I. INTRODUCTION

Study of the open string tachyon has led to some understanding of the nonperturbative aspects of string theory. Among many results related to the open string tachyon, rolling tachyons and tachyon solitons have been important issues.

When an unstable D-brane decays homogeneously, the real-time decay process of the D-brane can be described by a marginally-deformed boundary conformal field theory (BCFT). A. Sen found one parameter family of time-dependent solutions, referred to as rolling tachyon, which is identical with the marginal deformation parameter. This has been studied in Refs. [2, 3, 4] for the pure tachyon case and in Refs. [5, 6] in the presence of background electromagnetic fields. These time evolution properties of the unstable D-brane are also described in the other effective field theories, e.g., boundary string field theory (BSFT) [7, 8, 9], Dirac-Born-Infeld (DBI)-type effective field theory (EFT) [10, 11], and noncommutative effective field theory (NCFT) [12].

Simultaneously, much work has been done on tachyon solitons, for instance, tachyon kinks [13, 14, 15, 16, 17, 18, 19], tube [20], and vortices [13, 17, 21]. For the case of the pure tachyon, only the array of kink-antikink is found [14, 15, 16]. In the presence of constant electromagnetic fields, $F_{\mu\nu}$, rich spectra of the tachyon kinks are found and those configurations are classified according to the sign of $C^{11}$, the (11)-component of the cofactor for matrix $(\eta + F_{\mu\nu})$, where $x^1$ is the transverse direction to the kink [15]. When $C^{11}$ is negative, there are two species of kinks which are the array of kink-antikink and the single topological BPS kink. These two kinks are also identified in BCFT, NCFT, and BSFT [19]. (In addition to the above two, there is another solution in DBI-type EFT [15] and NCFT [18].) On the other hand, for $p \geq 2$, $C^{11}$ can be positive when there are nonzero components of electric field in both longitudinal and transverse directions as well as nonzero magnetic field. In this case, there are three more species of kinks named bounce, half-kink, and topological non-BPS kink in BSFT [19], DBI-type EFT [15], and NCFT [18]. These solutions have not been investigated in BCFT.

The purpose of this paper is to study the latter three types of solutions in the context of BCFT. We construct exactly marginal boundary tachyon operators in the presence of constant electromagnetic fields. Functional forms of these operators are hyperbolic cosine, hyperbolic sine, and exponential types. We compute the spacetime energy-momentum tensor and the source for the antisymmetric tensor field are computed in the path integral approach for the exponential-type tachyon vertex operator.

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II. TACHYON VERTICES FOR STATIC KINKS

Let us consider an unstable Dp-brane of tension $T_p$ ($p \leq 25$) with fundamental strings in bosonic string theory. Dynamics of the unstable Dp-brane is described in BCFT by introducing interaction terms of the tachyon $T(X)$ and the gauge field $A_{\mu}(X)$ on the boundary of the worldsheet. Specifically, in the flat spacetime, the worldsheet action...
is given by
\[ S_{BCFT} = S_0 + S_A + S_T \quad (1) \]
with
\[ S_0 = \frac{1}{2\pi} \int_{\Sigma} d^2w \eta_{\mu\nu} \partial X^\mu \partial X^\nu, \]
\[ S_A = \frac{i}{2\pi} \int_{\partial\Sigma} d\tau A_\mu(X) \partial_\tau X^\mu, \]
\[ S_T = \int_{\partial\Sigma} d\tau T(X), \quad (2) \]
where \( \Sigma \) denotes a worldsheet and \( \tau \) is the coordinate along the worldsheet boundary \( \partial\Sigma \). We also have the deformed boundary condition for \( X^\mu(w, \bar{w}) \) in the presence of the background electromagnetic field,
\[ (\partial - \bar{\partial})X^\mu \big|_{w=\bar{w}} + F^\mu_\lambda (\partial + \bar{\partial})X^\lambda \big|_{w=\bar{w}} = 0, \quad (4) \]
where \( w = \tau + i\sigma \) \((-\infty < \tau < \infty, \ 0 \leq \sigma < \infty) \) is the complex coordinate on the upper half plane (UHP). From the action (2) we read the worldsheet energy-momentum tensor:
\[ T_{ww} = -\eta_{\mu\nu} \partial X^\mu(w) \partial X^\nu(w) : . \quad (5) \]

The main purpose of this section is to find conformally-invariant tachyon vertices, \( S_T \), in the presence of a constant electromagnetic field \( F_{\mu\nu} \), which are appropriate for tachyon kinks. For convenience, we take symmetric \( S \)- and anti-symmetric \( A \)-gauge,
\[ A_{\mu} = -\frac{1}{2} F_{\mu\nu} X^\nu. \]

The general exponential-type tachyon vertex operator in open string theory is written as
\[ T(X) = \sum_i \left( \lambda_+^i e^{ik^\mu_+ X^\mu} + \lambda_-^i e^{-ik^\mu_- X^\mu} \right), \quad (6) \]
where \( \lambda_+^i \) and \( k^\mu_+ \) should be chosen to make \( T(X) \) real. Under the deformed boundary condition in Eq. (4), the correlation function on the upper half plane is obtained through operator product expansion (22):
\[ \langle X^\mu(w)X^\nu(w') \rangle_{\text{UHP}} = -\eta^{\mu\nu} \ln |w - w'| + \eta^{\mu\nu} \ln |w - \bar{w}'| - G^{\mu\nu} \ln |w - \bar{w}'|^2 - \theta^{\mu\nu} \ln \left( \frac{w - \bar{w}'}{w - \bar{w}} \right) \quad (7) \]
with open string metric \( G^{\mu\nu} \) and noncommutative parameter \( \theta^{\mu\nu} \) defined by
\[ G^{\mu\nu} = \left( \frac{1}{\eta + F} \right)^{\mu\nu}_S, \quad \theta^{\mu\nu} = \left( \frac{1}{\eta + F} \right)^{\mu\nu}_A, \quad (8) \]
where ‘S’(‘A’) denotes the symmetric (anti-symmetric) part of the matrix component. Then Eq. (1) on the boundary of coordinate \( \tau' \) provides
\[ X(w)X(\tau') \sim -G^{\mu\nu} \ln |w - \tau'|^2 - \theta^{\mu\nu} \ln \left( \frac{w - \tau'}{w - \tau} \right). \quad (9) \]

The conformal weight of the tachyon vertex operator at the boundary may be identified by considering the operator product expansion with the worldsheet energy momentum-tensor. By using Eq. (9) and the deformed boundary condition (1),
\[ T_{ww}(w) : e^{ik \cdot X(\tau')} : \sim \frac{G^{\mu\nu} k^\mu k^\nu}{(w - \tau')^2} : e^{ik \cdot X(\tau')} : + \frac{1}{w - \tau} \partial_{\tau'}: e^{ik \cdot X(\tau')} :. \quad (10) \]

Then the boundary operator has conformal weight \( h = G^{\mu\nu} k^\mu k^\nu \) and becomes marginal when
\[ G^{\mu\nu} k^\mu k^\nu = 1. \quad (11) \]

For the rest of the paper we will consider only the operators which depend on a single spatial coordinate, say \( X^1 \), in bosonic string theory. In this case, Eq. (11) reduces to
\[ k_1^2 = \frac{1}{G^{11}} = \frac{Y}{C^{11}}, \quad (12) \]
where \( Y \equiv \det(\eta + F) \) and \( C^{\mu\nu} \) is the cofactor of the matrix \( (\eta + F)_{\mu\nu} \). For physical electromagnetic fields, \( Y \) is nonpositive definite, but \( C^{11} \) can have both negative and positive values for \( p \geq 2 \). Thus depending on the signature of \( C^{11} \), \( k_1 \) can be real or imaginary. This is in contrast with the case of rolling tachyons, in which \( k_0^2 = Y/C^{00} \) is always negative, since \( Y \) is negative and \( C^{00} \) is kept positive.

When \( C^{11} \) is negative \((k_1 \text{ is real})\), the exactly marginal tachyon vertex operator has the form, up to a translation in \( X^1 \),
\[ T(X) = \lambda \cos \left( \sqrt{\frac{Y}{1 - C^{11}}} X^1 \right), \quad (13) \]
where we set \( \lambda_+ = \lambda_- = \frac{1}{\sqrt{2}} \lambda \) and \( k^1_+ = k^1_- = k_1 \) in Eq. (6). The resulting configuration for the pure tachyon case \( (Y = C^{11} = -1) \) is interpreted as an array of D(p−1)D(p−1) [23].

Eq. (13) was also discussed in the presence of an electric field in superstring theory \( (Y = -1 + E^2, C^{11} = -1) \) [12].

When \( C^{11} \) is positive \((k_1 \text{ is pure imaginary})\), the reality condition for the tachyon vertex operators in Eq. (9) requires that both \( \lambda_+ \) and \( \lambda_- \) should be real. According to the boundary values of \( T \), the tachyon vertices are classified by three types: (i) \( T(-\infty) = \mp \infty \) and \( T(+\infty) = \pm \infty \); (ii) \( T(-\infty) = 0 \) and \( T(+\infty) = \pm \infty \); (iii) \( T(-\infty) = \pm \infty \) and \( T(+\infty) = \pm \infty \):
\[ T(X^1) = \begin{cases} (i) \lambda \sinh (\kappa X^1) \\ (ii) \lambda \exp (\pm \kappa X^1) \\ (iii) \lambda \cosh (\kappa X^1) \end{cases}, \quad (14) \]
where
\[ ik_1 = \sqrt{\frac{Y}{C^{11}}} \equiv \kappa. \quad (15) \]
Now we check the exact marginality of $T(X^1)$. The operator product expansion of $T(X^1)$ with itself contains only pole singularity at most,

$$T(X^1(\tau_1))T(X^1(\tau_2)) = \frac{\tilde{K}}{\epsilon_{\tau_1-\tau_2}^2} + \text{(regular terms)},$$

where $\tilde{K} = -\frac{1}{2} \lambda^2$, 0, and $\frac{1}{7} \lambda^2$ for (i), (ii), and (iii) in Eq. (14), respectively. (Note that in this calculation the noncommutative parameter $\theta_{\mu \nu}^2$ plays no role, since $T(X)$ consists only of a single field $X^1$.) This implies that the operator $T(X)$ is self-local [25], which guarantees that it is exactly marginal.

### III. TACHYON KINKS WITH ELECTROMAGNETIC FIELD

In the $\sigma$-model approach to the string theory, the partition function of the worldsheet action with exact marginal couplings is identified as the spacetime action, and the couplings are interpreted as spacetime fields [26]. In relation to rolling tachyons, the energy-momentum tensor in BCFT can be read from the partition function of the worldsheet action [17] without gauge field interaction [4]. Here we take into account the gauge field interaction on the worldsheet boundary induced by the constant electromagnetic fields in addition to the tachyon vertex term. The energy-momentum tensor in BCFT can be read from the partition function of worldsheet theory coupled to background gravity [4]:

$$S = Z_{\text{disk}} \sim \int [dX^\mu] e^{-S_{\text{BCFT}}},$$

where $Z$ is the disk partition function, and we have replaced $\eta_{\mu \nu}$ in $S_{\text{BCFT}}$ with the generic curved spacetime metric $\eta_{\mu \nu}$. From the spacetime action [17] we read the energy-momentum tensor in flat spacetime:

$$T_{\mu \nu} = -\frac{2}{\sqrt{-g}} \frac{\delta S}{\delta g_{\mu \nu}} |_{\eta_{\mu \nu} = \eta_{\mu \nu}} = K \{ \eta_{\mu \nu} B(x) + A_{\mu \nu}(x) \},$$

where

$$B(x) = \langle [dX^\mu] e^{-[S_0(X^\mu) + S_A(X^\mu) + S_T(x^\mu + \theta x^\nu)]} \rangle_A,$$

$$A_{\mu \nu}(x) = \int [dX^\mu] \left( \langle : \partial X^\mu(0) \tilde{\partial} X^\mu(0) : \rangle + \langle : \partial X^\mu(0) \tilde{\partial} X^\mu(0) : \rangle \right)$$

$$\times e^{-[S_A(X^\mu) + S_T(x^\mu + X^\nu)]} \right) \times e^{-[S_A(X^\mu) + S_T(x^\mu + X^\nu)]} \right) \times e^{-[S_A(X^\mu) + S_T(x^\mu + X^\nu)]} \right),$$

for (i), (ii), and (iii) in Eq. (18), respectively. (Note that in this calculation the noncommutative parameter $\theta_{\mu \nu}^2$ plays no role, since $T(X)$ consists only of a single field $X^1$.) This implies that the operator $T(X)$ is self-local [25], which guarantees that it is exactly marginal.

The notations used here are as follows. $K$ in Eq. (18) is an overall constant. The symbol $: :$ denotes the normal ordering defined by

$$: \partial X^\mu(z) \tilde{\partial} X^\nu(z') : = \partial X^\mu(z) \tilde{\partial} X^\nu(z') + \eta_{\mu \nu} \tilde{\partial} \tilde{\partial} \log |z - z'|.$$

$X^\mu$ is split into the center of mass coordinate $x^\mu$ and fluctuations $X^\mu$, i.e., $X^\mu = x^\mu + \theta x^\nu$. Finally, $\langle \cdot \cdot \cdot \rangle_A$ denotes the vacuum expectation value in the presence of the gauge field. Proper normalization is made by choosing the vacuum-to-vacuum expectation value for the unit operator in the presence of the constant electromagnetic fields as

$$\langle 1 \rangle_A = \sqrt{-\det(\eta_{\mu \nu} + F_{\mu \nu})}.$$

Now we assume the case where $p \geq 2$ and $C_{11}^{11} > 0$. Then, the operators in Eq. (14) are exactly marginal and we can consider the deformation by them. Here we will consider only the exponential type,

$$T(X^1) = \lambda e^{\kappa X^1},$$

and compute the spacetime energy-momentum tensor in Eq. (18). After some calculations following the method in [4], we find

$$B(x) = \sqrt{-Y} f(x),$$

$$A_{\mu \nu}(x) = 2 \sqrt{-Y} \left( G^{\mu \nu} - \frac{1}{2} \eta^{\mu \nu} \right) f(x) + \kappa^2 \left( G^{11} G^{11} - \theta^{11} \theta^{11} \right) f(x) - \kappa^2 \left( G^{11} G^{11} - \theta^{11} \theta^{11} \right),$$

where

$$f(x) = \frac{1}{1 + 2 \pi \lambda e^{\kappa x_+}}.$$

Combining Eqs. (24), (25), and (18), we finally obtain

$$T^{\mu \nu} = T_p \sqrt{-Y} \kappa^2 \left[ \frac{1}{\kappa^2} \left( G^{\mu \nu} + G^{11} G^{11} - \theta^{11} \theta^{11} \right) f + G^{11} G^{11} - \theta^{11} \theta^{11} \right],$$

where we have determined the overall normalization constant, $K = -\frac{1}{2} T_p$, by comparison with the static limit,

$$T^{00}(\lambda = 0, F_{\mu \nu} = 0) = T_p.$$

Note that the energy-momentum tensor (27) satisfies the conservation law,

$$\partial_\mu T^{\mu \nu} = 0.$$
In addition, we compute the source for the antisymmetric tensor field from the spacetime action \(\Pi^{\mu
u}\) as
\[
\Pi^{\mu
u} \equiv 2 \frac{\delta S}{\delta F_{\mu
u}} = \frac{T_p}{2} \left\langle \left( \partial X'_\mu(0) \partial X'_\nu(0) : - \partial X'_\mu(0) \partial X'_\nu(0) : \right) \right\rangle e^{-\lambda f d\tau e^{\kappa(x^1+x')}},
\]
\[
= T_p \sqrt{-Y} k^2 \left\{ \frac{1}{k^2} \theta^{\mu
u} - \left( G^{\mu\nu} \theta^{\rho\sigma} - G^{\nu\rho} \theta^{\mu\sigma} \right) \right\} f + G^{\mu\nu} \theta^{\rho\sigma} - G^{\nu\rho} \theta^{\mu\sigma},
\]
\[
= T_p \sqrt{-Y} k^2 \left\{ \left[ \frac{1}{k^2} - \left( G^{\mu\nu} \theta^{\rho\sigma} - G^{\nu\rho} \theta^{\mu\sigma} \right) \right] f + G^{\mu\nu} \theta^{\rho\sigma} - G^{\nu\rho} \theta^{\mu\sigma} \right\}.
\]

From Eqs. (27) and (30), we see that both \(T^{\mu\nu}\) and \(\Pi^{\mu\nu}\) have essentially the same \(x^1\)-dependent part given by Eq. (20). For some components, the \(x^1\)-dependent part vanishes. Let us consider the case \(p = 2\) for definiteness. Then, \(-Y = 1 - E_2^2 - E_2^2 + B^2 (> 0)\) and \(C^{11} = 1 + E_2^2 (> 0)\) and, we find constant components
\[
\frac{\Pi^{01}}{E_1} = \frac{\Pi^{12}}{B} = \frac{T^{01}}{E_2 B} = \frac{T^{11}}{1 - E_2^2} = -\frac{T^{02}}{E_1 B} = -\frac{T^{12}}{E_1 E_2},
\]
\[
= \frac{T_2}{\sqrt{-Y}},
\]

while the other components are given by the sums of an \(x^1\)-dependent piece and a constant,
\[
\Pi^{02} = \frac{T_2 \sqrt{-Y} E_2}{1 - E_2^2} f + \frac{E_2 (E_1^2 - B^2)}{E_1 (1 - E_2^2)} \Pi^{01},
\]
\[
T^{00} = \frac{T_2 \sqrt{-Y}}{1 - E_2^2} f + \frac{E_2^2 + E_2^2 B^2}{E_1 (1 - E_2^2)} \Pi^{01},
\]
\[
T^{22} = \frac{T_2 \sqrt{-Y}}{1 - E_1^2} f - \frac{E_2^2 - B^2}{E_1 (1 - E_1^2)} \Pi^{01}.
\]

Note that for positive \(\lambda\), \(f(x^1)\) is regular everywhere, while it has a singular point if \(\lambda\) is negative. This is a well-known feature of bosonic theory and implies an instability of bosonic theory \([2]\) in which the effective tachyon potential is unbounded from below in the negative tachyon direction. For positive \(\lambda\), the static tachyon configuration connects the perturbative vacuum string at \(T = 0\) and the vacuum at \(T = \infty\), where the unstable D-brane disappears. Therefore, it may be interpreted as a half brane \(1/2(D(p-1))\)-brane similar to a bubble wall at a given time connecting two phases. This configuration has already been obtained as a half kink in DBI-type EFT \([15]\), NCFT \([18]\), and BSFT \([19]\). Since these effective actions can be derived from the superstring theory, we expect that the half brane can also be obtained in super BCFT.

In addition to the exponential-type vertex, there exist two more boundary tachyon operators, hyperbolic sine and cosine in Eq. (14). They could also be studied in the context of BCFT.

IV. CONCLUDING REMARKS

We have investigated tachyon vertex operators which are exactly marginal and depend on only one spatial direction \(X^1\) in the presence of background electromagnetic fields in bosonic string theory.

We found that, when \(\text{det}(\eta + F) < 0\) and \(C^{11} > 0\), there are hyperbolic sine, hyperbolic cosine, and exponential-type exactly marginal operators. This case can be realized depending on the magnitude of electric and magnetic fields along the transverse direction of \(x^1\). Therefore, the dimension of the unstable D-brane should be equal to or larger than two.

We obtained the general form of the energy-momentum tensor for the exponential-type operator by using the path integral approach developed in Ref. \([4]\). It may be interpreted as a half brane, which has been found in other approaches such as DBI-type EFT, NCFT, and BSFT.

Our analysis may be extended to the superstring case and also to other marginal operators, i.e., the hyperbolic sine and hyperbolic cosine cases. It may also be intriguing to include the R-R coupling \([27]\).

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