Exact Solutions in Open Bosonic String Field Theory and Marginal Deformation in CFT

J. Klusoň *

Institute of Theoretical Physics, University of Stockholm, SCFAB
SE- 106 91 Stockholm, Sweden

and

Institutionen för teoretisk fysik
BOX 803, SE- 751 08 Uppsala, Sweden
E-mail: josef.kluson@teorfys.uu.se

ABSTRACT: In this paper we continue our study of the exact solution in open bosonic string field theory. We present new solution in the string field theory defined on the background corresponding to the boundary conformal field theory describing D25-brane. Then we will study the fluctuation modes around this solution and we determine their basic properties from the linearised equation of motion of the string field theory defined above the classical solution.

KEYWORDS: String field theory

*On leave from Masaryk University, Brno
1. Introduction

The construction of time dependent classical solutions that describe the decay of an unstable D-brane as the tachyon rolls down toward the minimum of the potential has been given in series of seminal papers by A. Sen [1, 2, 3, 4, 5]. The study of time dependent solutions representing of tachyons in $p$-adic string theory was performed in [6] and from different point of views in [7, 8, 9, 10, 11, 12]. Application of the tachyon matter in cosmology has been discussed in [13, 14, 15, 16, 17, 18, 19, 20, 21, 22, 23, 24, 25, 26, 27, 28, 29, 30, 31, 32, 33, 34, 35, 36]. In [1, 2, 3, 4] the rolling tachyon solution was mainly discussed from the point of view of the world-sheet boundary conformal field theory ($BCFT$). More precisely, this solution was constructed by perturbing $BCFT$ describing the original D-brane by an exactly marginal perturbation. Since the Wick rotated version is exactly solvable [51, 52, 53, 54, 55] we can obtain lot of exact results considering the rolling of the tachyon to its vacuum value. More precisely, the strength $\lambda$ of the perturbation labels the initial value of the tachyon $T$. The corresponding boundary state gives information about the time evolution of various closed string sources, for recent discussion, see [5, 12]. In particular, it was shown that at $\lambda = 1/2$ which corresponds to placing the tachyon at the minimum of the potential, the whole boundary state vanishes at all time as conjectured in [37, 38, 39, 40, 41, 42]. It can be also shown that the rolling tachyon solution has deep connection with spacelike branes [56, 57, 58, 59]. In fact, it was stressed recently [12, 60] that this picture is greatly modified for nonzero string coupling constant $g_s \neq 0$.\footnote{For review of the tachyon condensation in string field theory and extensive list of references, see [43, 44, 45].}
In this paper we will try to present the modest contribution to this very reach and fascinating subject. We will continue our study of the exact solutions of open bosonic string field theory (SFT) and their possible relations to the time dependent solutions\(^3\). In some sense this paper can be regarded as an application of the general analysis performed in \([34]\), for similar approach of the searching of the exact solutions in SFT, see \([33, 64, 63, 60, 77, 78]\). More precisely, we start with the open bosonic string field theory defined on the background \(BCFT\) corresponding to space-time filling D25-brane. Then we find such a solution of the SFT equation of motion that when we expand string field around this solution the new BRST operator will have form that we could expect for the string field theory defined around the new conformal field theory background \(BCFT'\) that is related to the original one by marginal deformation in the original \(BCFT\). It is important to stress that we do not claim that the SFT for fluctuation modes is SFT formulated around the background \(BCFT''\). In particular, the fluctuation modes still belong to the Hilbert space of the original \(BCFT\) and the correlation functions are also calculated using the original \(BCFT\). In order to really show that the SFT for fluctuation modes corresponds to the SFT formulated around the new \(BCFT''\) we should present such a careful analysis as was done in \([78, 79, 80, 81, 82]\) in case of closed string field theory. In fact this approach is currently under study and we hope to return to this question in future.

In order to study the fluctuation spectrum around the classical solution we will proceed in different way. Using the form of the shifted BRST operator we will propose such a form of the string field fluctuation modes around the classical solution that will allow straightforward analysis of the linearised string field theory equation of motion which these modes have to obey. According to this analysis we will present arguments that these modes correspond to the states of the \(BCFT''\) with the marginal perturbation inserted on the boundary that was studied in \([51, 52, 53, 54, 55]\) (See also \([1, 2, 3, 4, 5]\) for recent application of this \(BCFT''\) in the case of the rolling tachyon solution.).

However we must mention one remarkable fact. As we will see explicitly in the next section, our solution has pure gauge form which could then suggest that our approach does not lead to any new physical situation. But this conclusion seems to be with contradiction with our explicit calculations that will show that the original BRST operator is really deformed by our solutions and has exactly the same form as we could expect from deformed \(BCFT\) theories studied in \([51, 52, 53, 54, 55]\). In order to resolve this puzzle we will argue that there is a difference between large gauge transformations that really reflect the symmetry of the gauge theory, which in our case is string field theory, and under which the string field theory theory action is

\(^3\)Alternative version of the exact solution of the open bosonic string field theory was given recently in \([31]\).
not generally invariant and under small gauge transformations that leave the action invariant and that really does not give new physical configurations.

This paper is organized as follows. In the next section we will calculate an exact solution of the SFT equation of motion in the spirit of our general analysis given in [65]. Then we perform an expansion in the SFT action around this solution and we will study the fluctuation modes around the new solution that will obey the linearised equation of motion. We will show that the spectrum of the modes is the same as in [1, 2, 3, 4, 5] in the case of the rolling tachyon solution or as in [51, 52, 53, 54, 55] in case of the spatial dependent solution. In the conclusion we will outline our results and we will discuss other problems that deserve further study.

2. Marginal solution

In this section we will find exact solution of the SFT equation of motion that can be directly related to the rolling tachyon solution, or more generally to BCFT studied in [51, 52, 53, 54, 55].

To begin with, we firstly review basic facts about bosonic string field theory, following mainly [43, 44, 45]. Gauge invariant string field theory is described with the full Hilbert space of the first quantized open string including $b, c$ ghost fields subject to the condition that the states must carry ghost number one, where $b$ has ghost number $-1$, $c$ has ghost number 1 and $SL(2, C)$ invariant vacuum $|0\rangle$ carries ghost number 0. We denote $\mathcal{H}$ the subspace of the full Hilbert space carrying ghost number 1. Any state in $\mathcal{H}$ will be denoted as $|\Phi\rangle$ and corresponding vertex operator $\Phi$ is the vertex operator that creates state $|\Phi\rangle$ out of the vacuum state $|0\rangle$

$$|\Phi\rangle = \Phi |0\rangle .$$

(2.1)

Since we are dealing with open string theory, the vertex operators should be put on the boundary of the world-sheet. The string field theory action is given

$$S = -\frac{1}{g_0^2} \left( \frac{1}{2\alpha'} \int I \circ \Phi(0) Q \Phi(0) + \frac{1}{3} \langle f_1 \circ \Phi(0) f_2 \circ \Phi(0) f_3 \circ \Phi(0) \rangle \right) ,$$

(2.2)

where $g_0$ is open string coupling constant, $Q$ is BRST operator and $<>$ denotes correlation function in the combined matter ghost conformal field theory. $I, f_1, f_2, f_3$ are conformal mapping exact form of which is reviewed in [43] and $f_i \circ \Phi(0)$ denotes the conformal transformation of $\Phi(0)$ by $f_i$. For example, for $\Phi$ a primary field of dimension $h$, then $f_i \circ \Phi(0) = (f'_i(0))^h \Phi(f_i(0))$. In the abstract language pioneered in [69] the open string field theory action (2.2) is

$$S = -\frac{1}{g_0^2} \left( \frac{1}{2\alpha'} \int \Phi \ast Q \Phi + \frac{1}{3} \int \Phi \ast \Phi \ast \Phi \right)$$

(2.3)
from which we immediately get an equation of motion
\[ \frac{1}{\alpha'} Q\Phi_0 + \Phi_0 \star \Phi_0 = 0 . \] (2.4)

It is easy to see that the string field in the form
\[ \Phi_0 = e^{-K_L(I)} \star \frac{1}{\alpha'} Q(e^{K_L(I)}) \] (2.5)
is solution of (2.4) for any ghost number zero operator $K_L$ acting on the string field theory algebra $\star$ identity $I$ which is ghost number zero field that obeys \[ I \star X = X \star I = X, \] (2.6)
for any string field $X$ \footnote{For recent study of the identity element $I$, see \cite{70, 71, 72, 73, 74}.}. The operator $K$ in (2.5) can be written as
\[ K = \frac{A}{2\pi i} \int_C dz k(z) + \frac{B}{2\pi i} \int_C d\overline{\zeta} \overline{k}(\overline{\zeta}), \] (2.7)
where $k(z), \overline{k}(\overline{\zeta})$ are holomorphic and anti holomorphic fields of conformal dimensions $(1, 0), (0, 1)$ respectively that are defined in the upper-half plane $\text{Im} \ z \geq 0$. The field $k(z)$ transforms under conformal transformation $z \to f(z)$ as
\[ U_f k(z)U_f^{-1} = \frac{df(z)}{dz} k(f(z)) \equiv f'(z) k(f(z)) , \] (2.8)
with the same expression for anti holomorphic field $\overline{k}(\overline{\zeta})$. In general, the integral in (2.7) depends on the integration contour $C$ which we choose as a semicircle in the upper half-plane $C : \{ z; z = -e^{-i\sigma+i\tau}, \sigma \in (0, \pi/2) \}$ and we also choose $\tau = 0$. The fact that $\sigma \in (0, \pi/2)$ means that we perform integration over the left side of the string so that we label corresponding operator with the subscript $L$. Operators without any subscript correspond to the integration of any local density over the whole string so that $\sigma \in (0, \pi)$ and finally operator corresponding to the integration over right side of the string will be labeled by the subscript $R$ with $\sigma \in (\pi/2, 0)$.

Now we will study the transformation properties of $K$ (2.7) under general conformal transformations. Generally $K$ is not invariant under conformal transformation under which it transforms as
\[ U_f KU_f^{-1} = \frac{A}{2\pi i} \int_C dw f'(w) k(f(w)) + \frac{B}{2\pi i} \int_C d\overline{w} f'(\overline{w}) \overline{k}(f(\overline{w})) = \frac{A}{2\pi i} \int_{f(C)} df k(f) + \frac{B}{2\pi i} \int_{f(C)} d\overline{f} \overline{k}(\overline{f}) \neq K , \] (2.9)
which is not equal to \( K \) since, as we have argued above, the integration in \( K \) explicitly depends on the integration contour \( C \). However we can easily prove that for the BCF T with standard Neumann boundary conditions

\[ k(z) = \tilde{k}(\bar{z}), \text{ Im } z = 0 \]  

(2.10)

the operator \( K \) with \( A = -B \) will be invariant under conformal transformations since then the definition of \( K \) does not depend on the integration contour \( C \). To see this let us consider following integral over closed curve \( C \) in the upper half-plane

\[ F(C) = \frac{1}{2\pi i} \oint_C dz k(z) - \frac{1}{2\pi i} \oint_C d\bar{z} \tilde{k}(\bar{z}). \]  

(2.11)

This integral is equal to zero thanks to the holomorphicity and antiholomorphicity of \( k, \tilde{k} \) respectively. Let us consider \( C \) as

\[ C = C_1 + C_{r1} + C_{r2} + C'_2, \]

\[ C_1 = \{ z; z = -r_1 e^{-i\sigma}, \sigma \in (0, \pi) \}, \]

\[ C_{r1} = \{ z \in (r_1, r_2) \text{, Im } z = 0 \text{, } r_1 > r_2 \}, \]

\[ C'_2 = \{ z; z = -r_2 e^{-i\sigma}, \sigma \in (\pi, 0) \}, \]

\[ C_{r2} = \{ z \in (r_2, r_1) \text{, Im } z = 0 \}. \]  

(2.12)

Then it is easy to see that the integrals of \( k, \tilde{k} \) performed over \( C_{r1}, C_{r2} \) cancel each other since \( z = \overline{z}, \ t(z) = \overline{\tilde{k}(\bar{z})} \). Consequently we get

\[ F(C) = 0 = K(C_1) + K(C'_2) = K(C_1) - K(C_2) \Rightarrow K(C_1) = K(C_2), C'_2 = -C_1 \]  

(2.13)

so that \( K \) does not depend on the integration contour \( C \) in the upper-half plane. For such an operator (2.9) gives

\[ U_f K U_f^{-1} = K. \]  

(2.14)

Now we prove that the operator \( K \) defined in (2.7) obeys \( \delta \)

\[ K_L(A) \ast B = -A \ast K_R(B) \]  

(2.15)

for any string fields \( A, B \). We follow mainly \( \delta \) which is based on fundamental papers \( \delta, \delta \). The \( N \)-string vertex of a midpoint interaction is defined by gluing the boundaries \( |w_i| = 1, (i = 1, \ldots, N) \) of \( N \) unit half disks in the upper-half plane \( \text{Im } z \geq 0 \) with the identifications

\[ w_i w_{i+1} = -1, \text{ for } |w_i| = 1, \text{ Re } w_i \leq 1. \]  

(2.16)

Under this identification the primary field of conformal dimension \( h \) transforms as

\[ U_f \phi(w_{i+1}) U_f^{-1} = \left( \frac{df(w_{i+1})}{dw_{i+1}} \right)^h \phi(f(w_{i+1})), f(w_{i+1}) = w_i \]  

(2.17)
so that for \( \phi(w) \) that is primary field of conformal dimension 1 we get

\[
U_f dw_{i+1} \phi(w_{i+1}) U_f^{-1} = dw_{i+1} \left( \frac{dw_i}{dw_{i+1}} \right) \phi(w_i) = dw_i \phi(w_i)
\]  

(2.18)

and similar for complex conjugate \( \overline{w}, \overline{w}_{i+1} \). Using this result and the fact that under identification (2.16) we change the range of integration from \( \sigma \in (0, \pi/2) \) to \( (\pi, \pi/2) \) we immediately get (2.15). For the spatial case \( A = B = I \) (2.15) gives

\[
K_L(I) = K_L(I) \ast I = -I \ast K_R(I) = -K_R(I) .
\]  

(2.19)

It is important to stress that we have got this result (2.19) without presumption that \( K \) is the derivation of the star algebra

\[
K(A \ast B) = K(A) \ast B + A \ast K(B) .
\]  

(2.20)

On the other hand it is well known that the BRST operator obeys

\[
Q(A \ast B) = Q(A) \ast B + (-1)^{|A|} A \ast Q(B)
\]  

(2.21)

which for \( A = I \) implies \( Q(I) = Q_L(I) + Q_R(I) = 0 \) [62].

In the following we restrict ourselves to the case \( B = 0 \) in (2.7) since we will not demand invariance of \( K \) under conformal transformation. In fact in our further discussion we will consider two particular examples of \( k \)

\[
K_t = \frac{A}{2\pi i} \int_C dz k_t(z) = \frac{A}{2\pi i} \int_C dz \exp \left( \frac{2}{\sqrt{\alpha'}} X^0_L(z) \right) ,
\]  

(2.22)

\[
K_s = \frac{A}{2\pi i} \int_C dz k_s(z) = \frac{A}{2\pi i} \int_C dz \exp \left( i \frac{2}{\sqrt{\alpha'}} X^1_L(z) \right) .
\]  

(2.23)

Now the fields \( k_t(z), k_s(z) \) have conformal dimension equal to 1 as can be seen from the following OPE

\[
T(z) k_{s,t}(w) = \frac{1}{(z - w)^2} k_{s,t}(w) + \frac{1}{z - w} \partial k_{s,t}(w) ,
\]  

(2.24)

where in our calculation we use convention from the very nice review [K3].

\[
T_m(z) = -\frac{1}{\alpha'} \partial _\zeta X^*_{L}(z) \partial _\zeta X^\nu_{L}(z) \eta_{\mu\nu} ,
\]

\[
\tilde{T}_m(\overline{z}) = -\frac{1}{\alpha'} \partial _\zeta X^*_{R}(\overline{z}) \partial _\zeta X^\nu_{R}(\overline{z}) \eta_{\mu\nu} ,
\]

\[
X^\mu_{L}(z) X^\nu_{L}(w) \sim -\frac{\alpha'}{2} \eta^{\mu\nu} \ln(z - w) ,
\]

\[
X^\mu_{R}(\overline{z}) X^\nu_{R}(\overline{w}) \sim -\frac{\alpha'}{2} \eta^{\mu\nu} \ln(\overline{z} - \overline{w}) ,
\]

\[
X^\mu_{L}(z) X^\nu_{R}(\overline{w}) \sim -\frac{\alpha'}{2} \eta^{\mu\nu} \ln(z - \overline{w})
\]  

(2.25)

\footnote{In the following we will often omit the subscripts \( t, s \) in situations, where we will not need explicit form of \( k \).}
with the BRST operator

\[ Q = \frac{1}{2\pi i} \int_C dz j_B(z) - \frac{1}{2\pi i} \int_C d\bar{z} \tilde{j}_B(\bar{z}) , \]

\[ j_B(z) = c(z) \left[ T_m(z) + \frac{1}{2} T_{gh}(z) \right] , \]

\[ \tilde{j}_B(\bar{z}) = \tilde{c}(\bar{z}) \left[ \tilde{T}_m(\bar{z}) + \frac{1}{2} \tilde{T}_{gh}(\bar{z}) \right] , \]

(2.26)

where \( j(z) \) is holomorphic and \( \tilde{j}(\bar{z}) \) is anti holomorphic current and where \( T_{ghost} \) is the stress energy tensor for the ghost field. In what follows we will not need to know the explicit form of the ghost contribution.

In summary, we have shown in the previous part that there are exact solutions of the SFT equation of motion that are based on the existence of primary fields of conformal dimension 1. Using these fields we have constructed operators in BCFT that are invariant under conformal transformations. This implies, since these operators do not contain ghost contribution, that their commutator with the BRST charge is nonzero. This important property the fact that these operators obey (2.13) leads to the modification of the BRST operator in the original BCFT theory to the new form of the BRST operator \( Q' \). It is important to stress that this operator \( Q' \) is still defined in the original BCFT theory, however the form of the action for fluctuation modes suggests that the string field action formulated around the classical solution corresponds to the string field theory action formulated around new BCFT' that arises from the original BCFT by marginal interaction inserted on the boundary of the world-sheet.

To see this more explicitly, let us expand the string field \( \Phi \) in (2.3) around the solution (2.5) as

\[ \Phi = \Phi_0 + \Psi . \]

(2.27)

and insert it in (2.3). Then we obtain an action for the fluctuation field \( \Psi \) in the same form as the original one (2.3)

\[ S = -\frac{1}{g_0^2} \left( \frac{1}{2\alpha'} \int \Psi \star Q' \Psi + \frac{1}{3} \int \Psi \star \Psi \star \Psi \right) , \]

(2.28)

where the new BRST operator \( Q' \) is introduced

\[ Q'(X) = Q(X) + \Phi_0 \star X - (-1)^{|X|} X \star \Phi_0 . \]

(2.29)

In order to obtain the new form of the BRST operator (2.29) we will follow the calculation outlined in [65] and which we briefly review below. We start with the function

\[ F(t) = \frac{1}{\alpha'} e^{-K_L(t)} \star Q(e^{K_L(t)}) , F(1) = \Phi_0 , F(0) = 0 \]

(2.30)
and perform Taylor expansion around the point $t = 1$

$$\Phi_0 = F(1) = F(0) + \sum_{n=1}^{\infty} \frac{1}{n!} \frac{d^n F}{dt^n}(0),$$  \hspace{1cm} (2.31)

where

$$\frac{dF}{dt} = \frac{1}{\alpha'} e^{-K_L(\mathcal{I})t} \ast [Q, K]_{L}(\mathcal{I}) \ast e^{K_L(\mathcal{I})t},$$

$$\frac{d^2F}{dt^2}(0) = \frac{1}{\alpha'}[Q, K]_{L}(\mathcal{I}) \equiv D_L(\mathcal{I}),$$

$$\frac{d^3F}{dt^3}(0) = [K_L, [K, D_L]] \ldots \frac{d^n F}{d^n t}(0) = \underbrace{[K, [K, \ldots, [Q, K]]]}_{L} L(\mathcal{I})$$

and consequently

$$\Phi_0 = \frac{1}{\alpha'} \sum_{n=1}^{\infty} \frac{1}{n!} \underbrace{[K, [K, \ldots, [Q, K]]]}_{L}(\mathcal{I}) \equiv D_L(\mathcal{I}).$$  \hspace{1cm} (2.32)

From the upper expression we see that we can express $\Phi_0$ as an action of the ghost number one operator $D_L$ acting on the identity field. Then we immediately obtain

$$Q'(X) = Q(X) + D_L(\mathcal{I}) \ast X - (-1)^{|X|} X \ast D_L(\mathcal{I}) =$$

$$= Q(X) - \mathcal{I} \ast D_R(X) - D_L(X) \ast \mathcal{I} = Q(X) - D(X).$$  \hspace{1cm} (2.34)

In the calculation performed above it is important to stress that if two operators $Q, K$ are derivatives of the string field algebra

$$Q(A \ast B) = Q(A) \ast B + (-1)^{|A|} A \ast Q(B),$$

$$K(A \ast B) = K(A) \ast B + A \ast K(B)$$  \hspace{1cm} (2.35)

then so is its commutator. To see this, let us define $C \equiv [Q, K]$. Then we can write

$$C(A \ast B) = Q(K(A) \ast B + A \ast K(B)) - K(Q(A) \ast B + (-1)^{|A|} A \ast Q(B)) =$$

$$= QK(A) \ast B + (-1)^{|A|} K(A) \ast Q(B) + Q(A) \ast K(B) + (-1)^{|A|} A \ast QK(B) -$$

$$-KQ(A) \ast B - Q(A) \ast K(B) - (-1)^{|A|} K(A) \ast Q(B) - (-1)^{|A|} A \ast KQ(B) =$$

$$= [Q, K](A) \ast B + (-1)^{|A|} A \ast [Q, K](B) = C(A) \ast B + (-1)^{|A|} A \ast C(B).$$  \hspace{1cm} (2.36)
Now it is clear that the BRST operator $Q'$ is completely determined by commutators of $K$ with $Q$. This calculation can be easily performed using

$$\left[ \frac{1}{2\pi i} \oint dz c(z) T_m(z) , \frac{A}{2\pi i} \int_C dwk(w) \right] = \frac{A}{2\pi i} \int_C d\omega \left[ \frac{1}{2\pi i} \oint C' dz c(z) T_m(z) k(w) \right] = \frac{A}{2\pi i} \int_C d\omega \left( c(w) k(w) \right) \frac{A}{2\pi i} \left( c(\epsilon) k(\epsilon) - c(-\epsilon) k(-\epsilon) \right).$$

(2.37)

In the previous expression $C'$ is a small circle around the point $w$ on $C$. Note that the operator $K$ explicitly depends on the integration contour which in our case is semicircle of radius $\epsilon \ll 1$ in the upper half-plane $\text{Im} \ z \geq 0$. In fact, we have defined operator $K$ by action on the general string field which is functional of string coordinates $X(\sigma, \tau = 0)$ evaluated at time $\tau = 0$ so that $C$ is a semicircle of radius $\epsilon$ around the point $z = 0$. As usual we take the limit $\epsilon \to 0$ in the end of calculation. On the other hand $Q$ is conserved charge so that we can deform the integration contour on which is defined and then using OPE between $T_m$ and $k$ we can calculate the integral over $C'$.

We can also see from (2.23) that there is nontrivial OPE between $X_L, X_R$ as a consequence of the boundary conditions at Im $z = 0$ imposed on the fields $X = X_L + X_R$. This suggests that we should calculate following term as well

$$\left[ -\frac{1}{2\pi i} \oint d\zeta \zeta c(\zeta) \tilde{T}_m(\zeta) , \frac{A}{2\pi i} \int_C dwk(w) \right].$$

(2.38)

However it is easy to see that this expression is equal to zero. This follows from the fact that the OPE between $\tilde{T}(\zeta) k(w)$ contains terms such as $1/(\zeta - w)$. On the other hand, if $Q$ is given as an integral of holomorphic and anti holomorphic currents defined in the region $\text{Im} \ z \geq 0$ then the integral over $\zeta$ lies in the lower half-plane $\text{Im} \ z \leq 0$. Then it is clear that the integral over complex conjugate of $C'$ does not encircle any singularity and hence it is equal to zero. The only possible dangerous terms could arise from the real axis however we can avoid these problems by defining the integration contour $C$ in $K$ starting slightly above the real axis, say at $z = -e^{-i\epsilon} \epsilon \ll 1$ and then after working out the previous commutator we take the limit $\epsilon \to 0$. As a result, the first term in (2.33) is equal to

$$[Q, K] = \frac{A}{2\pi i} \left( c(\epsilon) k(\epsilon) - c(-\epsilon) k(-\epsilon) \right).$$

(2.39)

Next term is equal to

$$[K, [Q, K]]$$

(2.40)
Now it is easy to see that further commutators in (2.33) are equal to zero thanks to the trivial OPE between $k(z)k(w)$. Finally we obtain the resulting BRST operator

$$Q' = Q - \frac{A}{2\pi i} c(\epsilon)k(\epsilon) + \frac{A}{2\pi i} c(-\epsilon)k(-\epsilon) .$$

(2.41)

Now we come to the interpretation of this result. It is clear from (2.41) that SFT expanded around the classical solution is in some sense related to the deformed $BCFT''$ that arises from the original one by marginal deformation inserted on the boundary of the world-sheet

$$S_{boud} = -A \int dt k(t) .$$

(2.42)

The question is what means the claim “in some sense related”. First of all, the fluctuation string field still belongs to the Hilbert space of the original $BCFT$. The correlation functions which define SFT action in the CFT formalism are also evaluated in the original $BCFT$ so it is certainly false to claim that we now have SFT formulated around the new background $BCFT''$. To find relation between SFT action for fluctuation field and SFT action defined on the $BCFT''$ theory we should perform the same analysis as in very nice papers in \[82, 81, 79, 78, 80\]. Since this analysis is rather complicated it deserves separate publication and we hope to obtain some explicit results in the near future.

We must also stress that the previous example is strictly speaking valid for rolling tachyon solution since only in this case the operator $K$ is real. For that reason we will rather consider following operators $K_{2s,t}$

$$K_{2s,t} = \frac{A}{4\pi i} \int_C dz k^+(z) + \frac{A}{4\pi i} \int_C dz k^-(z) ,$$

$$k^+_t(z) = \exp \left( \frac{2}{\sqrt{\alpha'}} X^0_{L}(z) \right) , \quad k^-_t(z) = \exp \left( -\frac{2}{\sqrt{\alpha'}} X^0_{L}(z) \right) ,$$

$$k^+_s(z) = \exp \left( \frac{2i}{\sqrt{\alpha'}} X^1_{L}(z) \right) , \quad k^-_s(z) = \exp \left( -\frac{2i}{\sqrt{\alpha'}} X^1_{L}(z) \right) .$$

(2.43)

Since the operators $k^+_s, k^-_s$ have conformal weights 1 we immediately see that $K_{2s,t}$ obeys the same rules as $K$ in (2.7), so that the string field (2.7) with $K$ given above is the exact solution of SFT equation of motion.

To obtain the form of the shifted BRST operator we proceed as in the previous case. Now the first commutator is equal to

$$[Q, K_{2s,t}] = \frac{A}{4\pi i} \left( c(\epsilon)k^+_s(\epsilon) - c(-\epsilon)k^+_s(-\epsilon) \right) + \frac{A}{4\pi i} \left( c(\epsilon)k^-_s(\epsilon) - c(-\epsilon)k^-_s(-\epsilon) \right) .$$

(2.44)
The second term in (2.33) is equal to
\[
\frac{1}{2\alpha'} [K_{2s,t}, [Q, K_{2s,t}]]
\] (2.45)
which using (2.44) gives
\[
\left[K_{s,t}, c(\tau) A \frac{1}{4\pi i} k^+_{s,t}(\tau) \right]_{\tau = \epsilon, -\epsilon} = \frac{A}{4\pi i} \left[ \frac{1}{4\pi i} \oint_C dw k^-_{s,t}(w), c(\tau) k^+_{s,t}(\tau) \right]_{\tau = \epsilon, -\epsilon} = \frac{A}{8\pi i} \oint_C dw c(\tau) \left( \frac{1}{(w - \tau)^2} - \frac{1}{w - \tau} k^3_{s,t}(\tau) \right)_{\tau = \epsilon, -\epsilon}
\]

and
\[
\left[K_{s,t}, A \frac{1}{4\pi i} c(\tau) k^-_{s,t}(\tau) \right]_{\tau = -\epsilon, \epsilon} = \left[ \frac{A}{4\pi i} \int dw k^+_{s,t}(w), \frac{A}{4\pi i} c(\tau) k^-_{s,t}(\tau) \right]_{\tau = -\epsilon, \epsilon} = \frac{A}{8\pi i} \int dw c(\tau) \left( \frac{1}{(w - \tau)^2} + \frac{1}{w - \tau} k^3_{s,t}(\tau) \right)_{\tau = -\epsilon, \epsilon} = \frac{A}{8\pi i} c(\tau) k^3_{s,t}(\tau)_{\tau = -\epsilon, \epsilon}
\] (2.46)

In the previous calculations we used
\[
k^+_{s,t}(z) k^-_{s,t}(w) = \frac{1}{(z - w)^2} + \frac{1}{z - w} k^3_{s,t},
\]
\[
k^3_{s,t}(z) = \frac{2}{\sqrt{\alpha'}} \partial_z X^0_L(z), k^3_{s,t}(z) = \frac{2i}{\sqrt{\alpha'}} \partial_z X^1_L(z).
\] (2.48)

From (2.46), (2.47) we immediately see that the second term in (2.33) is equal to zero and consequently (2.44) is the only nonzero term in (2.33). If we choose
\[
A = -2\pi \lambda
\] (2.49)
we obtain the shifted BRST operator
\[
Q'_{s,t} = Q + i\lambda c(\epsilon) k^+_{s,t}(\epsilon) - \frac{\lambda}{2} c(-\epsilon) k^+_{s,t}(-\epsilon) + i\lambda c(-\epsilon) k^-_{s,t}(-\epsilon) \Rightarrow \]
\[
Q'_{s} = Q + i\lambda c(\epsilon) \cosh \left( \frac{2}{\sqrt{\alpha'}} X^0(\epsilon) \right) - i\lambda c(-\epsilon) \cosh \left( \frac{2}{\sqrt{\alpha'}} X^0(-\epsilon) \right),
\]
\[
Q'_{t} = Q + i\lambda c(\epsilon) \cos \left( \frac{2}{\sqrt{\alpha'}} X^1(\epsilon) \right) - i\lambda c(-\epsilon) \cos \left( \frac{2}{\sqrt{\alpha'}} X^1(-\epsilon) \right),
\] (2.50)
where $Q_t$ corresponds to the $k_t^{+,-}$ in (2.43). With the analogy with the previous calculation we would like to interpret this deformation of the BRST operator as the result of the marginal interaction inserted on the boundary of the world-sheet

$$S_{\text{bound}} = \lambda \int d\tau \cosh \left( \frac{2}{\sqrt{\alpha'}} X^0(\tau) \right) .$$

(2.51)

This $BCFT$ was extensively studied in [1, 2, 3, 4, 5] and corresponds to the rolling tachyon solution. Once again we must stress that this claim should be proved in the same way as in [82, 81, 79, 78, 80]. At this place we have only shifted BRST operator in the Hilbert space of the original $BCFT$ and with the same correlations functions.

In the same way we can expect that $Q'_s$ corresponds to the deformed $BCFT''$ with the boundary interaction term

$$S_{\text{bound}} = \lambda \int d\tau \cos \left( \frac{2}{\sqrt{\alpha'}} X^1(\tau) \right) .$$

(2.52)

This $BCFT$ was extensively studied in [51, 52, 53, 54, 55].

It is remarkable that (2.5) can be expressed in the pure gauge form which could lead to the suspicion that this solution does not give any new physical information. We would like to argue that this is not the case. As is well known the string field theory action (2.3) is invariant under the small gauge transformations

$$\delta \Phi = Q \Lambda - \Lambda \star \Phi + \Phi \star \Lambda ,$$

(2.53)

where $\Lambda$ is ghost number zero string field. On the other hand this action is not generally invariant under the large gauge transformations

$$\Phi' = e^{-\Lambda} \star Q(e^\Lambda) + e^{-\Lambda} \star \Phi \star e^\Lambda .$$

(2.54)

As is well known there is a sharp distinction between the small gauge transformations and the large ones, for very nice discussion, see [83]. As was argued there, the small gauge transformation describes redundancy in our description of the theory. On the other hand, large gauge transformations are true symmetries that relate different solutions in given gauge theory which in our case is the open bosonic string field theory.

3. Spectrum of the fluctuations around the new background

In the previous section we have found an exact solution of the string field theory and after expansion above this solution an action for fluctuation modes that has the same form as the original one however with new, shifted BRST operator $Q'$. The form of this operator suggests that the string field theory formulated around this classical solution should be equivalent to the string field theory formulated around the new
that arises from the original BCFT by exactly marginal deformation. In this section we will try to give arguments that could support this claim. More precisely, we will try to determine spectrum of fluctuation modes around the classical solution and these modes should correspond to the states in the deformed BCFT. 

We start with the well known fact that fluctuation modes around the classical solution obey linearised SFT equation of motion

\[ Q'\Psi = Q(\Psi) - D(\Psi) = 0 \]  

To proceed observe that the new BRST operator \( Q' \) can be written as

\[ Q'(X) = e^K(Q(e^{-K}(X)) \]  

We see that it is natural to propose following form of the fluctuation field

\[ \Psi = e^K(\Phi) \]  

since then we get

\[ Q'(\Psi) = 0 = e^KQ(e^{-K}e^K(\Phi)) = 0 \Rightarrow Q(\Phi) = 0 \]  

We will see that the consistency of the definition of the action of \( K \) on the state \( \Psi \) and the requirement of the well defined vertex operator corresponding to the string field \( (3.3) \) will lead to the constraints on the allowed spectrum of the fluctuation field.

In this section we restrict ourselves to operator \( K_{2s} \equiv K \). The extension to time dependent solution would be straightforward, following [4]. So that let us consider following operator

\[ K = \frac{1}{2\pi i} \int_C dz \left( : \exp \left( i \frac{2}{\sqrt{\alpha'}} X_1^1(z) \right) : + : \exp \left( -i \frac{2}{\sqrt{\alpha'}} X_1^1(z) \right) : \right) \equiv K_+ + K_- \]  

Let us consider the ground state vertex operator of original BCFT \( 6 \)

\[ \Phi(0, 0) = c(0) : e^{ik_\mu X^\mu(0,0)} := c(0) : e^{2ik_\mu X^\mu_L(0)} : , \]

\[ Q(\Phi(0, 0)) = 0 \Rightarrow -k_0^2 + k_1^2 = \frac{1}{\alpha'} , \]

where we consider the dependence on \( k_1 \) only. In the following we will use these formulas

\[ : \exp \left( i \frac{2}{\sqrt{\alpha'}} X_1^1(z) \right) : e^{2ik_\mu X^\mu_L(0)} := e^{2ik(0)/\alpha} \]  

\[ : \exp \left( -i \frac{2}{\sqrt{\alpha'}} X_1^1(z) \right) : e^{2ik_\mu X^\mu_L(0)} := e^{-2ik(0)/\alpha} \]

\[ : \exp \left( i \frac{2}{\sqrt{\alpha'}} X_1^1(z) \right) : e^{2ik_\mu X^\mu_L(0)} := e^{2ik(0)/\alpha} \]  

\[ : \exp \left( -i \frac{2}{\sqrt{\alpha'}} X_1^1(z) \right) : e^{2ik_\mu X^\mu_L(0)} := e^{-2ik(0)/\alpha} \]  

\[ \]  

\[ \]
We use these OPE in the calculation of the commutators $[K, \Phi(0, 0)]$ that arise from the formula
\[
\Psi = e^K(\Phi) = e^K(\Phi(0, 0)|0) = e^K\Phi(0, 0)e^{-K}e^K|0 = e^K\Phi(0, 0)e^{-K}|0 .
\] (3.8)

Now we have
\[
\Psi(0, 0) = e^K\Phi(0, 0)e^{-K} = \Phi(0, 0) + \sum_{N=1}^{\infty} \frac{1}{N!} [K, \ldots, [K, \Phi(0, 0)]] .
\] (3.9)

The commutator $[K, \Phi]$ is given by contour integral around $z = 0$. To perform this calculation we use OPE (3.7). However in order these integrals to be well defined we see that the power of $z$ should be integer. This requirement together with (3.6) gives quantisation of the momentum $k_1 \equiv k$
\[
k = \frac{n}{2\sqrt{\alpha'}}, n = \pm 2, \pm 3, \pm 4 \ldots .
\] (3.10)

We can interpret this result as a selection rule for momentum states. In fact the same situation arises in the study of the BCFT with the boundary integration given (2.52).

Let us start with the first allowed state with $n = 2$
\[
\Phi_2(0, 0) = e^{\frac{\sqrt{\alpha'}}{\alpha'}X_L(0)} .
\] (3.11)

For this state we get
\[
[K, \Phi_2(0, 0)] = [K, , \Phi_2(0, 0)] = \frac{1}{2\pi i} \oint dw \frac{1}{w^2} : e^{\frac{\sqrt{\alpha'}}{\alpha'}X_L(w)} e^{\frac{\sqrt{\alpha'}}{\alpha'}X_L(0)} :=
\]
\[
= \frac{\partial}{\partial z} \left[ \frac{1}{2\pi i} \oint dw \frac{1}{w - z} : e^{\frac{\sqrt{\alpha'}}{\alpha'}X_L(w)} \right]_{z=0} e^{\frac{\sqrt{\alpha'}}{\alpha'}X_L(0)} :=
\]
\[
= -\frac{2i}{\sqrt{\alpha'}} \partial_w X_L(0) : e^{\frac{\sqrt{\alpha'}}{\alpha'}X_L(0)} e^{\frac{\sqrt{\alpha'}}{\alpha'}X_L(0)} := -\frac{2i}{\sqrt{\alpha'}} \partial_w X_L(0)
\] (3.12)

and also
\[
[K, [K, \Phi_2]] = [K, -\frac{2i}{\sqrt{\alpha'}} \partial_w X_L(0)] =
\]
\[
= -\frac{2i}{\sqrt{\alpha'}} \frac{1}{2\pi i} \oint dw \left( \frac{i\sqrt{\alpha'}}{w} e^{\frac{\sqrt{\alpha'}}{\alpha'}X_L(w)} - \frac{i\sqrt{\alpha'}}{w} e^{-\frac{\sqrt{\alpha'}}{\alpha'}X_L(w)} \right) =
\]
\[
= -\frac{2i}{\sqrt{\alpha'}} \frac{1}{2\pi i} \oint dw \left( e^{\frac{\sqrt{\alpha'}}{\alpha'}X_L(0)} - e^{-\frac{\sqrt{\alpha'}}{\alpha'}X_L(0)} \right),
\]
\[
[K, [K, [K, \Phi_2(0, 0)]]] = -\frac{8i}{\sqrt{\alpha'}} \partial_w X_L(w) .
\] (3.13)
We see that further commutators will give the same components as the three ones listen above. In other words the fluctuation field is given as

\[ \Psi = \Phi_2 + [K, \Phi_2] + \frac{1}{2} [K, [K, \Phi_2]] + \frac{1}{6} [K, [K, [K, \Phi_2]]] + \ldots = \]

\[ = e^{\sqrt{\alpha}} X_L(0) - \frac{i}{\sqrt{\alpha}} \partial_w X_L(0) + e^{\sqrt{\alpha}} X_L(0) - \frac{8i}{6 \sqrt{\alpha}} \partial_w X_L(0) + \ldots . \]

(3.14)

We observe that the fluctuation field \( \Psi \) is given as linear combination of the modes that belong to the representation of the \( SU(2) \) group with the spin \( j = 1 \). Exactly the same spectrum of the states arise in \( BCFT'' \) with the marginal interaction (2.52).

Generally we can argue as follows. Let us consider operator \( \Phi_n(z) = e^{\sqrt{\alpha}} X_L(z) \) with the spin equal to \( n/2 \) as can be seen from

\[ [J_3, \Phi_n(0)] = \frac{i}{\sqrt{\alpha}} \int \frac{dw}{2\pi i} \partial_w X_L(w) e^{\sqrt{\alpha}} \]

\[ = \frac{n}{2} \frac{i}{2\pi i} \int \frac{dw}{w - z} e^{\sqrt{\alpha}} X_L(z) = \frac{n}{2} \Phi_n(0,0) , \]

(3.15)

where

\[ J_3 = \frac{i}{\sqrt{\alpha}} \int dw \partial X_L^1(w) . \]

(3.16)

From (3.7) we get that \( \Phi_n \) is annihilated by \( K_+ \) for \( n \geq 0 \) and also following commutation relations

\[ [K_+, K_-] = 2J_3 , [J_3, K_+] = K_+ , [J_3, K_-] = -K_- . \]

(3.17)

Then \( [K, \Phi_n] = [K_-, \Phi_n] \) and the resulting operator will have the spin equal to \( n/2 - 1 \) as follows from

\[ [J_3, [K_-, \Phi_n]] + [K_-, [\Phi_n, J_3]] + [\Phi_n, [J_3, K_-]] = 0 \Rightarrow \]

\[ [J_3, [K_-, \Phi_n]] = \left( \frac{n}{2} - 1 \right) [K_-, \Phi_n] . \]

(3.18)

Then the successive action of \( K \) on \( \Phi_n \) will give states with lower spin and hence the fluctuation field is linear combination of states that form an irreducible representation of \( SU(2) \) gauge group with the spin \( n/2 \). The same spectrum of states arises in the \( BCFT'' \) with the marginal interaction on the boundary (2.52). We think that this coincidence supports our conjecture based on the form of the shifted BRST operator that the string field theory for fluctuation modes around the classical solution is related to he string field theory defined on the background \( BCFT'' \) that arises from the original one by marginal interaction inserted on the real line. Note that we do not claim that the string field action for fluctuation modes around the classical solution is exactly the same as the string field theory action formulated around the new background BCFT. To find relation between these two actions we should perform the same profound analysis as in [82, 81, 79, 78, 80]. We return to this question in the future.
4. Conclusion

In this short note we have proposed an exact solution of the open bosonic string field theory. The basic idea of our approach is to find such an operator of conformal weight one that has nontrivial commutation relation with the BRST operator. The condition that this operator has conformal dimension equal to one is important since only in this case this operator obeys \((2.13), (2.19)\) and then we can easily find the form of the shifted BRST operator. The form of shifted BRST operator suggests that SFT for fluctuation field above the classical solution is related to the SFT defined on \(BCFT''\) that arises from the original \(BCFT\) by insertion of the marginal deformation on the boundary of the world-sheet. In order to support this claim we have performed an analysis of the linearised equation of motion of the SFT around the classical solution and we have argued that the spectrum of the fluctuation modes agree with the spectrum of states in \(BCFT''\).

Of course, as we have also stressed in many places in this paper, we cannot claim that we have really prove that the SFT action for fluctuation modes is exactly the same as the SFT action formulated around background \(BCFT''\). This can be easily seen from the fact that the fluctuation states and correlation functions in the shifted action are still formulated using the original \(BCFT\). In order to find exact relation between action for fluctuation modes and SFT action defined on \(BCFT''\) we should perform the same careful analysis as in \([22, 81]\) \([73, 78, 80]\). This problem is currently under investigation and we hope to report our result in near future.

We must also stress that there is important limitation of our solution when we consider it in the context rolling tachyon \([1, 2, 3, 4, 5]\). Since this is solution of the SFT in the limit \(g_s \to 0\) we do not take into account the coupling between closed a open strings in the process of the rolling tachyon. When this coupling is considered the situation is much more involved and interesting, as was shown in recent papers \([12, 60, 84]\). It would be very nice to extend our analysis to the case of open-closed string field theory \([55]\). It would be also interesting to search for an exact solution in the string field theory that would lead to the more general boundary perturbation as was studied in \([4]\).

To conclude, we hope that our modest contribution presented in this paper could be helpful in solving these intriguing problems.

Acknowledgment I would like to thank Ulf Danielsson and Ulf Lindström for their support in my research. This work is partly supported by EU contract HPRN-CT-2000-00122.

References

[1] A. Sen, “Rolling tachyon,” JHEP 0204, 048 (2002) [arXiv:hep-th/0203211].

[2] A. Sen, “Tachyon matter,” arXiv:hep-th/0203265.
[3] A. Sen, “Field theory of tachyon matter,” arXiv:hep-th/0204143.

[4] A. Sen, “Time evolution in open string theory,” arXiv:hep-th/0207105.

[5] P. Mukhopadhyay and A. Sen, “Decay of unstable D-branes with electric field,” arXiv:hep-th/0208142.

[6] N. Moeller and B. Zwiebach, “Dynamics with infinitely many time derivatives and rolling tachyons,” arXiv:hep-th/0207107.

[7] K. Hashimoto, “Dynamical decay of brane-antibrane and dielectric brane,” JHEP 0207, 035 (2002) [arXiv:hep-th/0204203].

[8] S. Sugimoto and S. Terashima, “Tachyon matter in boundary string field theory,” JHEP 0207, 025 (2002) [arXiv:hep-th/0205085].

[9] J. A. Minahan, “Rolling the tachyon in super BSFT,” JHEP 0207, 030 (2002) [arXiv:hep-th/0205098].

[10] K. Ohta and T. Yokono, “Gravitational approach to tachyon matter,” arXiv:hep-th/0207004.

[11] T. Mehen and B. Wecht, “Gauge fields and scalars in rolling tachyon backgrounds,” arXiv:hep-th/0206212.

[12] T. Okuda and S. Sugimoto, “Coupling of rolling tachyon to closed strings,” arXiv:hep-th/0208196.

[13] G. W. Gibbons, “Cosmological evolution of the rolling tachyon,” Phys. Lett. B 537, 1 (2002) [arXiv:hep-th/0204008].

[14] M. Fairbairn and M. H. Tytgat, “Inflation from a tachyon fluid?,” arXiv:hep-th/0204070.

[15] S. Mukohyama, “Brane cosmology driven by the rolling tachyon,” Phys. Rev. D 66, 024009 (2002) [arXiv:hep-th/0204084].

[16] S. Mukohyama, “Inhomogeneous tachyon decay, light-cone structure and D-brane network problem in tachyon cosmology,” arXiv:hep-th/0208094.

[17] A. Feinstein, “Power-law inflation from the rolling tachyon,” arXiv:hep-th/0204140.

[18] T. Padmanabhan, “Accelerated expansion of the universe driven by tachyonic matter,” Phys. Rev. D 66, 021301 (2002) [arXiv:hep-th/0204150].

[19] A. Frolov, L. Kofman and A. A. Starobinsky, “Prospects and problems of tachyon matter cosmology,” arXiv:hep-th/0204187.

[20] D. Choudhury, D. Ghoshal, D. P. Jatkar and S. Panda, “On the cosmological relevance of the tachyon,” arXiv:hep-th/0204204.
[21] X. z. Li, J. g. Hao and D. j. Liu, “Can quintessence be the rolling tachyon?,” arXiv:hep-th/0204252.

[22] G. Shiu and I. Wasserman, “Cosmological constraints on tachyon matter,” Phys. Lett. B 541, 6 (2002) [arXiv:hep-th/0205003].

[23] T. Padmanabhan and T. R. Choudhury, “Can the clustered dark matter and the smooth dark energy arise from the same scalar field?,” arXiv:hep-th/0205055.

[24] L. Kofman and A. Linde, “Problems with tachyon inflation,” JHEP 0207, 004 (2002) [arXiv:hep-th/0205121].

[25] H. B. Benaoum, “Accelerated universe from modified Chaplygin gas and tachyonic fluid,” arXiv:hep-th/0205140.

[26] M. Sami, “Implementing power law inflation with rolling tachyon on the brane,” arXiv:hep-th/0205146.

[27] M. Sami, P. Chingangbam and T. Qureshi, “Aspects of tachyonic inflation with exponential potential,” Phys. Rev. D 66, 043530 (2002) [arXiv:hep-th/0205179].

[28] G. Shiu, S. H. Tye and I. Wasserman, “Rolling tachyon in brane world cosmology from superstring field theory,” arXiv:hep-th/0207119.

[29] Y. S. Piao, R. G. Cai, X. m. Zhang and Y. Z. Zhang, “Assisted tachyonic inflation,” arXiv:hep-ph/0207143.

[30] X. z. Li, D. j. Liu and J. g. Hao, “On the tachyon inflation,” arXiv:hep-th/0207146.

[31] J. M. Cline, H. Firouzjahi and P. Martineau, “Reheating from tachyon condensation,” arXiv:hep-th/0207156.

[32] G. N. Felder, L. Kofman and A. Starobinsky, “Caustics in tachyon matter and other Born-Infeld scalars,” arXiv:hep-th/0208019.

[33] B. Wang, E. Abdalla and R. K. Su, “Dynamics and holographic discreteness of tachyonic inflation,” arXiv:hep-th/0208023.

[34] M. C. Bento, O. Bertolami and A. A. Sen, Tachyonic inflation in the braneworld scenario,” arXiv:hep-th/0208124.

[35] A. Buchel, P. Langfelder and J. Walcher, “Does the tachyon matter?,” arXiv:hep-th/0207235.

[36] G. R. Dvali and S. H. Tye, “Brane inflation,” Phys. Lett. B 450, 72 (1999) [arXiv:hep-ph/9812483].

[37] A. Sen, “Stable non-BPS bound states of BPS D-branes,” JHEP 9808, 010 (1998) [arXiv:hep-th/9805019].
[38] A. Sen, “SO(32) spinors of type I and other solitons on brane-antibrane pair,” JHEP 9809, 023 (1998) [arXiv:hep-th/9808141].

[39] A. Sen, “BPS D-branes on non-supersymmetric cycles,” JHEP 9812, 021 (1998) [arXiv:hep-th/9812031].

[40] A. Sen, “Descent relations among bosonic D-branes,” Int. J. Mod. Phys. A 14, 4061 (1999) [arXiv:hep-th/9902105].

[41] A. Sen, “Universality of the tachyon potential,” JHEP 9912, 027 (1999) [arXiv:hep-th/9911116].

[42] A. Sen and B. Zwiebach, “Tachyon condensation in string field theory,” JHEP 0003, 002 (2000) [arXiv:hep-th/9912249].

[43] K. Ohmori, “A review on tachyon condensation in open string field theories,” arXiv:hep-th/0102085.

[44] I. Y. Arefeva, D. M. Belov, A. A. Giryavets, A. S. Koshelev and P. B. Medvedev, “Non-commutative field theories and (super)string field theories,” arXiv:hep-th/0111208.

[45] P. J. De Smet, “Tachyon condensation: Calculations in string field theory,” arXiv:hep-th/0109182.

[46] P. Di Vecchia and A. Liccardo, “D branes in string theory. I,” arXiv:hep-th/9912161.

[47] P. Di Vecchia and A. Liccardo, “D-branes in string theory. II,” arXiv:hep-th/9912275.

[48] M. R. Gaberdiel, “Lectures on non-BPS Dirichlet branes,” Class. Quant. Grav. 17, 3483 (2000) [arXiv:hep-th/0005029].

[49] A. Lerda and R. Russo, “Stable non-BPS states in string theory: A pedagogical review,” Int. J. Mod. Phys. A 15, 771 (2000) [arXiv:hep-th/9905006].

[50] A. Sen, ‘Non-BPS states and branes in string theory,” arXiv:hep-th/9904207.

[51] C. G. Callan and I. R. Klebanov, “Exact C = 1 boundary conformal field theories,” Phys. Rev. Lett. 72, 1968 (1994) [arXiv:hep-th/9311092].

[52] C. G. Callan, I. R. Klebanov, A. W. Ludwig and J. M. Maldacena, “Exact solution of a boundary conformal field theory,” Nucl. Phys. B 422, 417 (1994) [arXiv:hep-th/9402113].

[53] J. Polchinski and L. Thorlacius, “Free Fermion Representation Of A Boundary Conformal Field Theory,” Phys. Rev. D 50, 622 (1994) [arXiv:hep-th/9404008].

[54] A. Recknagel and V. Schomerus, “Boundary deformation theory and moduli spaces of D-branes,” Nucl. Phys. B 545, 233 (1999) [arXiv:hep-th/9811237].

[55] M. R. Gaberdiel and A. Recknagel, “Conformal boundary states for free bosons and fermions,” JHEP 0111, 016 (2001) [arXiv:hep-th/0108238].
[56] M. Gutperle and A. Strominger, “Spacelike branes,” JHEP 0204, 018 (2002) [arXiv:hep-th/0202210].

[57] S. Roy, “On supergravity solutions of space-like Dp-branes,” JHEP 0208, 025 (2002) [arXiv:hep-th/0205198].

[58] M. Kruczenski, R. C. Myers and A. W. Peet, “Supergravity S-branes,” JHEP 0205, 039 (2002) [arXiv:hep-th/0204144].

[59] C. M. Chen, D. V. Gal’tsov and M. Gutperle, “S-brane solutions in supergravity theories,” Phys. Rev. D 66, 024043 (2002) [arXiv:hep-th/0204071].

[60] A. Strominger, “Open String Creation by S-Branes,” arXiv:hep-th/0209090.

[61] J. Kluson, “Time dependent solution in open bosonic string field theory,” arXiv:hep-th/0208028.

[62] G. T. Horowitz, J. Lykken, R. Rohm and A. Strominger, “A Purely Cubic Action For String Field Theory,” Phys. Rev. Lett. 57, 283 (1986).

[63] T. Takahashi and S. Tanimoto, “Wilson lines and classical solutions in cubic open string field theory,” Prog. Theor. Phys. 106, 863 (2001) [arXiv:hep-th/0107046].

[64] T. Takahashi and S. Tanimoto, “Marginal and scalar solutions in cubic open string field theory,” JHEP 0203, 033 (2002) [arXiv:hep-th/0202133].

[65] J. Kluson, “Some solutions of Berkovits’ superstring field theory,” arXiv:hep-th/0201054.

[66] J. Kluson, “Exact solutions of open bosonic string field theory,” JHEP 0204, 043 (2002) [arXiv:hep-th/0202045].

[67] J. Kluson, “Marginal deformations in the open bosonic string field theory for N D0-branes,” arXiv:hep-th/0203089.

[68] J. Kluson, “New solution of the open bosonic string field theory,” arXiv:hep-th/0205294.

[69] E. Witten, “Noncommutative Geometry And String Field Theory,” Nucl. Phys. B 268, 253 (1986).

[70] I. Ellwood, B. Feng, Y. H. He and N. Moeller, “The identity string field and the tachyon vacuum,” JHEP 0107, 016 (2001) [arXiv:hep-th/0105024].

[71] Y. Matsuo, “Identity projector and D-brane in string field theory,” Phys. Lett. B 514, 407 (2001) [arXiv:hep-th/0106027].

[72] I. Kishimoto, “Some properties of string field algebra,” JHEP 0112, 007 (2001) [arXiv:hep-th/0110124].
[73] M. Schnabl, “Wedge states in string field theory,” arXiv:hep-th/0201095.

[74] I. Kishimoto and K. Ohmori, “CFT description of identity string field: Toward derivation of the VSFT action,” JHEP 0205, 036 (2002) [arXiv:hep-th/0112169].

[75] A. LeClair, M. E. Peskin and C. R. Preitschopf, “String Field Theory On The Conformal Plane. 1. Kinematical Principles,” Nucl. Phys. B 317, 411 (1989).

[76] A. LeClair, M. E. Peskin and C. R. Preitschopf, “String Field Theory On The Conformal Plane. 2. Generalized Gluing,” Nucl. Phys. B 317, 464 (1989).

[77] K. Bardakci and A. Konechny, Nucl. Phys. B 598 (2001) 427 [arXiv:hep-th/0009214].

[78] A. Sen, “On The Background Independence Of String Field Theory,” Nucl. Phys. B 345, 551 (1990).

[79] A. Sen, “On The Background Independence Of String Field Theory. 2. Analysis Of On-Shell S Matrix Elements,” Nucl. Phys. B 347, 270 (1990).

[80] A. Sen, “On the background independence of string field theory. 3. Explicit Field redefinitions,” Nucl. Phys. B 391, 550 (1993) [arXiv:hep-th/9201041].

[81] A. Sen and B. Zwiebach, “A Proof of local background independence of classical closed string field theory,” Nucl. Phys. B 414, 649 (1994) [arXiv:hep-th/9307088].

[82] A. Sen and B. Zwiebach, “Background independent algebraic structures in closed string field theory,” Commun. Math. Phys. 177, 305 (1996) [arXiv:hep-th/9408053].

[83] J. A. Harvey, “Magnetic monopoles, duality, and supersymmetry,” arXiv:hep-th/9603086.

[84] A. Sen, “Time and Tachyon,” arXiv:hep-th/0209122.

[85] B. Zwiebach, “Oriented open-closed string theory revisited,” Annals Phys. 267, 193 (1998) [arXiv:hep-th/9705241].

[86] A. Strominger, “Lectures On Closed String Field Theory,” IASSNS-HEP-87/28 [SPIRES entry] Delivered at ICTP School on Superstrings, Trieste, Italy, Apr 1-15, 1987.