Interpolating between open and closed strings – a BSFT approach

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Abstract

We address the conjecture that at the tachyonic vacuum open strings get transformed into closed strings. We show that it is possible in the context of boundary string field theory to interpolate between the conventional open string theory, characterized by having the D25 brane as the boundary state, and an off-shell (open) string theory where the boundary state is identified with the closed string vacuum, where holomorphic and antiholomorphic modes decouple and where bulk vertex operator correlation functions are identical to those of the closed string.
Introduction

In recent years there has been a significant progress in the understanding of open string tachyon condensation\textsuperscript{4}, i.e. the passage from the perturbative unstable open string vacuum to the “true” tachyonic vacuum with a non-vanishing tachyon condensate\textsuperscript{5}.

The tachyon condensation of the open bosonic string has been studied using cubic open string field theory (OSFT) \textsuperscript{3}, boundary string field theory (BSFT) \textsuperscript{4, 5} and most recently vacuum string field theory (VSFT) \textsuperscript{6}. The use of OSFT typically involves level truncation and quite formidable calculations. VSFT is the conjectured form of OSFT when expanded around the tachyonic vacuum.

Of the three approaches, BSFT is on the least firm footing as it has problems with unrenormalizable boundary interactions, but it allows for the most explicit verification of Sen’s conjecture about the relation between the tension of $D$-branes and the string field theory action \textsuperscript{7, 8}.

In this paper we discuss the conjecture that in the tachyonic vacuum the open strings disappear and one is left with the conventional closed string theory. There are two ways in which this can be realized. The simplest (and least interesting way) is the one where the starting point is open and closed string theory in co-existence. In the process of open string tachyon condensation, all possible excitations of the open string disappear and one is left with the closed string theory. The other possibility (which we will investigate here) is that one starts with just the theory of open strings. The disappearance of the open string then implies the following: as one moves toward the tachyonic vacuum the excitations of the open string (even at tree level) should become equivalent to closed string excitations. Of course it is well known that closed strings appear in open string loop diagrams, but the string world sheets considered here will always be those of open strings, in particular the above should also be true for the disk.

The majority of works studying the appearance of closed strings around the tachyonic vacuum have been based on considerations of space-time effective actions \textsuperscript{9, 10, 11}. Apart from that, the problem of the disappearance of open strings has been studied using level truncation methods in OSFT \textsuperscript{12, 13}, and is in a way built into the foundation of VSFT from the start (see also \textsuperscript{14}). In this paper we adopt a world-sheet perspective and use the formalism of BSFT.

The configuration space of BSFT is the space of boundary interactions (boundary field theories) in a certain (closed string) background, which we here take to be flat space-time. Since a general boundary interaction will not result in a conformally invariant world-sheet theory, this is indeed a way of probing general off-shell open strings. Once the boundary interactions are parameterized, there is a concrete prescription for calculating the value of the BSFT action in terms of these parameters

\textsuperscript{4}For old work on tachyon condensation, see \textsuperscript{1, 2}.

\textsuperscript{5}Since we consider only the bosonic string the effective action is of course unbounded from below. What we call the “true” vacuum is thus only a local minimum.
(up to renormalization ambiguities). A minimization of the BSFT action gives a flow from the unstable perturbative vacuum to the true tachyonic vacuum.

The points in the BSFT configuration space along this flow correspond to two-dimensional field theories which encode how the properties of the open string get modified in the process of decay to the tachyonic vacuum. These theories are in general not conformal and usually mix holomorphic and antiholomorphic fields through the boundary interactions.

The main question that we address in this paper is whether in the BSFT configuration space one can find a point which can be identified with the CFT of ordinary closed string theory. In particular, for such a theory there should be a decoupling between holomorphic and antiholomorphic field modes, contrary to the generic situation in open string theory.

We will show that one can find a continuous interpolation between the ordinary open string theory and an open string theory with special boundary interactions, where all bulk excitations (vertex operators) are identical to closed string excitations (see below for a precise statement), and where holomorphic and antiholomorphic fields indeed decouple. Finally we perform a first analysis of the behavior of the BSFT action for such a family of boundary interactions.

**Setup**

Let us consider a correlation function with a number of (closed string) on-shell vertex operators on a disk (entering the amplitude of scattering of closed string states off an open string). For the simplest case of (closed string) tachyons we have the well known formula

\[ \langle e^{ik_1X(z_1,\bar{z}_1)} \cdots e^{ik_nX(z_n,\bar{z}_n)} \rangle_{\text{open}} \sim \prod |z_i - z_j|^{2k_i k_j} |1 - z_i \bar{z}_j|^{2k_i k_j} (1 - |z_i|^2)^{k_i^2} \]  

The presence of the open string boundary can be seen through the last two terms which mix holomorphic and antiholomorphic coordinates. Once we go off-shell (with respect to the open string) and add a boundary interaction eq. (1) is modified. A necessary condition for the transmutation of open string excitations into closed string excitations is that (1) gets modified to the standard closed string correlation function

\[ \langle e^{ik_1X(z_1,\bar{z}_1)} \cdots e^{ik_nX(z_n,\bar{z}_n)} \rangle_{\text{closed}} \sim \prod |z_i - z_j|^{2k_i k_j} \]  

The analysis of the tachyon condensation is usually performed using just the quadratic tachyon profile \( T(X) = a + u X^2 \) and the end-point of tachyon condensation is for \( a \to \infty \) and \( u \to 0 \). At this point the amplitude (1) remains unmodified. Additional boundary interactions are necessary in order to reach (2).

The goal of this paper is to construct a family of boundary interactions which interpolates smoothly between (1) and (2).
Matter boundary interaction

We consider a quadratic nonlocal interaction of the general form considered by Li and Witten [15]:

\[ S_B = a + \frac{1}{8\pi} \int d\theta d\theta' X(\theta)u(\theta - \theta')X(\theta') \]  
(3)

where

\[ u(\theta) = \frac{1}{2\pi} \sum_n u_n e^{in\theta} \]  
(4)

Our interpolating boundary interaction is defined by the choice

\[ u_n = t|n|e^{-|n|\epsilon} \]  
(5)

where \( t \) is a coupling constant and \( \epsilon \) is an UV cut-off. In the final expressions we should set the cut-off \( \epsilon \rightarrow 0 \).

It is convenient for later use to express the boundary interaction in terms of Fourier components of the field \( X(\theta) \)

\[ X(\theta) = \sum_{n=-\infty}^{\infty} X_n e^{in\theta} \]  
(6)

We have

\[ S_B = a + \frac{t}{2} \sum_{n=1}^{\infty} nX_{-n}X_n \]  
(7)

This is an ‘almost local’ interaction. Decomposing \( X = X_+ + X_- + x_0 \) into positive, negative and zero modes we have a local form:

\[ S_B = a + \frac{t}{8\pi} i \int d\theta \left( \frac{\partial X_-}{\partial \theta} - X_- \frac{\partial X_+}{\partial \theta} \right) \]  
(8)

The boundary action (3) leads to modified boundary conditions, whose general form is presented in [15].

We will now analyze the transmutation of open strings into closed strings through the boundary interaction (3) from three different points of view.

I — Green’s function

The Green’s function for (3) has been derived by Li and Witten [15]. For our choice of couplings \( u_k = t \cdot |k| \), suppressing the divergent part coming from the zero-mode, one is left with

\[ G(z, w) = -\left( \log |z - w|^2 + \frac{1 - t}{1 + t} \log |1 - \bar{z}w|^2 \right) \]  
(9)

It is seen that (3) interpolates between the propagator for the open string (\( t=0 \)) and the propagator for the closed string (\( t=1 \)). It follows that all correlation functions will satisfy

\[ \langle V_1(z_1) \ldots V_n(z_n) \rangle_{\text{closed}} = \langle V_1(z_1) \ldots V_n(z_n) \rangle_{\text{open with } t = 1} \]  
(10)
II — Cutting and patching holes

Let us now consider, from a different point of view, the special relation between open strings with $t=1$ boundary interactions and closed strings.

Any closed string correlation function with arbitrary insertions inside the unit disk may be written as a path integral

$$\int D X V(z_1) \ldots V(z_n) e^{-S_{\text{closed}}} \tag{11}$$

where the fields $X(z, \bar{z})$ are defined on the complex sphere. We can factorize the path integral into an integral over the fields inside the unit disk with fixed boundary values $X(\theta) = \sum_n X_n e^{in\theta}$, an integral over the values of the fields outside the disk with the boundary values $X(\theta)$ and finally an integral over the boundary values themselves. Let us perform first the integral outside the disk. Following [16] we split the fields outside the unit disk into a classical piece (regular at infinity) which solves the equations of motion and saturates the boundary conditions, and a piece $x_D(z, \bar{z})$ with Dirichlet boundary conditions on the unit circle and regular at infinity:

$$X(z, \bar{z}) = X_0 + \sum_{n=1}^{\infty} X_{-n} z^{-n} + X_n \bar{z}^{-n} + x_D(z, \bar{z}) \tag{12}$$

Substituting it into the action we get

$$\text{Normalization} \cdot \exp\left(-\frac{1}{2} \sum_{n=1}^{\infty} nX_{-n}X_n\right), \tag{13}$$

where Normalization comes from the path integral over $x_D$ (and is equal to the partition function for the open string with Dirichlet boundary conditions). The boundary term in eq. (13) is exactly our boundary action (7) with $t=1$. At this stage we are left with an integral over the fields on the disk with this additional boundary action, i.e. precisely an open string correlation function with boundary action $S_B(t=1)$.

By construction, the expression (13) is of course just the Schrödinger representation of the closed string ground state wave function and is the simplest example of the operator–state mapping of conformal field theory [1].

III — Boundary state

We now construct the boundary states corresponding to our family of boundary interactions. It will enable us to look at the meaning of the disappearance of the open string (D25-brane) from yet another point of view.

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6 In this sense one can in principle obtain effective boundary interactions mimicking the insertion of arbitrary vertex operators outside the disk, but they will in general correspond to singular boundary interactions because the corresponding wave functions will have zeros.
The prescription for associating a boundary state with a given, not necessarily conformal, boundary interaction is [17]

\[ |B\rangle = \int dx_n d\vec{x}_n e^{-S_b} |x, \vec{x}\rangle \] (14)

where

\[ |x, \vec{x}\rangle = \prod_{n=1}^{\infty} e^{-\frac{1}{2} x_n^2 - a_n a_n^\dagger + x_n a_n^\dagger + \vec{x}_n \vec{a}_n^\dagger} |0\rangle \] (15)

Here \( x_m = X_m / \sqrt{m} \) and \( \vec{x}_m = X_{-m} / \sqrt{-m} \). The integration in (14) is Gaussian and the result is

\[ |B\rangle = \sqrt{1 + t} e^{-a} \prod_{n=1}^{\infty} e^{\frac{1 + t}{1 + t} a_n a_n^\dagger} |0\rangle \] (16)

The factor \( \sqrt{1 + t} \) comes from the infinite product \( \prod_n (1 + t)^{-1} \) regularized using the \( \zeta \) function prescription.

For \( t = 0 \) the boundary state can be understood as defining the D25 brane, or equivalently the conventional open string. When we increase \( t \) the properties of the D25 brane get modified as the boundary interaction of open string is turned on. In particular, the form of the boundary states shows that for \( t < 1 \) there is a coupling between holomorphic and antiholomorphic modes. When \( t \) reaches one the holomorphic and antiholomorphic modes decouple and all creation operators disappear: we are left with the closed string oscillator out-vacuum! Thus the interpolation between \( t = 0 \) and \( t = 1 \) may be viewed as the disappearance of the D25 brane, and the emergence of the closed string vacuum.

**Ghost boundary interaction**

Until now we have completely ignored the ghost sector of the bosonic string theory. As the matter sector is changing with increasing \( t \) in the direction of decoupling holomorphic and antiholomorphic modes, it is natural to assume that the ghost sector should also be changing. Indeed the arguments which show \( S_B(t = 1) \) to originate from cutting out a disk in the closed string world-sheet and integrating over fields outside the disk can also be applied to the ghost fields. Rather than reporting on this construction (which only applies for \( t = 1 \)) we use the general formalism of boundary states to construct a suitable ghost action which interpolates between an ordinary open string ghost sector and one in which the holomorphic and antiholomorphic ghosts are decoupled.

The standard ghost field mode expansions are

\[ b(z) = \sum_n b_n z^{-n-2} \quad c(z) = \sum_n c_n z^{-n+1} \] (17)
and analogous ones for $\tilde{b}$ and $\tilde{c}$. In the open string the only fields surviving on the boundary are $b_{ab}^a b^b$ and $c_t \equiv c_\theta t^a$ (denoted by $c(\theta)$ in [15]) and which we will denote here by $B(\theta)$ and $C(\theta)/(2i)$ respectively. These fields have the following mode expansions

$$B(z) = \sum_n (b_n + \tilde{b}_{-n}) e^{-in\theta} \quad C(z) = \sum_n (c_n - \tilde{c}_{-n}) e^{-in\theta}$$

(18)

The starting point of the construction of boundary states is to introduce classical anti-commuting (Grassmann) fields on the boundary with Fourier components $B_n$, $C_n$ and form coherent states (similarly to the construction in [17]):

$$(b_n + \tilde{b}_{-n}) |B_{coh}\rangle = B_n |B_{coh}\rangle$$

$$(c_n - \tilde{c}_{-n}) |B_{coh}\rangle = C_n |B_{coh}\rangle$$

(19)

Then the boundary state corresponding to a boundary action $S_g(B_n, C_n)$ will be given by

$$|B_g\rangle = \int \prod_n dB_n dC_n e^{-S_g(B_n, C_n)} |B_{coh}\rangle$$

(20)

where $|B_{coh}\rangle$ depends on $B_n$ and $C_n$.

The solution of equations (19) is

$$|B_{coh}\rangle = \mathcal{N} \prod_{n=1}^\infty e^{-\tilde{c}_n^b b_n^b - C_n b_n^b - B_n b_n^b} e^{-\tilde{c}_n^c b_n^c + C_n b_n^c - B_n b_n^c} |Z\rangle$$

(21)

where $|Z\rangle = \frac{c_0 + \tilde{c}_0}{2} c_1 \tilde{c}_1 |q = 0\rangle$ (see e.g. [18]). We fix the normalization factor $\mathcal{N}$ so that the boundary state from (20) with zero action will give the ordinary open string ghost boundary state:

$$\exp \left( \sum_{n=1}^\infty \left( \tilde{c}_n^b b_n^b + c_n^c b_n^c \right) \right) |Z\rangle$$

(22)

With this choice we have

$$\mathcal{N} = \prod_n 4(1 - \frac{1}{2} C_n B_n)(1 + \frac{1}{2} C_n B_n)$$

(23)

We want now to find a ghost boundary action which will decouple the left- and right- moving ghosts, similarly to what happened in the matter sector.

The simplest action that satisfies our requirements is

$$S_g = \frac{g}{2} \sum_{n=1}^\infty (C_n B_n - C_n B_{-n})$$

(24)

\footnote{We write now explicitly the tensor indices of the ghost fields, and contract them with normal $n^a$ or tangent $t^a$ vectors to the boundary.}
With this choice the ghost boundary state following from (20) is

$$|B_g\rangle = \left(\prod_n (1 + g)^2\right) \cdot \exp\left(\frac{1 - g}{1 + g} \sum_{n=1}^{\infty} \left(\tilde{c}^+_n b^+_n + c^+_n \tilde{b}^+_n\right)\right) |Z\rangle$$

(25)

The action (24) can be rewritten as a nonlocal action of the form

$$\frac{1}{4\pi} \int d\theta d\theta' B(\theta) v(\theta - \theta') C(\theta')$$

(26)

with $v(\theta - \theta') = \frac{1}{2\pi} \sum_n v_n e^{in(\theta - \theta')}$ where

$$v_n = -g \quad n > 0$$

$$v_n = g \quad n < 0$$

(27)

(28)

The above action is quite natural and involves just the combinations of ghost fields which are present on the boundary for the ordinary open string. The nonlocal interaction is of the same type as (3), and (25) is the “ghost version” of (16). However, a priori, the matter boundary interaction with $u(\theta)$ does not determine a specific ghost interaction with some $v(\theta)$. Similarly to the situation for matter fields we have a decoupling between holomorphic and antiholomorphic ghost modes for $g=1$.

**Conformal properties**

It is interesting to consider the conformal properties of the boundary interactions. A conformally (Virasoro-) invariant boundary state $|B\rangle$ in a boundary conformal field theory satisfies

$$L_n |B\rangle = \tilde{L}_{-n} |B\rangle$$

(29)

where $L_n = \sum_k a_k a_{n-k}$. This is of course true for $t=0$. Our choice of matter action (7) provides a minimal deviation from this situation in the sense that the state (16) is still $SL(2,R)$ invariant, but not Virasoro invariant as one can easily check. Indeed, (29) is only satisfied for $n = 0, \pm 1$, in accordance with the fact that the boundary interaction will move us off-shell. The point $t=1$ is special by restoring the $SL(2,C)$ invariance of the closed string vacuum.

Correspondingly, if we consider the inclusion of ghost interactions the point $t=g=1$ is singled out by being BRST invariant:

$$Q |B_{t=1}\rangle |B_{g=1}\rangle = 0$$

(30)

where $Q$ is the closed string BRST operator

$$Q = \sum_{n=-\infty}^{\infty} c_n L_{-n}^X + \sum_{n=-1}^{\infty} c_{-n} L_n^{gh} + \sum_{n=2}^{\infty} L_{-n}^{gh} c_n + c.c.$$
BSFT analysis

BSFT is defined on the space of all boundary interactions

\[ S_B = \sum_i \int d\theta \lambda_i \mathcal{V}_i(X, b, c) \]  

(32)

through a choice of a corresponding family of ghost number 1 operators:

\[ O(\theta) = \sum_i \lambda_i O_i(\theta) \]  

(33)

where \( b_{-1}O_i = \mathcal{V}_i \). Here \( b_{-1} \) is an operator which when expressed in terms of closed string modes is \( i(b_0 - \bar{b}_0) \). As emphasized in [4] there may be an ambiguity in the choice of \( O_i \) for a given \( \mathcal{V}_i \). The BSFT action is defined as a function of the parameters \( \lambda_i \) through the differential equations

\[ \frac{\partial S}{\partial \lambda_i} = \frac{1}{2} \int d\theta \int d\theta' \langle O_i(\theta) \{ Q, O \}(\theta') \rangle, \]  

(34)

where \( \langle \ldots \rangle \) is the unnormalized correlation function. Since \( d^2S = 0 \) the above equations determine, at least locally, a well defined action. When the \( O_i \) come from matter weight 1 primary operators, the r.h.s. of (34) can be rewritten as \( \beta^j G_{ij}(\lambda) \), where \( \beta^j(\lambda) \) are the \( \beta \)-functions for the renormalization group flow in the set of boundary field theories, and the fixed points are defined by \( \beta^j(\lambda^*) = 0 \). Our choice of coupling constants (3) was partially motivated by the fact that a calculation using the boundary interaction (3)-(4) gives [15]

\[ \beta_n(\{ u_k \}) \sim \frac{1}{2} n(u_{n+1} - u_{n-1}) - u_n, \quad n > 0. \]

which is zero for the choice (3) of couplings in the limit \( \varepsilon = 0 \). Indeed we found earlier that the boundary state is \( SL(2,R) \) invariant and hence scale invariant.

In principle our program is as follows: calculate the renormalization flow (34) in terms of the coupling constants \( a, t, g \) and show that \( t, g = 1 \) is a fixed point which can be reached along a trajectory with decreasing BSFT action. Ultimately this fixed point should have the following property: for generic boundary operators one has \( \partial S/\partial \lambda_i = 0 \) at \( t = g = 1 \). This would mean that the boundary really disappears and such perturbations could be considered as symmetries of the theory.

Unfortunately the calculations are non-trivial for two reasons. As noticed already in [13] there is a non-trivial UV cut-off dependence in the theory (here explicitly present in the innocently looking \( \varepsilon \) in the definition (3) of \( S_B \)), as well as the need to provide an explicit IR regularization. The other complication is that the modified ghost sector should be included and it will couple in a non-trivial way to the matter sector. The \( \partial S/\partial g \) correlator is also very complicated. While the calculation involving the ghosts is quite lengthy and will not be attempted here, let us just highlight
the ambiguities mentioned by referring to the matter sector (we thus set \( g = 0 \)).

Following [15] (except for a slightly different choice of ghost number 1 operator \( g \) which is free from spurious IR divergences\(^8\)) one obtains from the definition (34)

\[
S = \left( -\sum_{m=1}^{\infty} \frac{1}{2} m (u_{m+1} - u_{m-1}) - u_m \frac{1}{m + u_m} - \frac{1}{2} \frac{u_0 + u_1}{1 + u_1} + a + 1 \right) \cdot Z, \tag{35}
\]

where the partition function \( Z \) is given by

\[
\log Z = -\sum_{k=1}^{\infty} \log \left( 1 + \frac{u_k}{k} \right) - a \tag{36}
\]

After inserting \( u_k = t k e^{-k \varepsilon} \) and isolating poles in \( \varepsilon \) we obtain

\[
\log Z = \frac{1}{\varepsilon} \sum_{n=1}^{\infty} \frac{(-1)^n}{n^2} t^n + \frac{1}{2} \log(1 + t) + O(\varepsilon) - a \tag{37}
\]

We can choose to renormalize \( a \) through \( a = a_R + \Delta a \):

\[
\Delta a = \frac{1}{\varepsilon} \sum_{n=1}^{\infty} \frac{(-1)^n}{n^2} t^n + \Delta a_{\text{finite}} \tag{38}
\]

The finite term should be fixed by some renormalization scheme prescription (a physical definition of \( a_R \)). Unfortunately, as already emphasized in [15, 19], we are lacking such a prescription in the context of BSFT. We thus put

\[
Z = \sqrt{1 + t} \cdot e^{-a_R - \Delta a_{\text{finite}}(a_R, t)} \tag{39}
\]

For \( D \) scalar fields we get \( \prod_{i=1}^{D} \sqrt{1 + t_i} \cdot e^{-a_R - \Delta a_{\text{finite}}}. \)

In an analogous way as in [15] the counterterm (38) also removes the divergence in the BSFT action (35), and we obtain

\[
S = \left( a_R + \Delta a_{\text{finite}}(t, a_R) + 1 - \frac{1}{2} \sum_{i=1}^{D} \frac{t_i}{1 + t_i} \right) \prod_{i=1}^{D} \sqrt{1 + t_i} e^{-a_R - \Delta a_{\text{finite}}(a_R, t)} \tag{40}
\]

From this expression it is clear that a well motivated renormalization prescription for \( \Delta a_{\text{finite}}(a_R, t) \) is needed in order to use \( S \) to study the renormalization group flow as a function of the parameters \( a, t \). Indeed, by looking at the behaviour of the BSFT action close to the \( a = t = 0 \) point, one can give arguments that \( \Delta a_{\text{finite}}(a_R, t) \) should be in general non-zero.

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\(^8\)We take \( O(\theta) = \frac{1}{2\pi} \epsilon_1(\theta)(X(\theta) - x_0) \int d\theta' u(\theta - \theta')(X(\theta') - x_0) \) where \( x_0 = \frac{1}{2\pi} \int d\theta X(\theta) \) is the zero mode. This choice of \( O(\theta) \) ensures that the ordinary open string with Neumann b.c. is a solution of the BSFT equations of motion \( \partial S/\partial a = \partial S/\partial t = 0 \) at \( t = a = 0 \).
Discussion

We have tried to identify the boundary interaction $S_B$ which corresponds to the tachyonic vacuum from the requirement that the correlation functions in the bulk agree with those of the closed string. The corresponding (open) string theory has indeed a decoupling of holomorphic and antiholomorphic modes and the associated boundary state can be identified with the closed string oscillator vacuum state. All correlation functions of bulk vertex operators coincide with those computed in an ordinary closed string CFT.

It would be desirable to obtain this boundary interaction from a renormalization group flow of the string action, i.e. to show explicitly that the BSFT action decreases when turning on the boundary interaction. As described above this requires a better understanding of the regularization of BSFT and the corresponding renormalization conditions, as well as a complete treatment of the ghost sector.

A better understanding of the ghost interaction and the relevant ghost zero-mode structure might also throw some light on one missing step in the above construction, namely how one obtains the closed string $S$-matrix elements from the open string theory with an off-shell boundary interaction. Although all closed string correlation functions of any vertex operators are exactly reproduced with (3) at $t = 1$, we have of course to integrate over the positions of the vertex operators to obtain the $S$-matrix elements and we encounter here a global mismatch between the integration regions of the closed and the open string which might be related to a proper definition of the coupling of on-shell closed strings to the off-shell open string theory. In view of the simplicity of the proposed boundary interaction, and its very close link with the closed string CFT’s we feel that these problems are worth further investigation.

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