

Constructing Open String Field Theory Solutions from Boundary Condition Changing Operators

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This is a short review of the results of the work [arXiv:1009.6185] with Michael Kiermaier and Yuji Okawa, which was presented at the *String Field Theory and Related Aspects 2010* conference. We construct an analytic solution of open string field theory which is given in terms of boundary condition changing (bcc) operator. The solution is valid for bcc operators with vanishing conformal weight such as those corresponding to regular marginal deformations of the background, but looks like a good starting point to look for solutions corresponding to more general deformations. Applications to the rolling tachyon profile are also presented.

§1. Introduction

Witten's open string field theory (OSFT)²⁾ was originally introduced as a candidate for a non-perturbative completion of perturbative open bosonic string theory. As such, it is expected to encode all possible consistent backgrounds of the perturbative theory as solutions to its classical equations of motion.

From the world-sheet perspective, consistent backgrounds for open strings are given by boundary conformal field theories (BCFT), and different open string backgrounds correspond to different boundary conditions on the world-sheet. A change of the boundary conditions can be described by the introduction of boundary condition changing (bcc) operators in the original BCFT. For example, the insertion of a pair of bcc operators $\sigma_L(a)$ and $\sigma_R(b)$ corresponds to a modification of the boundary conditions on the segment of the boundary stretched between the points a and b .

Therefore, the correspondence between conformal invariance in the world-sheet and the equations of motion in the spacetime theory is equivalent, in the case of open bosonic string theory, to the correspondence between pairs of bcc operators and solutions to the equations of motion of OSFT

$$Q\Psi + \Psi * \Psi = 0. \tag{1.1}$$

Here Ψ is an open string field of ghost number one, Q is the BRST operator, and the symbol $*$ denotes multiplication of string fields using Witten's star product.

Since the construction of an analytic solution to (1.1) by Schnabl,³⁾ an impressive amount of analytic results for string field theory has been obtained, and part of this correspondence has been clarified. In particular, OSFT solutions have been obtained which correspond to exactly marginal deformations of the original BCFT.⁴⁾⁻⁶⁾ Marginal deformations of the boundary conditions are those that gener-

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ate new BCFT that are continuously connected to the original one by a one parameter family of BCFTs. A primary field $V(t)$ of weight one in the matter sector generates an infinitesimal deformation of the BCFT, and conformal invariance is preserved to linear order in the deformation parameter which we denote by λ . When operator products of the marginal operator $V(t)$ are regular, finite deformations also preserve conformal invariance and thus the operator $V(t)$ generates a family of boundary conditions parameterized by λ . In this case, the change of boundary conditions on a segment $[a, b]$ can be implemented by

$$\sigma_L(a)\sigma_R(b) = \exp \left[\lambda \int_a^b dt V(t) \right]. \quad (1.2)$$

The bcc operators associated with such regular marginal deformations have vanishing conformal weights and satisfy

$$\lim_{\epsilon \rightarrow 0} \sigma_L(0)\sigma_R(\epsilon) = 1, \quad \lim_{\epsilon \rightarrow 0} \sigma_L(a)\sigma_R(b)\sigma_L(b+\epsilon)\sigma_R(c) = \sigma_L(a)\sigma_R(c). \quad (1.3)$$

All solutions constructed so far for these simple type of deformations rely heavily on expansions on the deformation parameter λ . Therefore, they do not seem to be the best starting point to look for solutions which correspond to general deformations, which are not necessarily continuously connected to the original BCFT, and for which no deformation parameter exists. Furthermore, these solutions rely on the existence of a BRST-closed operator such as $cV(t)$. For more general deformations, no natural choice of BRST-closed operator exists.

In this work we review a different solution recently constructed for regular marginal deformations, that can be expressed only in terms of bcc operators.¹⁾ No reference to the parameter λ or the deformation operator $V(t)$ is needed. Therefore, this solution looks like a natural starting point to look for solutions corresponding to general boundary condition deformations. In §2 we present the construction of the solution as a gauge-equivalent form of a previous result. In §3 we study the solution in the language of conformal field theory (CFT) and present an extremely simple form for its contraction $\langle \phi, \Psi \rangle$ with an arbitrary state ϕ in the Fock space. This determines the behavior of the component fields of the solution. The results obtained in this section are applied in §4, where we study the time dependence of the tachyon field in the specific example of the dynamical process of D-brane decay.

§2. A solution in terms of boundary condition changing operators

Our starting point is the analytic solution introduced in 4) and 5) for regular marginal deformations. The basic building blocks of this solution are the so-called wedge states W_α , defined implicitly by their BPZ inner product with a generic Fock-space state ϕ

$$\langle \phi, W_\alpha \rangle = \langle f \circ \phi(0) \rangle_{\mathcal{W}_\alpha}. \quad (2.1)$$

Here $\phi(\xi)$ is the operator associated to the state ϕ under the state-operator mapping, and $f \circ \phi(\xi)$ denotes its conformal transformation under the map

$$f(\xi) = \frac{2}{\pi} \arctan \xi. \quad (2.2)$$

The correlation function on the right-hand side of (2.1) is evaluated on the wedge surface \mathcal{W}_α with $\alpha \geq 0$, which is the semi-infinite strip on the upper half-plane of z between the vertical lines $\Re(z) = -\frac{1}{2}$ and $\Re(z) = \frac{1}{2} + \alpha$ with these lines identified by translation.

We will also make use of wedge-based states, i.e., wedge states with operator insertions. One of the operators we will find useful is the line integral of the b ghost operator defined by

$$\mathcal{B} = \int_{i\infty}^{-i\infty} \frac{dz}{2\pi i} b(z), \quad (2.3)$$

where we used the doubling trick. Its BRST transformation is given by

$$\mathcal{K} = \int_{i\infty}^{-i\infty} \frac{dz}{2\pi i} \oint \frac{dw}{2\pi i} j_B(w) b(z) = \int_{i\infty}^{-i\infty} \frac{dz}{2\pi i} T(z), \quad (2.4)$$

where j_B is the BRST current, T is the energy-momentum tensor, and the contour of the integral over w encircles z counterclockwise.

Let us denote by K , B , c and V , wedge-based states of zero width with insertions of \mathcal{K} , \mathcal{B} , $c(\frac{1}{2})$, and $V(\frac{1}{2})$ on the wedge surface \mathcal{W}_0 , respectively. For example, the state K is defined by the CFT correlators

$$\langle \phi, K \rangle = \langle f \circ \phi(0) \mathcal{K} \rangle_{\mathcal{W}_0}. \quad (2.5)$$

The states B , c and V are defined in a similar way. In fact, K is the string field that generates the wedge states W_α through the relation

$$W_\alpha = e^{\alpha K}. \quad (2.6)$$

The algebra satisfied by these states is given by

$$[K, B] = 0, \quad [B, V] = 0, \quad [c, V] = 0, \quad c^2 = 0, \quad B^2 = 0, \quad \{B, c\} = 1. \quad (2.7)$$

The BRST transformation Q acts on these states in the following way:

$$QB = K, \quad QK = 0, \quad Qc = cKc, \quad QV = [K, cV]. \quad (2.8)$$

It follows that $Q(cV) = 0$, which expresses the marginality of the operator V .

With these ingredients at hand, we can now present a solution to the OSFT equations of motion (1.1) corresponding to any marginal deformation operator $V(t)$ with regular operator products:^{(4), (5), (7)}

$$\Psi = f(K) \lambda cV \left[1 - B \frac{f(K)^2 - 1}{K} \lambda cV \right]^{-1} f(K). \quad (2.9)$$

To avoid wedge states W_α with negative width α , we require the arbitrary function $f(K)$ to take the following form:

$$f(K) = \int_0^\infty dt \tilde{f}(t) e^{tK}. \quad (2.10)$$

It can be shown that Ψ represents a solution of the equations of motion (1.1) by using the relations (2.7) and (2.8), and no reference to explicit CFT correlators is needed.

Our goal now is to find a function $f(K)$ such that the solution (2.9) can be expressed entirely in terms of bcc operators. In order to do so, let us study how a wedge state with modified boundary conditions looks like in the algebraic language presented above. An insertion of

$$\exp \left[\lambda \int_a^b dt V(t) \right] = 1 + \lambda \int_a^b dt_1 V(t_1) + \lambda^2 \int_a^b dt_1 \int_{t_1}^b dt_2 V(t_1) V(t_2) + \dots \quad (2.11)$$

corresponds to

$$\begin{aligned} e^{\alpha K} + \lambda \int_0^\alpha dt_1 e^{t_1 K} V e^{(\alpha-t_1)K} + \lambda^2 \int_0^\alpha dt_1 \int_{t_1}^\alpha dt_2 e^{t_1 K} V e^{(t_2-t_1)K} V e^{(\alpha-t_2)K} + \dots \\ = e^{\alpha(K+\lambda V)} \end{aligned} \quad (2.12)$$

with $\alpha = b - a$. This means that $K + \lambda V$ is the string field that generates wedge states with modified boundary conditions. Therefore, we can trade any superposition $h(K + \lambda V)$ of deformed wedge states by

$$\dots h(K + \lambda V) \dots = \dots \sigma_L h(K) \sigma_R \dots, \quad (2.13)$$

where the dots \dots represent arbitrary wedge-based states with the boundary conditions of the undeformed BCFT. Furthermore, considering the BRST transformations of (2.13), we can see that

$$\begin{aligned} \dots h(K + \lambda V)(\lambda cV) \dots &= \dots \sigma_L h(K) (Q \sigma_R) \dots, \\ \dots (-\lambda cV) h(K + \lambda V) \dots &= \dots (Q \sigma_L) h(K) \sigma_R \dots \end{aligned} \quad (2.14)$$

We learn that any time the state V appears in the form $h(K + \lambda V)$, $h(K + \lambda V)(\lambda cV)$ or $(\lambda cV) h(K + \lambda V)$, it can be substituted by bcc operators.^{*})

Now we want to find a function $f(K)$ such that any appearance of the state V in the solution (2.9) can be replaced by bcc operators and their BRST transformations. It can be shown¹⁾ that there is a unique choice (up to rescaling of K) if one allows only for a finite number of insertions of bcc operators on the boundary, which is given by

$$f(K) = \frac{1}{\sqrt{1-K}} = \int_0^\infty ds \frac{e^{-s}}{\sqrt{\pi s}} e^{sK}. \quad (2.15)$$

^{*}) One must ensure that $h(K)$ does not contain wedge states of negative width, i.e. that it can be expressed as $h(x) = \int_0^\infty d\alpha \tilde{h}(\alpha) e^{\alpha x}$.

Substituting (2.15) into (2.9) we obtain, after some massage, the desired solution

$$\boxed{\begin{aligned}\Psi &= \frac{1}{\sqrt{1-K}} \lambda cV \frac{1}{1-K-\lambda V} (1-K) Bc \frac{1}{\sqrt{1-K}} \\ &= -\frac{1}{\sqrt{1-K}} (Q\sigma_L) \frac{1}{1-K} \sigma_R (1-K) Bc \frac{1}{\sqrt{1-K}}.\end{aligned}} \quad (2.16)$$

It is interesting to note that exactly the same function $f(K)$ was used before⁸⁾ to obtain a simple form of the analytic solution corresponding to tachyon condensation.⁹⁾

Since (2.16) was constructed as a special case of the general solution (2.9), it automatically satisfies the equations of motion. It also satisfies the reality condition on the string field,¹⁰⁾ but this is not manifest in its present form. We can, however, rewrite (2.16) as

$$\boxed{\Psi = -\frac{1}{\sqrt{1-K}} (Q\sigma_L) \sigma_R \frac{1}{\sqrt{1-K}} - \frac{1}{\sqrt{1-K}} (Q\sigma_L) \frac{B}{1-K} (Q\sigma_R) \frac{1}{\sqrt{1-K}}.} \quad (2.17)$$

This form is manifestly symmetric when we reverse the order of multiplication of string fields and thus satisfies the reality condition,¹⁰⁾ which guarantees that the string field theory action is real.

Although we arrived at (2.17) from an expression that contained the marginal parameter λ and operator V explicitly, it is now written only in terms of K , B , σ_L , and σ_R . This is a solution to the equation of motion for any choice of bcc operators σ_L , σ_R in the matter sector that satisfy the operator products (1.3). More generally, (2.17) satisfies the equation of motion for any choice of three states B , σ_L , and σ_R satisfying the relations

$$B^2 = 0, \quad [B, \sigma_L] = 0, \quad [B, \sigma_R] = 0, \quad \sigma_L \sigma_R = 1, \quad \sigma_R \sigma_L = 1. \quad (2.18)$$

$QB = K$ can be thought of as a definition of the string field K . Equation (2.18) and their BRST transformations are the only ingredients needed to prove that (2.17) satisfies the equations of motion.

§3. CFT language and universal coefficients

In this section we translate the solution (2.17) to the CFT language, and study the behavior of its component fields. These are specified by giving the inner products $\langle \phi, \Psi \rangle$ for arbitrary states ϕ in the Fock space. We will exemplify the computation with the case of a state ϕ of the form $\phi = -c\partial\phi_m$, where ϕ_m is a matter primary field of weight h .

We can use the fact that the BRST transformations of bcc operators with vanishing conformal weight are given by $Q\sigma_L = c\partial\sigma_L$ and $Q\sigma_R = c\partial\sigma_R$ to write

$$\Psi = -\frac{1}{\sqrt{1-K}} (c\partial\sigma_L) \sigma_R \frac{1}{\sqrt{1-K}} - \frac{1}{\sqrt{1-K}} (c\partial\sigma_L) \frac{B}{1-K} (c\partial\sigma_R) \frac{1}{\sqrt{1-K}}. \quad (3.1)$$

The first term in (3.1) is a superposition of wedge states with a single insertion of $(c\partial\sigma_L)\sigma_R = -\lambda cV$, and its inner product with ϕ vanishes for a weight $h \neq 1$. Therefore, for $h > 1$, we can focus on the second term, and with the help of the Laplace transformations of $1/\sqrt{1-K}$ and $1/(1-K)$ we can write

$$\langle \phi, \Psi \rangle = -\int_{\frac{1}{2}}^{\infty} dx \int_0^{\infty} ds \int_{\frac{1}{2}}^{\infty} dy \frac{e^{1-x-s-y}}{\pi \sqrt{(x-\frac{1}{2})(y-\frac{1}{2})}} \left\langle f \circ \phi(0) c\partial\sigma_L(x) \mathcal{B} c\partial\sigma_R(x+s) \right\rangle_{\mathcal{W}_{L-1}}, \quad (3.2)$$

where we have defined $L = x + s + y$. Since we have chosen the state ϕ in a basis with factorized matter and ghost sectors, the correlator in (3.2) can be separated into

$$\begin{aligned} & \langle f \circ \phi(0) c\partial\sigma_L(x) \mathcal{B} c\partial\sigma_R(x+s) \rangle_{\mathcal{W}_{L-1}} \\ &= -\frac{\pi}{2} \langle c\partial c(0) c(x) \mathcal{B} c(x+s) \rangle_{\mathcal{W}_{L-1,g}} \times \partial_a \partial_b \langle f \circ \phi_m(0) \sigma_L(a) \sigma_R(b) \rangle_{\mathcal{W}_{L-1,m}} \Big|_{\substack{a=x \\ b=x+s}}, \end{aligned} \quad (3.3)$$

where $\langle \dots \rangle_g$ and $\langle \dots \rangle_m$ stand for ghost and matter correlators, respectively. Each correlator can then be computed separately by conformally mapping the wedge surface \mathcal{W}_{L-1} to the upper half plane (UHP). The results are

$$\langle c\partial c(0) c(x) \mathcal{B} c(x+s) \rangle_{\mathcal{W}_{L-1,g}} = -\frac{L^2}{\pi^3} \left[\theta_y \sin^2 \theta_x + \theta_x \sin^2 \theta_y - \sin \theta_x \sin \theta_s \sin \theta_y \right] \quad (3.4)$$

for the ghost sector and

$$\partial_a \partial_b \langle f \circ \phi_m(0) \sigma_L(a) \sigma_R(b) \rangle_{\mathcal{W}_{L-1,m}} \Big|_{\substack{a=x \\ b=x+s}} = -C_\phi \frac{\pi^2 h(h-1)}{L^2 \sin^2 \theta_s} \left| \frac{2 \sin \theta_s}{L \sin \theta_x \sin \theta_y} \right|^h \quad (3.5)$$

for the matter sector. Here we have defined

$$\theta_\ell = \frac{\ell}{L} \pi \quad (3.6)$$

and the matter three-point function

$$C_\phi = \langle \phi_m(0) \sigma_L(1) \sigma_R(\infty) \rangle_{\text{UHP},m}. \quad (3.7)$$

Combining these results we can obtain a very simple factorized form for the inner products

$$\boxed{\langle \phi, \Psi \rangle = C_\phi g(h)} \quad (3.8)$$

with

$$\begin{aligned} g(h) &= \frac{h(h-1)}{2\pi} \int_{\frac{1}{2}}^{\infty} dx \int_0^{\infty} ds \int_{\frac{1}{2}}^{\infty} dy \frac{e^{1-x-s-y}}{\sin^2 \theta_s \sqrt{(x-\frac{1}{2})(y-\frac{1}{2})}} \left| \frac{2 \sin \theta_s}{L \sin \theta_x \sin \theta_y} \right|^h \\ &\quad \times \left[\theta_y \sin^2 \theta_x + \theta_x \sin^2 \theta_y - \sin \theta_x \sin \theta_s \sin \theta_y \right]. \end{aligned} \quad (3.9)$$

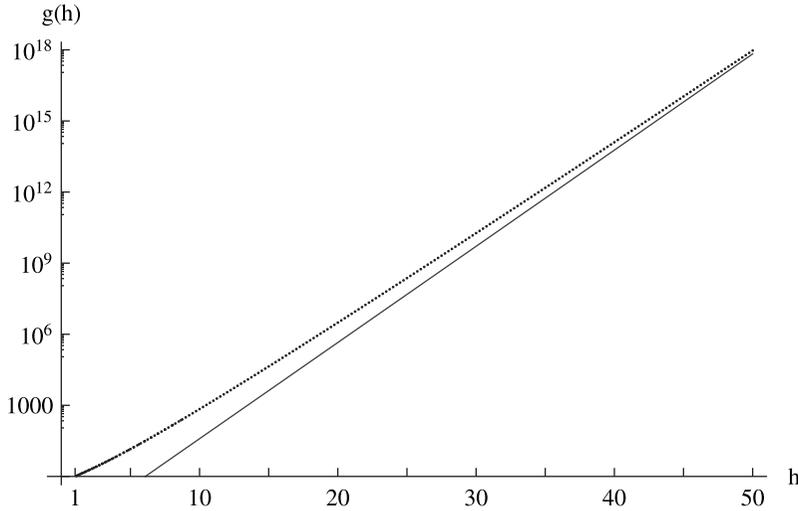


Fig. 1. The black dots represent numerical evaluations of the function $g(h)$ given in Eq. (3.9). The solid line is the function $a(8/\pi)^h$, with the coefficient a fitted to numerical evaluations of $g(h)$ in the range $1.25 \leq h \leq 60$.

Equation (3.8) shows the very remarkable factorization properties of the solution (2.17). Notice that all the information about the boundary deformation is encoded in the three point function C_ϕ , while the function $g(h)$ depends exclusively on the conformal weight of ϕ , and is a universal function that can be computed once and for all.*) From (3.9) we can see that, for $h \gg 1$, the dominant contribution to $g(h)$ comes from the factor $|\dots|^h$ which tends to its maximum $8/\pi$ when $x = y = \frac{1}{2}$ and $s \rightarrow \infty$. We conclude that

$$g(h) \sim \left(\frac{8}{\pi}\right)^h \quad \text{for } h \gg 1. \quad (3.10)$$

Figure 1 gives a plot of $g(h)$ on a logarithmic scale, together with the asymptote of its large- h behavior.

§4. Application to the rolling tachyon profile

The dynamical process of D-brane decay is described in open string field theory by the exactly marginal operator⁽¹¹⁾⁻¹³⁾

$$V(t) = e^{\frac{1}{\sqrt{\alpha'}} X^0}(t). \quad (4.1)$$

In this section, we will use the solution (2.17) associated to this deformation to study the behavior of its tachyon component $T(x)$ as the brane decays. A useful way to

*) The form of the function $g(h)$ depends also on the ghost sector of the state ϕ . For other choices different than $c\bar{c}$, other universal functions can be easily obtained in a similar way.

parametrize the time dependence of this field is^{4),5)}

$$T(x^0) = \sum_{n=1}^{\infty} \beta^{(n)} \lambda^n e^{\frac{1}{\sqrt{\alpha'}} n x^0}. \quad (4.2)$$

The absolute value of the deformation parameter λ can be rescaled by a shift in time and therefore does not have physical meaning. Its sign specifies the direction towards which the tachyon starts to roll. In our conventions, $\lambda < 0$ corresponds to the situation in which the tachyons rolls to the tachyon minimum. The coefficients $\beta^{(n)}$ determine completely the tachyon profile and can be obtained from the string field solution Ψ by studying^{*})

$$\beta^{(n)} \lambda^n = \langle \phi^{(n)}, \Psi \rangle, \quad (4.3)$$

where

$$\phi^{(n)}(t) = -c\partial c \phi_m^{(n)}(t), \quad \phi_m^{(n)}(t) = e^{-\frac{n}{\sqrt{\alpha'}} X^0}(t). \quad (4.4)$$

We note now that the states $\phi^{(n)}$ have conformal weight $h = n^2$, and take the same form as that of the states used in the previous section to study the solution (2.17). Therefore, we can borrow the results from that section (mainly the function $g(h)$ of (3.9)), and the computation of the coefficients $\beta^{(n)}$ reduces to the study of the three-point functions $C_{\phi^{(n)}}$. These can be determined by standard CFT methods and are given by¹⁾

$$C_{\phi^{(n)}} = \langle \phi_m^{(n)}(0) \sigma_L(1) \sigma_R(\infty) \rangle_{\text{UHP, m}} = \lambda^n \prod_{i=0}^{n-1} \frac{i!^4}{(2i+1)!(2i)!}. \quad (4.5)$$

Combining this result with the explicit evaluation of the $g(h)$ function (3.9), we can evaluate the coefficients $\beta^{(n)}$ for arbitrary n . In particular, we can study their large- n behavior, which determines the late time x_0 evolution of the tachyon field $T(x_0)$. Using Sterling's approximation one can see that

$$C_{\phi^{(n)}} = \lambda^n \exp\left[-2n^2 \log 2 + \mathcal{O}(n \log n)\right]. \quad (4.6)$$

This result, together with the asymptotic behavior in (3.10) leads to the large- n expression for the coefficients

$$\beta^{(n)} = \exp\left[-\gamma n^2 + \mathcal{O}(n \log n)\right] \quad \text{with } \gamma = \log \frac{\pi}{2}. \quad (4.7)$$

The convergence of the series (4.2) for any time x_0 is manifest from this expression. The simplicity of the solution (2.17) has allowed us to confirm analytically this result which had been previously suggested from numerical results.

The profile of the tachyon field exhibits a wild oscillatory behavior at late time already encountered in the early analysis by level truncation,^{14),15)} and later confirmed by the analytic results in Schnabl gauge.^{4),5)} This somewhat surprising behavior is

^{*}) Henceforth we will assume that all correlators have been normalized by the spacetime volume.

related to the appearance of the exponential suppression coefficient γ of Eq. (4.7). It would be interesting to find a solution for which the oscillations were absent, i.e., for which $\gamma = 0$. Having an analytic evaluation of the coefficient γ for the solution (2.17) described in this work, one can use reparameterizations to generate new solutions and study the late time behavior of the tachyon field. This can be done explicitly¹⁾ for the family of solutions associated to special projectors.¹⁶⁾ The result is that the coefficient γ is strictly positive for this kind of solutions, and thus all of them exhibit the highly oscillatory behavior mentioned before. It would be interesting to find out whether this characteristic is intrinsic to the rolling tachyon profile, and therefore exhibited by any solution of OSFT associated to the marginal deformation (4.1).

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