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Brane decay and death of open strings

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Abstract: We show how open strings cease to propagate when unstable D-branes decay. The information on the propagation is encoded in BSFT two-point functions for arbitrary profiles of open string excitations. We evaluate them in tachyon condensation backgrounds corresponding to (i) static spatial tachyon kink (= lower dimensional BPS D-brane) and (ii) homogeneous rolling tachyon. For (i) the propagation is restricted to the directions along the tachyon kink, while for (ii) all the open string excitations cease to propagate at late time and are subject to a collapsed light cone characterized by Carrollian contraction of Lorentz group.

Keywords: D-branes, Tachyon Condensation, String Field Theory.

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Contents

1. Open strings in tachyon condensation

What happens to open strings when D-branes decay by tachyon condensation? Although much interesting structures of off-shell string theory has been revealed since Sen's conjectures [1] on disappearance of unstable D-branes by open string tachyon condensation, the above question which is indispensable for truly understanding the decay of D-branes is still unanswered in a satisfactory manner. At vanishing string coupling, the D-branes are by themselves defined as hypersurfaces on which open strings end, thus the question looks tautological and ill-defined — however, this should be the place where intriguing physics is hidden behind. In spite of the fact that the degrees of freedom of dynamical D-branes are defined through the open strings, the theory should also describe physics without D-branes. In this article, we provide a description of open strings in brane decay backgrounds.

The tachyon condensation is involved intrinsically with off-shell physics in string theory such as its vacuum structure, one has to employ string field theories to make an advance to answer the above question. The string field theories known so far consist mainly of two categories — boundary string field theory (BSFT) [2] and cubic string field theory (CSFT) [3]. Depending on which one uses, the approaches to the above question might be different. Here we take the first one since it enables us to observe a direct relation to worldsheet properties, as we shall see in this article. In the approaches using the BSFT, the essential point was provided in [4],[1] in which the Sen's conjecture on the disappearance of unstable D-branes by the tachyon condensation [1] is realized in such a way that the D-brane worldvolume action (BSFT action) itself vanishes. This is a general property of BSFT actions concerning the tachyon potential in string theory. When a constant tachyon (which is off-shell) is put to its vacuum $|T| = \infty$, the BSFT action vanishes due to the overall potential factor going like $\sim e^{-T^2}$ or $\sim e^{-T\bar{T}}$ in superstring theory.

[1] See also [5]–[8]. The observation in [4] is purely classical, while the confinement mechanism proposed earlier by Yi [7] is a quantum effect.

The achievements in the BSFT is not only for the constant tachyon. Tachyon profiles linear in the target space coordinates provide intriguing results, including a verification of the Sen's conjecture on the D-brane descent relations [1], for example. For a non-BPS D9-brane in type-IIA superstring theory, a linear tachyon profile

$$T = u_9 X^9 \tag{1.1}$$

with $u_9 = \pm\infty$ solves the equation of motion of the BSFT. This kink solution connecting the two vacua $T = \pm\infty$ represents a BPS D8-brane localized in the x^9 direction and in fact its tension is precisely reproduced in the analysis of the super BSFT [9] (for the bosonic case and the D-$\overline{\mathrm{D}}$ case, see [5] and [10] respectively).

Recent development on the tachyon condensation is mostly on its time-dependent process whose study was initiated in [11] with use of boundary deformations of conformal field theories, of the type $T(X) \sim e^{X^0}$, which is called rolling tachyon. An interesting outcome was that at late time there remains the "tachyon matter" with finite energy but vanishing pressure. Description of the tachyon matter in the super BSFT was provided [12, 13] where the tachyon profile is linear,

$$T = u_0 X^0 \, . \tag{1.2}$$

(Note that this looks quite different from the boundary deformation above.) Taking $u_0 \to \pm 1$, one finds that only the late time limit ($x^0 \to \infty$) of (1.2) is a solution of equations of motion for the BSFT lagrangian derived for linear tachyon profiles. There is a small correction to (1.2) which vanishes exponentially at the late time. The pressure computed from the super BSFT action with the above solution vanishes at the late time limit.[2]

At the late time of the rolling tachyon, there should be no degree of freedom of propagating open strings, since the brane should have been annihilated by the tachyon condensation. To show this is the aim of this article. In [14], a tachyon effective field theory whose solution describes the tachyon matter was introduced, and it was shown that around that solution there is no plane wave fluctuation. Following it, [15] showed that this is also the case for gauge fluctuations in the BSFT action of [16] with the above linear tachyon profiles. An essential feature was found in [17] that in fact the propagation of the fluctuations in the BSFT action derived with the linear tachyon profiles is subject to a collapsed light cone given by *Carrollian contraction* of Lorentz group in which the speed of light becomes zero. But the problem is that, the BSFT action used in [17] is valid only for the linear profiles, while generic fluctuations are not of this type.

In this article, we show the collapse of the light cone of all the excitations of the open strings. We allow arbitrary fluctuations around the background (1.1) or (1.2), in the scheme of BSFT. The light cone structure can be read off from two-point functions in BSFT. For superstrings, the BSFT action is just a partition function of the worldsheet theory with generic boundary interactions [18], and in this sense the two-point function for gauge fields in BSFT was first computed in [19]. More systematically the tachyon two-point function

[2] We provide a description of the tachyon matter in bosonic BSFT in appendix A.

was computed in the appendix of [5], and also in [20]. (We will explore the two- and three-point functions in BSFT and its relation to the tachyon condensation, in our forth-coming paper [21].)

For example, the two-point function for the gauge fields in super BSFT can be computed in the following manner. The worldsheet boundary term for the gauge field is written as

$$I_{\rm B} = -i\int_{\partial\Sigma} d\tau \int dk \left(a_\mu(k) : \dot{X}^\mu e^{ik_\nu X^\nu} : -2f_{\mu\nu}(k) : e^{ik_\rho X^\rho} : \psi^\mu\psi^\nu\right), \tag{1.3}$$

where $a_\mu(k)$ is the momentum representation of the target space gauge field, $A_\mu(x) = \int dk\, a_\mu(k)e^{ikx}$, and $f_{\mu\nu}(k) = ik_\mu a_\nu - ik_\nu a_\mu$ is that of the field strength. The two point function of the gauge fields in the target space is just $\int dx\langle I_{\rm B} I_{\rm B}\rangle$, which can be evaluated with use of worldsheet propagators on the boundary of a unit disk. For the flat background with no tachyon condensation, they are[3]

$$\left\langle \hat{X}^\mu(\tau)\hat{X}^\nu(0)\right\rangle = -4\eta^{\mu\nu}\log\left|2\sin\frac{\tau}{2}\right|, \qquad \langle\psi^\mu(\tau)\psi^\nu(0)\rangle = \frac{1}{2\sin\frac{\tau}{2}}, \tag{1.5}$$

where the zero mode for X is already subtracted: $X^\mu = x^\mu + \hat{X}^\mu$. A straightforward calculation of $\int dx\langle I_{\rm B} I_{\rm B}\rangle$ exhibits the on-shell condition [21]. Massive excitations can be treated in the same manner, at least for two-point functions in the BSFT. We apply this strategy to extract the light cone structure of the open string excitations in the background linear profiles of the tachyon. Since we are interested in small fluctuations around the backgrounds, the two-point functions in the BSFT contain enough information on our concern. The open string excitations we treat here are of the standard normalization in no background coupling, and we would like to see how they behave once they are in the nontrivial tachyon background. What we shall find is that peculiar behavior of the worldsheet propagators, such as vanishing / diverging, results in the death of open strings — singular structures of the light cones.

2. Death of open strings

2.1 Spatial tachyon kink and descent relation

As a warm-up, we consider a spatial kink of a linear profile of the tachyon in a non-BPS D9-brane (1.1). In this background which is off-shell except $u_9 = 0$ or $\pm\infty$, the relevant worldsheet boundary propagator is written as [9]

$$\left\langle X^9(\tau)X^9(0)\right\rangle = 2\sum_{m\in\mathbf{Z}}\frac{1}{|m|+u_9^2}e^{im\tau}, \qquad \langle\psi^9(\tau)\psi^9(0)\rangle = -\frac{i}{2}\sum_{r\in\mathbf{Z}+1/2}\frac{r}{|r|+u_9^2}e^{ir\tau}. \tag{2.1}$$

[3] The worldsheet action we use here is

$$\frac{1}{4\pi}\int d^2z\left[\frac{2}{\alpha'}\partial_z X^\mu\partial_{\bar z}X_\mu + \psi^\mu\partial_{\bar z}\psi_\mu + \tilde{\psi}^\mu\partial_z\tilde{\psi}_\mu\right]. \tag{1.4}$$

The spacetime metric is taken to be diag$(-1,1,1,\ldots,1)$, and we put $\alpha' = 2$ in this article.

Here we included the zero mode in the propagators. When we take the $u_9 = \pm\infty$ solution representing a BPS D8-brane, the correlators for X^9 and ψ^9 vanish completely. Let us consider this effect on the two-point functions in the super BSFT in detail. First, the k^9 dependence in the resultant two-point functions in BSFT disappears. This is because the momentum k^9 is always coupled to $\hat{X}^9$ and so its contractions in $: e^{ik\cdot\hat{X}} :$ disappear when $|u_9| \to \infty$. The momentum representation of the BSFT action for the fluctuations is independent of k_9, which means that in the coordinate representation the action doesn't contain ∂_9 and especially any kinetic term along the x^9 direction. Therefore, in this background, any excitation of an open string cannot propagate in the x^9 direction. The relevant light cone collapses to the shape of a fan: the effective inverse metric appearing in the kinetic term becomes

$$G^{\mu\nu} = \lim_{u_9\to\infty} \mathrm{diag}\left(-1, 1, \ldots, 1, \mathcal{O}(1/u_9^2)\right) . \tag{2.2}$$

Physics at different values of x^9 are decoupled from each other.

An additional fact is that for this background the zero mode integral in the partition function gives a localization of the worldvolume [2, 5, 9], that is, an overall delta function $\delta(x^9)$ in front of the BSFT lagrangian. This means that the propagation only at the selected value $x_9 = 0$ survives, while the restricted propagation at $x_9 \neq 0$ has a trivial vanishing action. Therefore, the gauge fields exist only at the D8-brane worldvolume $x^9 = 0$ and the propagation is only along the D8-brane, which is consistent with the Sen's conjecture.

Furthermore, for the massless gauge bosons for example, the polarization parallel to X^9 should disappear in this background since the resultant brane has no worldvolume direction for that. This in fact occurs, since the kinetic term for this polarization comes from the contraction of $\dot{X}^9$ or ψ^9 in the vertex operators of the gauge fields (1.3).

The same mechanism is applicable to massive excitations, and all the excitations are subject to these constraints. They cannot propagate along x^9, they live at $x^9 = 0$, and the polarizations are restricted in such a way that the worldvolume theory is just that of the 8+1 dimensions. We have seen here that the vanishing of the worldvolume propagator results in the collapse of the light cone for the open string excitations.[4]

2.2 Rolling tachyon

Physics in the background of tachyon profile linear in time is, as we shall see, different from the above static case. In [12, 13], a time-dependent spatially-homogeneous tachyon decay was analysed in the scheme of BSFT, by adopting the linear tachyon profile (1.2). This profile is on-shell for $|u_0| = 1$ only at the late time limit, $x^0 \to \infty$, in the following sense: the BSFT equations of motion derived with the linear profile has a solution which differs from (1.2) by a small correction dumping exponentially in time. Writing $u_0^2 = 1 - \kappa$ with a small positive parameter κ, one obtains at the late time $\kappa \propto e^{-(x^0)^2/8}$ [12]. (Although this explicit form of κ may be corrected, we expect $\kappa \to 0$ in the late time behavior of the rolling tachyon as disscussed in [12].) This is the rolling tachyon in the super BSFT.

[4] In the renormalization group approach [22] or the boundary state approach [23], it was shown that the Neumann boundary condition for the open string turned to the Dirichlet boundary condition due to the tachyon condensation. This is also consistent with our result.

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For bosonic strings, see appendix A. The correction κ is time-dependent, but at late time its time-dependence becomes very small and we may treat this as a constant perturbation from $|u_0| = 1$.

The nature of the propagating fluctuations in this late-time rolling tachyon background is encoded in the BSFT two-point functions. We evaluate a super BSFT two-point function for the gauge field excitation of an arbitrary profile (1.3), around this background.[5] It is enough to consider the two-point functions since we are interested in a small fluctuation of a given open string field. The kinetic structure of the fluctuations may largely depend on the worldsheet propagators, as we have seen in the previous subsection. So let us look closely at the worldsheet propagators. The propagators in this background include

$$\langle X^0(\tau)X^0(0)\rangle = -2\sum_{m\in\mathbf{Z}} \frac{1}{|m|-u_0^2} e^{im\tau}\,, \qquad \langle \psi^0(\tau)\psi^0(0)\rangle = \frac{i}{2}\sum_{r\in\mathbf{Z}+1/2} \frac{r}{|r|-u_0^2} e^{ir\tau}\,. \tag{2.3}$$

Note that the sign in front of u^2 is different from that of the previous subsection due to the sign of the metric. Consequently, the behavior of the correlators is quite different from that of (2.1). Let us expand the correlator for X's in (2.3) for a small κ. In (2.3), first the term with $m = \pm 1$ is the leading order and in fact divergent in the limit $\kappa \to 0$. The next-to-leading order term is given by effectively shifting m by one in the summation, thus we obtain the expansion

$$\left\langle \hat{X}^0(\tau)\hat{X}^0(0)\right\rangle = -2\left[\frac{1}{\kappa}(e^{i\tau}+e^{-i\tau}) - e^{i\tau}\log(1-e^{i\tau}) - e^{-i\tau}\log(1-e^{-i\tau}) + \mathcal{O}(\kappa)\right]. \tag{2.4}$$

Note here that we have already subtracted the zero mode part, $m = 0$. As we shall see, this divergence is related to the collapse of the light cone.

For this X^0 direction, the change in the propagator forces us to redefine the normal ordering of operators. Denote the normal ordering with the propagator (2.3) as ${}^{\circ}_{\circ}\, e^{ik_\mu \hat{X}^\mu}\, {}^{\circ}_{\circ}$, then the relation to the usual normal ordering with the propagator (1.5) is

$$\begin{aligned} :e^{ik_\mu \hat{X}^\mu}: &= \exp\left[\frac{1}{2}(k_0)^2\left(\left\langle \hat{X}^0(\epsilon)\hat{X}^0(0)\right\rangle_{u_0=0,\epsilon=0} - \left\langle \hat{X}^0(\epsilon)\hat{X}^0(0)\right\rangle_{u_0,\epsilon=0}\right)\right] {}^{\circ}_{\circ}\, e^{ik_\mu \hat{X}^\mu}\, {}^{\circ}_{\circ} \\ &= \exp\left[(k_0)^2 \sum_{0\neq m\in\mathbf{Z}} \frac{u_0^2}{|m|(|m|-u_0^2)}\right] {}^{\circ}_{\circ}\, e^{ik_\mu \hat{X}^\mu}\, {}^{\circ}_{\circ}\,. \end{aligned} \tag{2.5}$$

This momentum-involved redefinition can be evaluated for the small κ as

$$:e^{ik_\mu \hat{X}^\mu}:= \exp\left[\frac{2}{\kappa}(k_0)^2 - \frac{\pi^2\kappa}{3}(k_0)^2 + \mathcal{O}(\kappa^2)\right] {}^{\circ}_{\circ}\, e^{ik_\mu \hat{X}^\mu}\, {}^{\circ}_{\circ}\,. \tag{2.6}$$

This additional factor can be absorbed into the definition of $a_\mu(k)$, if one wants. But we leave it because we would like to see what happens to the gauge field with the standard normalization defined with the normal ordering $:e^{ik_\mu \hat{X}^\mu}:$.

[5] The reason why we use gauge excitations first is that we know $A_\mu = 0$ is on-shell around the rolling tachyon (1.2). Treatment of the tachyon fluctuation needs a precaution because the background (1.2) is on-shell only in the late time limit.

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The gauge $A_0 = 0$ reduces the computation of the two-point function,

$$\langle I_{\rm B} I_{\rm B}\rangle = -\int dk d\tilde{k} \int \frac{d\tau}{2\pi}\left[a_i(k)a_j(\tilde{k})\left(\frac{-\delta_{ij}}{\sin^2\frac{\tau}{2}} - 4\tilde{k}_i k_j \cot^2\frac{\tau}{2}\right) + f_{ij}(k)f_{ij}(\tilde{k})\frac{-1}{\sin^2\frac{\tau}{2}} - \right.$$
$$\left. - 8 f_{i0}(k)f_{i0}(\tilde{k})\frac{1}{\sin\frac{\tau}{2}}\langle\psi^0(\tau)\psi^0(0)\rangle\right]\langle : e^{ikX}(\tau) :: e^{i\tilde{k}X}(0) :\rangle\,.$$

The correlator for the directions other than X^0 can be evaluated easily with (1.5) as

$$\left\langle : e^{ik_iX^i(\epsilon)} :: e^{i\tilde{k}_iX^i(0)} :\right\rangle = |1 - e^{i\epsilon}|^{4k_i\tilde{k}_j\delta^{ij}} e^{i(k_i+\tilde{k}_i)x^i}\,. \tag{2.7}$$

Using the expansion (2.4), we can include the direction X^0 and obtain

$$\iint \frac{d\tau_1 d\tau_2}{(2\pi)^2}\left\langle {}^{\circ}_{\circ}\, e^{ik_\mu\hat{X}^\mu(\tau_1)}\,{}^{\circ}_{\circ}\,{}^{\circ}_{\circ}\, e^{i\tilde{k}_\mu\hat{X}^\mu(\tau_2)}\,{}^{\circ}_{\circ}\right\rangle = \tag{2.8}$$
$$= \int_0^{2\pi}\frac{d\tau}{2\pi}|1-e^{i\tau}|^{4k_i\tilde{k}_i} e^{\frac{4}{\kappa}k_0\tilde{k}_0\cos\tau}(1-e^{i\tau})^{-2k_0\tilde{k}_0 e^{i\tau}}(1-e^{-i\tau})^{-2k_0\tilde{k}_0e^{-i\tau}}e^{\mathcal{O}(\kappa)}\,.$$

To evaluate this, we assume $k_0\tilde{k}_0 < 0$.[6] The integrand has a sharp peak at $\cos\tau \sim -1$ when κ is very small. We may treat it as a delta function with an appropriate normalization dependent on κ. For a finite $k_0\tilde{k}_0$ and in the limit $\kappa \to 0$, we obtain a formula for any κ-independent smooth function $g(\tau)$,

$$\int\frac{d\tau}{2\pi}g(\tau)\langle\,{}^{\circ}_{\circ}\, e^{ik\hat{X}}(\tau)\,{}^{\circ}_{\circ}\,{}^{\circ}_{\circ}\, e^{i\tilde{k}\hat{X}}(0)\,{}^{\circ}_{\circ}\,\rangle = g(\pi)\sqrt{\frac{-\kappa}{8\pi k_0\tilde{k}_0}}\exp\left[-\frac{4}{\kappa}k_0\tilde{k}_0 + 4\log 2k_i\tilde{k}_i\right]. \tag{2.9}$$

Then finally the nonzero-mode part of the two-point function reads

$$\langle I_{\rm B}I_{\rm B}\rangle = \int dk d\tilde{k}\left[a_i(k)a^i(\tilde{k}) + 2f_{ij}(k)f^{ij}(\tilde{k}) - 4f_{i0}(k)f^{i0}(\tilde{k})\,(3+\pi)\right]\times$$
$$\times\sqrt{\frac{-\kappa}{8\pi k_0\tilde{k}_0}}\exp\left[-\frac{4}{\kappa}k_0\tilde{k}_0 + 4\log 2k_i\tilde{k}_i + \frac{2}{\kappa}(k_0^2+\tilde{k}_0^2)\right]. \tag{2.10}$$

In this expression, the limit $\kappa \to 0$ gives a divergent factor, therefore this shows that any nonzero k_0 results in no dynamics. For any non-vanishing k_0, the two point function diverges and then if this serves as a kinetic term in a quantum field theory the path integral with non-vanishing k_0 is highly oscillatory, with which any correlation function vanishes.[7]

For the massive excitations, the same physics applies, which implies that there is no propagation of open strings in the rolling tachyon background at the late time. This is simply because all the open string excitations are accompanied with the momentum eigenfunction $: e^{ik\cdot\hat{X}} :$ whose contraction necessarily gives the divergence whenever it has a non-vanishing k_0.

[6]Now the background is time-dependent, so the momentum is not conserved, $k_0 + \tilde{k}_0 \neq 0$ generically. The case of $k_0\tilde{k}_0 > 0$ can be treated in a similar manner but with a steepest-descent method.

[7]Or saying it differently, the magnitude of the BSFT kinetic term roughly corresponds to an effective "tension" of the brane felt by the string excitations. In the present case this diverges and resultantly the open string modes effectively disappear since it is too heavy to excite.

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Thus, whole the open string excitations are subject to a light cone of the so-called Carrollian contraction of Lorentz group [17].[8] The light cone becomes a singular half-line and no open string can move on the original worldvolume of the brane. The upper limit of the speed of the propagations becomes zero. This describes how the open strings die in the late-time rolling tachyon.

It was suggested from the effective field theory that the spatial distribution of the open string condensations becomes arbitrary, and so there appears a huge degeneracy in the spatial configurations [17]. This can also be seen in the above result. Since the dependence on the spatial momentum $k_i\tilde{k}_i$ is in the next-to-leading order in κ, we can always set the small $\mathcal{O}(\kappa)$ dependence in k_0 in such a way that it cancels arbitrary function of the sub-leading order written by k_i. Therefore in the limit $\kappa \to 0$, any dependence of k_i can be absorbed into the vanishing k_0 and so allowed. This is the reflection of the fact that there is no interaction between different points on the original brane worldvolume, in the late-time limit of the rolling tachyon.

2.3 Rolling tachyon in background string charge

We may extend this relation between the divergence in the worldsheet propagators and the light cone structure of the excitations on the decaying brane, to the situation with a background string charge on the D-brane [24]. This is the setup used for the tree level open-closed duality stated in [25]. The duality statement is supported by a Nambu-Goto analysis [26] in which the string oscillations can propagate with a reduced speed of light proportional to the background string density E. This reduction of the speed was first observed in a low energy effective field theory [17]. Indeed, we will see in the BSFT two point function that the light cone structure in the rolling tachyon with the presence of this electric field E is just that. So, in this case, the open string excitations don't cease to propagate but a propagation along the background electric field $\vec{E}$ is allowed with the reduced speed.

In the background electric field $F_{01} \equiv E$ with the tachyon profile linear in time, the worldsheet propagator is given by [27]

$$\left\langle \hat{X}^\mu(\tau)\hat{X}^\nu(0)\right\rangle = \sum_{m=1}^{\infty}\left(M_+^{\mu\nu}e^{im\tau} + M_-^{\mu\nu}e^{-im\tau}\right), \tag{2.11}$$

where the nontrivial parts of the matrices $M_\pm$ are

$$M_+ = \begin{pmatrix} -m+u_0^2 & E \\ -E & m \end{pmatrix}^{-1} = \frac{1}{m(m-u_0^2)-m^2E^2}\begin{pmatrix} -m & mE \\ -mE & m-u_0^2 \end{pmatrix}, \tag{2.12}$$

$$M_- = \begin{pmatrix} -m+u_0^2 & -E \\ E & m \end{pmatrix}^{-1} = \frac{1}{m(m-u_0^2)-m^2E^2}\begin{pmatrix} -m & -mE \\ mE & m-u_0^2 \end{pmatrix}. \tag{2.13}$$

[8]Let us briefly mention the relation to the computations in [17] where the BSFT action [16] derived by the linear tachyon background was adopted. In our respect, the computations in [17] corresponds to a different limit, that is, taking $k_0 \to 0$ first and then $\kappa \to 0$, since there only the constant gauge field strength was considered.

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Our concern is the divergence appearing for $m = 1$, at $u_0 = \pm\sqrt{1-E^2}$ which is the rolling speed of the tachyon [24, 17]. We perturb it as before, $u_0^2 = 1 - E^2 - \kappa$ with a small positive κ, then the divergent parts of the relevant propagators are

$$\langle X^0(\tau)X^0(0)\rangle = \frac{-4}{\kappa}\cos\tau\,, \qquad \langle X^1(\tau)X^1(0)\rangle = \frac{4E^2}{\kappa}\cos\tau\,, \tag{2.14}$$

$$\langle X^0(\tau)X^1(0)\rangle = -\langle X^1(\tau)X^0(0)\rangle = \frac{2iE}{\kappa}\sin\tau\,. \tag{2.15}$$

The change of the normal ordering is[9]

$$: e^{ik_\mu \hat{X}^\mu} := \exp\left[\frac{2}{\kappa}\left(k_0^2 - E^2 k_i^2\right) + \mathcal{O}(1)\right] {}^{\circ}_{\circ}\, e^{ik_\mu \hat{X}^\mu}\, {}^{\circ}_{\circ}\,. \tag{2.16}$$

The exponential factor appearing in the partition function from the correlators among vertex operators reads

$$\exp\left[\frac{4}{\kappa}\left(k_0\tilde{k}_0 - E^2 k_1\tilde{k}_1\right)\cos\tau - \frac{2iE}{\kappa}\sin\tau\right]. \tag{2.17}$$

This factor has again a sharp peak at $\tau = \pi$ (when $k_0\tilde{k}_0 - E^2 k_i\tilde{k}_i < 0$), thus after an integration over τ, it results in a factor

$$\exp\left[-\frac{4}{\kappa}\left(k_0\tilde{k}_0 - E^2 k_1\tilde{k}_1\right) + \frac{2}{\kappa}\left(k_0^2 - E^2 k_1^2\right) + \frac{2}{\kappa}\left(\tilde{k}_0^2 - E^2\tilde{k}_1^2\right) + \mathcal{O}(1)\right], \tag{2.18}$$

where we have included the factor coming from (2.16). This diverges for nonzero $k_0^2 - E^2 k_1^2$ (or one or all k replaced by $\tilde{k}$). Thus we obtain a constraint for propagating degrees of freedom

$$k_0^2 - E^2 k_1^2 = \mathcal{O}(\kappa) \to 0\,. \tag{2.19}$$

which shows that only the propagation along $\vec{E}$ can be allowed and its speed is in fact E. This is consistent with the analysis of highly-oscillated Nambu-Goto strings [26].

In [17], this reduction of the speed of light was obtained for low energy limit of gauge field excitation in the rolling tachyon background. The computation of [17] can be reproduced from our perspective if we consider a constant field strength, which was the starting point of [17], as the low energy limit. This means that one took the limit $k \to 0$ before taking $\kappa \to 0$. Our present calculation shows that the reduction of the light cone structure appears not only for the very low energy but also for all the open string excitations with any momentum. The light cone becomes the shape of a fan, which can also be read from the effective inverse metric $G^{\mu\nu}$ defined when the exponent of (2.18) is rewritten as

$$-\frac{2}{\kappa}\left(k_\mu - \tilde{k}_\mu\right) G^{\mu\nu}\left(k_\nu - \tilde{k}_\nu\right), \qquad G^{\mu\nu} = \mathrm{diag}(-1, E^2, \underbrace{0, 0, \ldots, 0}_{\mathcal{O}(\kappa)})\,. \tag{2.20}$$

[9] On-shell gauge field in no tachyon condensation satisfies $k_0^2 - k_i^2 = 0$, so the exponent seen here is usually positive.

3. Summary and discussions

The essential point of this article is that the divergence or disappearance of the worldsheet boundary propagators crucially affects propagation of string excitations in the target space. The connection between the worldsheet and the target space is provided in the scheme of the BSFT. For the linear profile $T = u_9 X^9$ ($u_9 \to \pm\infty$), the X^9 propagator vanishes and resultantly the 99 component of the target space effective inverse metric is eliminated. For the rolling tachyon $T = u_0 X^0$ ($u_0 \to \pm 1$), the X^0 propagator diverges, and the spatial component of the target space effective inverse metric vanishes. Therefore the information of the target space metric is encoded in the ratios among various worldsheet propagators.

This feature is in fact found also in the worldvolume of a BPS Dp-brane without the tachyon condensation but with a constant gauge field strength. The effective open string metric of this case appears as a coefficient of the log part of the boundary propagator of the worldsheet bosons [28]: $G^{\mu\nu} = [1/(\eta + F)]^{(\mu\nu)}$ where the indices are symmetrized. First, taking a limit $F_{12} \to \infty$, we have vanishing propagators for X^1 and X^2, so following the argument in section 2.1, we observe that the light cone structure of the fluctuations is reduced to that of a D$(p-2)$-brane. This is in fact expected, since in this limit the bound charge of the D$(p-2)$-branes, mesured by the magnitude of F_{12}, diverges and the system is expected to be saturated by the collection of the D$(p-2)$-branes. Second, let us consider another limit $F_{01} \to 1$. The propagators for X^0 and X^1 diverge. We find a light cone similar to that of section 2.3 but with $E = 1$. In this case the bound fundamental strings whose charge diverges saturate the brane system so that only propagations along the strings are allowed. (See [29] for a related study of the light cones.)

The Carrollian contraction of the Lorentz group found here is based on the divergence of the two point functions of standard open string vertex operators. One can argue that this divergence might be cured by "renormalization", that is, a momentum-dependent field redefinition of the field. In fact, as indicated in (2.6), a natural normalization in this late-time rolling tachyon background may be with the normal ordering $\circ\!\!\!\circ\, e^{ik\cdot X} \circ\!\!\!\circ$. This causes the field redefinition of the form $a'(k) \equiv \exp[2k_0^2/\kappa]a(k)$ ($\kappa \to 0$), but even with this normalization, the two-point function diverges, due to the contraction among the correlators: $\exp[-4k_0\tilde{k}_0/\kappa]$.

Supposing that k_0 and $\tilde{k}_0$ are of the same order, a stronger renormalization $a'(k) \equiv \exp[4k_0^2/\kappa]a(k)$ looks to make the divergent factor eliminated and might give a canonically normalized kinetic term for the fluctuations. However, more severe divergence can be found in general n-point functions with $n > 2$. Since the n-point function includes arbitrary contraction among n operators of the form $:e^{ik_0 X^0}:$, it may diverge much strongly than what can be eliminated with the above field redefinition. Thus the "renormalized" two-point function becomes rather meaningless. The resultant theory appears to be a strongly interacting theory and no perturbative state defined with the renormalized $:e^{ik\cdot X}:$ may appear as an asymptotic state, unless $k_0 = 0$. This is consistent with the picture of confinement of open strings at the tachyon vacuum [7].

In cubic string field theory [3], various attempts have been made to understand the open string excitations at the true vacuum [30, 31, 32]. These attempts are for static homogeneous tachyon condensation in which the constant tachyon is put to its true vacuum. Thus a direct relation to our BSFT approach is unclear. For the constant tachyon in BSFT, the worldsheet boundary receives a weight factor $\sim e^{-T^2}$ which results in the vanishing of the worldvolume D-brane action, while in CSFT similar structure was found [31]. Due to this boundary factor, the boundary shrinks effectively and disappear [5], which seems to be realized also in CSFT [32]. Further concrete relations among the two SFT's should be clarified in the future.

The rolling tachyon process has been considered much in deformed conformal field theories, especially boundary Liouville theories [33, 34]. It was pointed out in [35] that the late-time spectrum has vanishing wave functions while its k_0 dependence is diverging, which looks quite consistent with our results of the divergence of the two-point functions. It is interesting to seek for a relation of our BSFT approach to the creation of open strings due to the change of the vacua in the conformal field theories [33].

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A. Tachyon matter in bosonic BSFT

In this appendix, we present a description of tachyon matter in bosonic string theory, as a solution of a bosonic BSFT. This is an analogue of [12, 13] given for superstring. The analysis presented in this article for the light cone structure is applicable to the bosonic strings in the same manner.

BSFT action. First, let us construct a BSFT action by closely following the original construction [2] but slightly generalizing it to include a time-like target-space direction X^0. We introduce a tachyon profile at the boundary of the worldsheet,

$$S_{\rm bdry} = \int \frac{d\theta}{2\pi}\, T(X)\,, \qquad T(X) \equiv a + \sum_{i=1}^{25} \frac{u_i}{4}(X^i)^2 - \frac{u_0}{4}(X^0)^2\,, \tag{A.1}$$

where a, u_i and u_0 are boundary deformation parameters. We have defined this with the minus sign in front of the X^0 mass term so that the propagators for X's have the following common form:

$$\left\langle (X^i)^2(\theta)\right\rangle = \frac{2}{u_i} - \sum_{m=1}^{\infty} \frac{4u_i}{m(m+u_i)}\,, \qquad \left\langle (X^0)^2(\theta)\right\rangle = -\left(\frac{2}{u_0} - \sum_{m=1}^{\infty}\frac{4u_0}{m(m+u_0)}\right).$$

Following the definition of the partition function given with (A.1), we find

$$\frac{d}{du_i}\log Z_1 = -\frac{1}{8\pi}\int_0^{2\pi} d\theta \left\langle (X^i)^2(\theta)\right\rangle, \qquad \frac{d}{du_0}\log Z_1 = \frac{1}{8\pi}\int_0^{2\pi} d\theta \left\langle (X^0)^2(\theta)\right\rangle.$$

Obviously the partition function is a product of a function of u_i and that of u_0, so we obtain (up to a constant overall factor)

$$Z = e^{-a} Z_1(u_0) \prod_{i=1}^{25} Z_1(u_i)\,, \qquad Z_1(u) \equiv \sqrt{u}e^{\gamma u}\Gamma(u)\,. \tag{A.2}$$

In the derivation of the action from the partition function, which was described in [2] in detail, it is easy to find that the inclusion of X^0 is just done by regarding this X^0 as iX^{26}, and accordingly $u_0 = u_{26}$. Then

$$S = \left(1 + a + u_0 - u_0 \frac{\partial}{\partial u_0} + \sum_{i=1}^{25}\left(u_i - u_i \frac{\partial}{\partial u_i}\right)\right) Z(a, u_0, u_i)\,. \tag{A.3}$$

For our purpose, we turn on only u_0 and u_1. (Although a homogeneous rolling tachyon solution is obtained with $u_1 = 0$, we need this u_1 dependence to evaluate its pressure later.) We need a lagrangian written in terms of the target space tachyon field $T(x)$ defined in (A.1). Using the following formulas

$$\int dx^1 dx^0\, e^{-T} = e^{-a}\frac{4\pi}{\sqrt{-u_0 u_1}}\,, \qquad \int dx^1 dx^0\, T e^{-T} = e^{-a}\frac{4\pi}{\sqrt{-u_0 u_1}}(a+1)\,, \tag{A.4}$$

we obtain the lagrangian as

$$L = e^{-T}\frac{\sqrt{-u_0 u_1}}{4\pi}\left[T + u_0 - u_0\frac{\partial}{\partial u_0} + u_1 - u_1 \frac{\partial}{\partial u_1}\right] Z_1(u_0) Z_1(u_1)\,. \tag{A.5}$$

We may replace the parameters u by

$$u_1 \to 2\partial_1\partial_1 T\,, \qquad u_0 \to -2\partial_0\partial_0 T \tag{A.6}$$

which follow from the definition of the tachyon profile (A.1). With the simple substitution of these expression to (A.5), we obtain a lagrangian written solely by the tachyon field. Note that in this lagrangian the covariance is not manifest, since we have suppressed the dependence on $\partial_1\partial_0 T$ in the computation. However, the Lorentz invariance can be restored in the following way. The lagrangian (A.5) can be Taylor-expanded in terms of u_0 and u_1. The first nontrivial term comes like $u_0 + u_1$ times a factor of a function of T such as e^{-T}. We can rebuild its Lorentz-invariant form

$$u_0 + u_1 = 2\eta^{\mu\nu}\partial_\mu\partial_\nu T\,. \tag{A.7}$$

So we might expect that the lagrangian can be written only by this invariant, as is usually used for BSFT lagrangians for superstrings.[10] However this is not the case. The next-to-leading term has the following form

$$f(T)\left((u_0)^2 + (u_1)^2\right) + g(T) u_0 u_1\,. \tag{A.8}$$

This is lifted to its Lorentz-invariant expression,

$$4f(T)\eta^{\mu\nu}\eta^{\rho\sigma}\partial_\mu\partial_\rho T\partial_\nu\partial_\sigma T - 2g(T)\left((\eta^{\mu\nu}\partial_\mu\partial_\nu T)^2 - \eta^{\mu\nu}\eta^{\rho\sigma}\partial_\mu\partial_\rho T \partial_\nu\partial_\sigma T\right). \tag{A.9}$$

Higher order terms include much more intricate Lorenz contractions but they can be uniquely lifted to their Lorentz-invariant expressions.

[10] Note that even for superstrings we face a similar problem when we consider D-$\overline{\mathrm{D}}$ or several D-branes [36].

Rolling solution. Let us find a time-dependent homogeneous solution. Instead of solving an equation of motion, homogeneous time-dependent solution can be easily obtained by solving the energy-conservation condition. The easiest way to get the hamiltonian from the above lagrangian (A.5) is to couple the system to a background metric and derive the energy-momentum. In the present case all the nontrivial Lorentz contractions reduce to the simple single invariant $u_0 = -2(v\partial_0(v\partial_0 T))$ where $v \equiv \sqrt{-g^{00}}$. Our action is written as

$$S = \int dx^0 \, \frac{1}{v} e^{-T} (F(u_0) + TG(u_0)) \tag{A.10}$$

where

$$F(u_0) \equiv \sqrt{-u_0} \left(\frac{1}{2} + u_0 - u_0 \frac{\partial}{\partial u_0} \right) Z_1(u_0) \, , \qquad G(u_0) \equiv \sqrt{-u_0} Z_1(u_0) \, . \tag{A.11}$$

We find the hamiltonian $H = -[\partial S/\partial v]_{v=1}$ as

$$H = e^{-T}(F + TG) - 2\partial_0 \left(\partial_0 T e^{-T} (F' + TG') \right) + 4\partial_0^2 T \left(e^{-T} (F' + TG') \right) . \tag{A.12}$$

The tachyon matter is defined as the late time residual matter when the tachyon rolls to its true vacuum $T = \infty$. So we are interested in the late time behavior of the system given by the above hamiltonian (A.12). First, u_0 should be negative so that the tachyon may not roll down in the wrong side of the potential. The hamiltonian should be conserved, while the overall exponential factor in the above expression is vanishing as $e^{-T} \sim \exp[u_0 (x^0)^2/4]$. So we expect that the other factors in the hamiltonian diverge. In fact, from the explicit expression for Z_1, we find, at $u_0 \sim -1$,

$$F(u_0) \sim (1 + u_0)^{-2} \, , \qquad G(u_0) \sim (1 + u_0)^{-1} \, . \tag{A.13}$$

The most diverging factor comes from the second term in (A.12),

$$H \sim -2\partial_0 T e^{-T} (-2\partial_0^3 T) F'' \, . \tag{A.14}$$

We consider a small fluctuation around the tachyon profile of $u_0 = -1$,

$$T = a + \frac{1}{4}(x^0)^2 + \epsilon(x^0) \, . \tag{A.15}$$

Then the energy conservation implies

$$x^0 e^{-(x^0)^2/4} (\partial_0^3 \epsilon)(\partial_0^2 \epsilon)^{-4} = \text{const.} \tag{A.16}$$

This can be solved as

$$\ddot{\epsilon} \sim (x^0)^{2/3} \exp\left[-\frac{(x^0)^2}{12} \right] . \tag{A.17}$$

So the fluctuation vanishes in the late time limit, which shows the consistency of our ansatz.

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Pressure. Finally we would like to show that the pressure of the homogeneous solution obtained here is exponentially dumping. The pressure along x^1 direction should be defined as

$$p_1 = 2\frac{\partial L}{\partial g^{11}}\,. \tag{A.18}$$

To this end we have to work with nonzero u_1 since in it nontrivial dependence on g^{11} might appear in the covariant derivatives. As seen above, the whole structure of the Lorenz contraction is very nontrivial. But here we know that in the end we put the metric to be constant and also the tachyon to be dependent only on time, which results in a lot of simplification. The covariant derivatives appearing in the expression is

$$\nabla_\mu \partial_\nu T = \partial_\mu \partial_\nu T - \Gamma^\rho_{\nu\mu}\partial_\rho T \tag{A.19}$$

where the Christoffel symbol is defined as usual, $\Gamma^\rho_{\nu\mu} = (1/2)g^{\rho\sigma}(\partial_\mu g_{\sigma\nu} + \partial_\nu g_{\sigma\mu} - \partial_\sigma g_{\mu\nu})$. The index ρ will be contracted with $\partial_\rho T$ and so it becomes irrelevant unless $\rho = 0$. We want to see the dependence on g^{11}, then the only term which we find relevant is

$$\Gamma^0_{11} = \frac{1}{2}\partial_0\left(\frac{1}{g^{11}}\right). \tag{A.20}$$

As well as the Christoffel symbols, the metrics used for contracting the indices of the derivatives can be relevant in the calculation of the energy-momentum tensor, but this contribution turns out to be vanishing. This is because after the differentiation with respect to g^{11} the remaining term always include ∂_1 which is zero when acting on $T(x^0)$. Thus the covariant derivative ∇_1 can appear only once in the computation of the pressure, so we need only a linear term in u_1.

In the evaluation of the pressure, we may replace u_1 by

$$u_1 \sim -\partial_0\left(\frac{1}{g^{11}}\right)\partial_0 T\,. \tag{A.21}$$

The lagrangian can be written as, after a little calculation,

$$L = \frac{1}{4\pi}e^{-T}\sqrt{-u_0}\left(\frac{1}{2} + T + u_0 - u_0\frac{\partial}{\partial u_0} + u_1\right)Z_1(u_0) + \mathcal{O}(u_1^2)\,. \tag{A.22}$$

so, restoring the g^{11} dependence including the overall factor $\sqrt{-g} \sim (g^{11})^{-1/2}$, we obtain

$$p_1 = -L - 2\partial_0\left(\partial_0 T L\right). \tag{A.23}$$

The first term comes from $\sqrt{-g}$ while the second term comes from the Christoffel in u_1. Writing the lagrangian in terms of F and G again, we find that the most dominant part in the pressure comes from

$$-\frac{1}{2\pi}\partial_0 T e^{-T}\partial_0 F = \frac{1}{\pi}\partial_0 T e^{-T}\partial_0^3 T F' \sim \exp\left[-\frac{(x^0)^2}{6}\right]. \tag{A.24}$$

In the last part we substituted the solution (A.17). This shows that the pressure is vanishing at the late time of the rolling tachyon. Technically speaking, the reason why the hamiltonian is conserved while the pressure is dumping is that, in the expressions in terms of F, the number of the derivatives acting on F is less by one in the pressure compared to that of the hamiltonian.

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