \cdots, \quad (62) $$

where $C_n$ are coefficients whose form is not needed in what follows, but $\sum_{n=1}^{\infty} C_n a^n$ is convergent.

This integral can be done and the result is

$$ I = \lim_{\delta \to 0} - \frac{4}{y} \left( 1 - \sqrt{1 - \frac{y}{2} \ln \delta} \right) + \text{finite terms} \sim \frac{\ln \delta}{\sqrt{y}}, \quad (63) $$

therefore,

$$ Z \sim \left( \frac{1}{y} + 2 \ln 2 \right). \quad (64) $$

we obtain the tachyon action in the form

$$ S_{1-\text{loop}} \sim \int d^{10}X \ e^{-\frac{1}{2} R_{\mu T} T^{\mu T} [[2 \ln 2] \partial T \partial^{\mu T} + 1]} \quad (65) $$
Case 3. The form of tachyon action in case 2 is unexpected according to our early analysis (31). The reason is that we set $y_a = y_b$ after we integrated over $y_b$. In fact, this is the case which corresponds to an open string ending on two different branes. Another choice is to set $y_a = y_b = y$ before integrating over $y_b$. Since the ends are not independent in this case, $Z$ is only a function of $y$, and thus this corresponds to an open string ending on only a single brane. This should be the proper correction to the disc case. To do this, let’s go back to equation (47) and set $y_a = y_b = y$ there; we obtain

$$\frac{\partial}{\partial y} \ln Z(a/b) = -\frac{1}{8\pi} \int_0^{2\pi} d\theta \left( X^2(\theta) + \psi(\theta) \frac{1}{\partial_\theta} \psi(\theta) + \tilde{\psi}(\theta) \frac{1}{\partial_{\tilde{\theta}}} \tilde{\psi}(\theta) \right) \bigg|_{\rho=b}$$

$$= (2 \ln 2) - \frac{1}{2} F(y, y; a, b) + \frac{1}{2} F(2y, 2y; \sqrt{a}, \sqrt{b}). \quad (66)$$

Integrating over $y$, up to an integration constant, we get

$$\ln Z(a/b) = (2 \ln 2) y - \frac{1}{4} \ln \left[ 2y - y^2 \ln \left( \frac{a}{b} \right) \right]$$

$$+ \frac{1}{2} \sum_{k=1}^{\infty} \left\{ \ln \left[ \left( 1 + \frac{2y}{k} \right)^2 - \left( 1 - \frac{2y}{k} \right)^2 \left( \frac{a}{b} \right)^k \right] \right. $$

$$- 2 \ln \left[ \left( 1 + \frac{y}{k} \right)^2 - \left( 1 - \frac{y}{k} \right)^2 \left( \frac{a}{b} \right)^{2k} \right] \right\}, \quad (67)$$

Taking the outer radius $b = 1$ in the following. Integrating over the modulus $d\lambda = da/a$, we obtain

$$Z = \int d\lambda Z(a)$$

$$= \sqrt{2} Z' 4^y \int_0^1 \frac{da}{a} \left( 2y - y^2 \ln a \right)^{-1/4} \cdot \prod_{k=1}^{\infty} \left[ (1 + \frac{2y}{k})^2 - (1 - \frac{2y}{k})^2 (\frac{a}{b})^k \right]^{1/2}$$

$$= \sqrt{2} Z' \int_0^1 \frac{da}{a} \left[ 2y - y^2 \ln a \right]^{-1/4} \cdot \prod_{k=1}^{\infty} \left( \frac{1 - a^{2k-1}}{1 - a^{2k}} \right)^{1/2}$$

$$\cdot \left( 1 + \left[ 2 \ln 2 - 4 \sum_{n=1}^{\infty} \ln(1 - a^{2n-1}) \right] y + \cdots \right). \quad (68)$$
We handle the integral over $a$ by the same way as in case 2. We get

$$I = \lim_{\delta \to 0} \int_1^\delta \frac{da}{a} \left[2y - y^2 \ln a\right]^{-1/4} \cdot \prod_{k=1}^\infty \left(1 - \frac{a^{2k-1}}{1 - a^{2k}}\right)^{1/2}$$

$$= \lim_{\delta \to 0} -\frac{4}{3y^2} \left((2y)^{3/4} - (2y - y^2) \ln \delta)^{3/4}\right) + \text{finite terms}$$

$$\sim \frac{\ln \delta}{\sqrt{y}}$$  \hspace{1cm} (69)$$

Retaining the first two terms of the partition function,

$$Z \sim Z' \left(\frac{1}{\sqrt{y}} + 2 \ln 2 \sqrt{y}\right) \cdot \ln \delta \equiv (\Lambda + \Lambda_{\text{finite}})Z' \left(\frac{1}{\sqrt{y}} + 2 \ln 2 \sqrt{y}\right).$$  \hspace{1cm} (70)$$

we obtain the tachyon action in the form

$$S_{1\text{-loop}} = T_9(\Lambda + \Lambda_{\text{finite}}) \int d^{10}X e^{-\frac{1}{2}T^2}[(2 \ln 2)\partial_\mu T \partial^\mu T + 1].$$  \hspace{1cm} (71)$$

When we consider the D9-brane, which is actually the whole 10-dimensional spacetime, both ends of open strings must end on the same D9-brane. Thus we should use the one loop tachyon action of case 3 as the quantum correction of the tree level action. By the definition of the effective action of tachyonic fields (2), we get

$$S = S_{\text{tree}} + S_{1\text{-loop}}$$

$$= \frac{T_9}{g_s} \int d^{10}X e^{-\frac{1}{2}T^2}[(2 \ln 2)\partial_\mu T \partial^\mu T + 1]$$

$$+ T_9(\Lambda + \Lambda_{\text{finite}}) \int d^{10}X e^{-\frac{1}{2}T^2}[(2 \ln 2)\partial_\mu T \partial^\mu T + 1]$$

$$= T_9 \int d^{10}X e^{-\frac{1}{2}T^2} \left[\left(\frac{1}{\lambda} + \Lambda + \Lambda_{\text{finite}}\right) [(2 \ln 2)\partial_\mu T \partial^\mu T + 1]\right],$$  \hspace{1cm} (72)$$

where we define \[13\]

$$\lambda = g_s e^{-\frac{1}{2}T^2}$$  \hspace{1cm} (73)$$

as the effective string coupling. The infinite cut-off $\Lambda$ can be absorbed into the renormalized effective string coupling $\lambda'$ defined by

$$\frac{1}{\lambda'} = \frac{1}{\lambda} + \Lambda.$$  \hspace{1cm} (74)
Thus we can write the tachyon action up to the one loop quantum correction as

\[ S = \frac{T_9'}{\lambda'} \int d^{10}X \ e^{-\frac{1}{2} T^2} \left[ (2 \ln 2) \partial_\mu T \partial^\mu T + 1 \right], \]  

(75)

where the one loop corrected D9-brane tension is (in terms of \( \alpha' = 2 \))

\[ T_9' = (1 + \lambda')T_9 \]  

(76)

and the one loop quantum corrected tachyon potential

\[ V(T) = e^{-\frac{1}{2} T^2} \]  

(77)

which is the square of the tree level tachyon potential \( e^{-\frac{1}{4} T^2} \).

4 Conclusions

We have studied quantum effects for tachyon condensation using background independent string field theory methods. In particular, the contribution to open superstring partition function from the annulus diagram is calculated. We have shown that there are three important contributions to the tachyonic field action at one loop level. First, the tachyonic potential is proportional to \( e^{-\frac{1}{2} T^2} \) and not \( e^{-\frac{1}{4} T^2} \) as for the tree level. The second point is that there arises a renormalization of the effective string coupling given by \( \lambda' = g_{\text{ren}}^s e^{-T^2/4} \), where \( g_{\text{ren}}^s \) is the renormalized closed string coupling constant defined by (73) and (74). We note that this involves a field dependent renormalization. Finally, there is a finite one-loop contribution to the D9-brane tension given by (76).

Except for the renormalized quantities \( \lambda', T_9', \) and \( e^{-T^2/2} \), the one loop corrected tachyon action has the same form as the tree level one, which is expected based on general analysis of loop effects. Our assumption that the string field action is given by the partition function on the world sheet with the topology of an annulus is therefore justifiable.

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