New Solution of Open Bosonic String Field Theory

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ABSTRACT: We present example of exact solution to Witten’s open bosonic string field theory. We will analyse the new BRST operator and we will argue that the new solution describes the flow from zero slope limit ($\alpha' \to 0$) to the tensionless limit ($\alpha' \to \infty$) in the string world-sheet action.

KEYWORDS: String field theory

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1. Introduction

Recently there has been great interest in open bosonic string field theory [5], especially in the context of the tachyon condensation (For a review of string field theory, see [1, 2, 3, 4].) The remarkable success of string field theory in addressing this question leads us to believe that string field theory has a profound role to play in the formulation of string theory.

In some previous papers we have proposed a method for finding exact solutions to the open bosonic string field theory [6, 7, 8]. This method relies on the existence of a ghost number zero operator that does not commute with the BRST operator so that the string field theory formulated around the new background corresponds to a changed BRST operator.

In this short note we continue this investigation and find new exact solutions to the open bosonic string field theory action and the corresponding modified BRST operator which have very unusual properties. We show that the form of the new BRST operator depends on a single parameter $t$ that can vary from $0$ to $\infty$. For $t = 0$ the BRST operator is the same as the BRST operator for a massless particle. For $t = 1$ we have the original D25-brane, and finally, for $t = \infty$ the BRST operator corresponds to the BRST operator for a tensionless string.

A crucial point is that the effective tension that we introduce during the calculation appears to effectively depend on the form of the string field theory solution depends on the direction where we allow our operator $K$ applies.

This has the remarkable consequence in the emergence of a string with anisotropic string tension. A similar situation has previously been studied in the context of the noncommutative open string (NCOS) [15, 19, 20, 21, 22, 23, 24, 25, 26, 27, 28, 29].
Our result deserves further study in that context. We believe our results merit further study. Especially it would be nice to perform an explicit CFT analysis of a world-sheet theory with an-isotropic tension.

2. Review of the general formalism

In this section we briefly review the general formalism introduced in [6] and then applied in [7, 8].

The action of the open bosonic string field theory [5] has the form

\[ S = - \left( \frac{1}{2} \int \Phi \star Q \Phi + \frac{1}{3} \int \Phi \star \Phi \star \Phi \right), \] (2.1)

Varying the action leads to the equation of motion

\[ Q \Phi + \Phi \star \Phi = 0. \] (2.2)

As demonstrated in [6], it is easy to see that the field

\[ \Phi_0 = e^{-K_L(\mathcal{I})} \star Q(e^{K_L(\mathcal{I})}) \] (2.3)

is a solution of the equation of motion (2.2) since we have

\[ Q(\Phi_0) = Q(e^{-K_L(\mathcal{I})} \star Q(e^{K_L(\mathcal{I})})) = -e^{-K_L(\mathcal{I})} \star Q(e^{K_L(\mathcal{I})}) \star e^{-K_L(\mathcal{I})} \star Q(e^{K_L(\mathcal{I})}) \] (2.4)

and

\[ \Phi_0 \star \Phi_0 = e^{-K_L(\mathcal{I})} \star Q(e^{K_L(\mathcal{I})}) \star e^{-K_L(\mathcal{I})} \star Q(e^{K_L(\mathcal{I})}). \] (2.5)

where \( K \) is a ghost number zero operator that obeys

\[ K(X \star Y) = K(X) \star Y + X \star K(Y), \]
\[ \int K(X) \star Y = -\int X \star K(Y), \] (2.6)

and where \( \mathcal{I} \) is the identity ghost number zero field [9] that obeys

\[ \mathcal{I} \star \Psi = \Psi \star \mathcal{I} = \Psi. \] (2.7)

for any string field \( \Psi \). \(^2\) In (2.3) the index \( L \) means the integration of the local density over the left side of the string that is labeled by \( \sigma \in (0, \pi/2) \). Similarly, the notation \( K_R \) corresponds to the integration of the local density over right side of the string with \( \sigma \in (\pi/2, \pi) \). \(^3\)

\(^2\)A recent discussion of this special string field is given in [10, 11].
As was shown in [6], the solution (2.3) leads to a shifted form of the BRST operator

\[ \tilde{Q}(X) = Q(X) + \Phi_0 \star X - (-1)^{|X|} X \star \Phi_0 = e^{-K} \left( Q(e^K(X)) \right) \]  

(2.8)
in the action for the fluctuation field when we perform an expansion of string field $\Phi$ around the classical solution $\Phi_0$ in (2.1). In (2.8) the symbol $|X|$ denotes the ghost number of the string field $\Phi$.

The previous results were used in [7, 8]. In the next section we will consider another form of the operator $K$, one that resembles the form of the generator of dilation in the conformal field theory.

3. Scaling transformation

Let us start with the world-sheet action for the bosonic string

\[ S = -\frac{1}{4\pi \alpha'} \int d^2 \sigma \eta^{\alpha\beta} \partial_\alpha X^K \partial_\beta X^L \eta_{LK} , \]  

(3.1)

where $\eta^{\alpha\beta}$ is two-dimensional Minkowski metric with signature $\eta_{\alpha\beta} = \text{diag}(-1,1)$ and $\eta_{KL}, K, L = 0, \ldots, 25$ is a 26-dimensional target space Minkowski metric with signature $(-, +, \ldots, +)$.

From the action (3.1) we obtain the momentum conjugate to $X^K(\sigma)$

\[ P_K(\sigma) = \frac{\delta L}{\delta X^K(\sigma)} = \frac{1}{2\pi \alpha'} \dot{X}^L(\sigma) \eta_{LK} \]  

(3.2)

with the following commutation relation

\[ [P_K(\sigma), X^L(\sigma')] = -i \delta(\sigma, \sigma') \delta_K^L , \]  

(3.3)

where $\delta(\sigma, \sigma')$ is a delta function that obeys Neumann boundary conditions so that can be written as

\[ \delta(\sigma, \sigma') = \frac{1}{\pi} \sum_{n=-\infty}^{\infty} \cos n\sigma \cos n\sigma' . \]  

(3.4)

We take the ghost number zero operator $K$ to be of the form \footnote{In our convention: $K, L = 0, \ldots, 25$, $i, j = 0, \ldots, P$, $\mu, \nu = P + 1, \ldots, 25$.}

\[ K = i \int_0^\pi d\sigma X^i(\sigma) P_i(\sigma) , \]  

(3.5)

where we do not care about a possible constant factor arising from the fact that ordering of the operators is not unique. E.g. starting from a symmetric ordering.
description and using (3.3) would give the above $K$ multiplied by a constant which we discard. For later use we calculate the following commutators

\[
[P_K P_L(\sigma) \eta^{KL}, \epsilon K] = 2 \epsilon P_i P_j(\sigma) \eta^{ij}, \\
[X'^K X'^L(\sigma) \eta_{KL}, \epsilon K] = -2 \epsilon \eta_{ij} X'^i X'^j(\sigma), \\
[X'^K P_K(\sigma), \epsilon K] = 0.
\]

(3.6)

Start with $\epsilon \ll 1$ we have

\[
\Phi_0 = e^{-K_L(\mathcal{I})} \ast Q(e^{K_L(\mathcal{I})}) \sim (\mathcal{I} - K_L(\mathcal{I})) \ast Q(\mathcal{I} + K_L(\mathcal{I})) = Q(K_L(\mathcal{I}))
\]

(3.7)

which can be written as

\[
\Phi_0 = Q(K_L(\mathcal{I})) = Q_L(K_L(\mathcal{I})) + Q_R(K_L(\mathcal{I})) = \\
= Q_L(K_L(\mathcal{I})) + K_L(Q_R(\mathcal{I})) = Q_L(K_L(\mathcal{I})) - K_L(Q_L(\mathcal{I})) = [Q_L, K_L](\mathcal{I})
\]

(3.8)

using (3).

\[
Q_L(\mathcal{I}) = -Q_R(\mathcal{I}).
\]

(3.9)

Furthermore

\[
Q_R(K_L(\mathcal{I})) = \mathcal{I} \ast Q_R(K_L(\mathcal{I})) = -Q_L(\mathcal{I}) \ast K_L(\mathcal{I}) = \\
= Q_L(\mathcal{I}) \ast K_R(\mathcal{I}) = -K_L(Q_L(\mathcal{I})) \ast \mathcal{I} = -K_L(Q_L(\mathcal{I})).
\]

(3.10)

Let us finally define $D_L(\mathcal{I})$ by

\[
\Phi_0 = D_L(\mathcal{I}) = [Q_L, K_L](\mathcal{I}).
\]

(3.11)

We then obtain the quadratic term for the fluctuation fields $\Psi$ in the form

\[
\int \Psi \ast Q' \Psi = \int \Psi \ast Q(\Psi) + \int \Psi \ast D_L(\mathcal{I}) \ast \Psi + \int \Psi \ast \Psi \ast D_L(\mathcal{I}) = \\
= \int \Psi \ast Q(\Psi) - \int \Psi \ast D(\Psi)
\]

(3.12)

where we used

\[
\int \Psi \ast D_L(\mathcal{I}) \ast \Psi = -\int \Psi \ast \mathcal{I} \ast D_R(\Psi) = -\int \Psi \ast D_R(\Psi),
\]

\[
\int \Psi \ast \Psi \ast D_L(\mathcal{I}) = -\int \Psi \ast \Psi \ast D_R(\mathcal{I}) = -\int \Psi \ast D_L(\Psi),
\]

\[
D_L(X) \ast Y = (-1)^X X \ast D_R(Y).
\]

(3.13)
The last formula may be proven as in \cite{9} using the fact that the commutator \([Q, K]\)
has odd grading. Let us calculate this commutator. Since we have
\[
Q = \frac{1}{\pi} \int_0^\pi d\sigma J_0(\sigma) = \frac{1}{\pi} \int_0^\pi d\sigma c^0(\sigma) T_{a0}(\sigma) + Q_{\text{ghost}},
\]
\[
T_{00} = T_{11} = \frac{1}{2} \left( 4\pi^2 \alpha' P_K P_L \eta^{KL} + \frac{1}{\alpha'} X'^i K X'^j L \eta_{ij} \right),
\]
\[
T_{10} = T_{01} = 2\pi P_K X'^i K, X'^j K = \partial_\sigma X^K,
\]
(3.14)
where \(Q_{\text{ghost}}\) is the ghost contribution to the BRST charge whose explicit form we will not require. Using this form of the BRST operator we get
\[
D = [Q, K] = \frac{1}{\pi} \int_0^\pi d\sigma c^0(\sigma) \frac{1}{2} \left[ 4\pi^2 \alpha' 2\varepsilon P_i P_j \eta^{ij} - 2\varepsilon \frac{1}{\alpha'} X'^i X'^j \eta_{ij} \right].
\]
(3.15)
Consequently the new BRST operator has the form
\[
Q' = Q - D = \frac{1}{\pi} \int_0^\pi d\sigma \left\{ c^0(\sigma) \frac{1}{2} \left[ 4\pi^2 \alpha' P_\mu P_\nu \eta^{\mu\nu} + \frac{1}{\alpha'} X'^i X'^j \eta_{ij} + 4\pi^2 \alpha' (1 - 2\varepsilon) P_i P_j \eta^{ij} + (1 + 2\varepsilon) \frac{1}{\alpha'} X'^i X'^j \eta_{ij} \right] + \right.
\]
\[
\left. + c^1(\sigma) 2\pi X'^i K P_K(\sigma) \right\} + Q_{\text{ghost}}.
\]
(3.16)
We see from (3.16) that the numerical factors in front of \(P_i P_j, X'^i X'^j\) now depend on the parameter \(\varepsilon\). There is a suggestive physical interpretation of this phenomena which transpires after an analysis of the case of finite \(\varepsilon\). To this end, we calculate exact form of the new BRST operator given in (2.8). To do this calculation we define the function
\[
F(t) = e^{-\varepsilon K t} \left( Q(e^{\varepsilon K t}) \right), \quad F(t = 0) = Q, F(t = 1) = \tilde{Q}.
\]
(3.17)
It is easy to see that
\[
\tilde{Q} = F(0) + \sum_{n=1}^\infty \frac{1}{n!} \frac{d^n F}{d^n t} \bigg|_{t=1} = Q + \sum_{n=1}^\infty \frac{(-1)^n \varepsilon^n}{n!} [K, [K, \ldots, [K, Q]]].
\]
(3.18)
so to determine (3.18) we must calculate all commutators in (3.18). The second order commutator is equal to
\[
[K, [K, Q]] = -[K, D] = [D, K] =
\]
\[
= \frac{1}{2\pi} \int_0^\pi d\sigma c^0(\sigma) \frac{1}{2} \left( 4\pi^2 \alpha' 2\varepsilon P_i P_j \eta^{ij} + 2\varepsilon \frac{1}{\alpha'} X'^i X'^j \eta_{ij} \right),
\]
(3.19)
where we have used (3.17) and (3.4). We conclude that the new BRST operator will have the form
\[
Q' = \frac{1}{\pi} \int_0^\pi d\sigma \left\{ c^0(\sigma) \frac{1}{2} \left[ 4\pi^2 \alpha' P_\mu P_\nu \eta^{\mu\nu} + \frac{1}{\alpha'} X'^{\mu} X'^{\nu} \eta_{\mu\nu} \right] + \ight.
\left. + (1 - 2\epsilon + \frac{1}{2}(2\epsilon)^2 + \ldots) 4\pi^2 \alpha' P_i P_j \eta^{ij} + (1 + 2\epsilon + \frac{(2\epsilon)^2}{2} + \ldots) \frac{1}{\alpha'} X'^{R} X'^{ij} \eta_{ij} \right\} + c^1(\sigma) 2\pi X'^{iK} P_k(\sigma) \bigg) + Q_{\text{ghost}} =
\left. + e^{-2\epsilon} 4\pi^2 \alpha' P_i P_j \eta^{ij} + e^{2\epsilon} \frac{1}{\alpha'} X'^{ni} X'^{ij} \eta_{ij} \bigg] + c^1(\sigma) 2\pi X'^{iK} P_k(\sigma) \bigg) + Q_{\text{ghost}} . \right.
\]

(3.20)

In the next section we try to interpret this solution.

### 4. Interpretation of the solution

In this section we interpret the solution (3.20). As we have seen, the string field theory action expanded around the new solution has the form
\[
S = -\frac{1}{2} \int \Psi \ast Q' \Psi - \frac{1}{3} \int \Psi \ast \Psi \ast \Psi \tag{4.1}
\]
with
\[
Q' = \frac{1}{\pi} \int_0^\pi d\sigma \left\{ c^0(\sigma) \frac{1}{2} \left[ 4\pi^2 \alpha' P_\mu P_\nu \eta^{\mu\nu} + \frac{1}{\alpha'} X'^{\mu} X'^{\nu} \eta_{\mu\nu} \right] + \ight.
\left. + \left( e^{-2\epsilon} 4\pi^2 \alpha' P_i P_j \eta^{ij} + e^{2\epsilon} \frac{1}{\alpha'} X'^{ni} X'^{ij} \eta_{ij} \bigg] + \ight.
\left. + c^1(\sigma) 2\pi X'^{iK} P_k(\sigma) \bigg) + Q_{\text{ghost}} . \right.
\]

(4.2)

In what follows we will discuss mainly the shifted part of the BRST operator. We also introduce the parameter
\[
e^\epsilon = t \tag{4.3}
\]
We see that for \( t = 1 \) we regain the original action. For \( t^2 \rightarrow 0 \) we rewrite the BRST operator as
\[
Q' = \frac{1}{\pi} \int_0^\pi d\sigma \left\{ c^0(\sigma) \frac{1}{2} \left[ e^{-2\epsilon} 4\pi^2 \alpha' P_i P_j \eta^{ij} + e^{2\epsilon} \frac{1}{\alpha'} X'^{ni} X'^{ij} \eta_{ij} \bigg) + c^1(\sigma) 2\pi X'^{ni} P_i(\sigma) \bigg) = \ight.
\left. = \frac{1}{\pi} \int_0^\pi d\sigma \left\{ c^0(\sigma) \frac{1}{2} \left[ 4\pi^2 \alpha' P_i P_j \eta^{ij} + \frac{1}{\alpha'} X'^{ni} X'^{ij} \eta_{ij} \bigg) + c^1(\sigma) 2\pi X'^{ni} P_i(\sigma) \bigg) = \ight.
\left. = \frac{1}{\pi} \int_0^\pi d\sigma \left\{ c^0(\sigma) \frac{1}{2} \left[ 4\pi^2 \alpha_{eff} P_i P_j \eta^{ij} + \frac{1}{\alpha_{eff}} X'^{ni} X'^{ij} \eta_{ij} \bigg) + c^1(\sigma) 2\pi X'^{ni} P_i(\sigma) \bigg) , \right.
\right.
\]

(4.4)

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where we have introduced
\[ \alpha_{eff}' = \frac{\alpha'}{t^2}. \]  
(4.5)

Now we have
\[ t^2 \to \infty \Rightarrow \alpha_{eff}' \to 0, \]
\[ t^2 \to 0 \Rightarrow \alpha_{eff}' \to \infty. \]  
(4.6)

The first limit corresponds to the zero slope limit in which case the ordinary string action reduces to the action for massless particle. It is remarkable that in our solution this limit can be anisotropic. The second limit corresponds to the tensionless limit [12, 13, 14, 15, 16, 17].

To see this more directly, let us follow [31, 32, 33, 34]. The first limit \( \alpha_{eff}' \to 0 \) is the well known zero slope limit where the string effectively collapses to a point. From (4.4) it is clear that to have a finite expression the terms proportional to \( X' \) should be equal to zero, in other words we consider the zero modes of the field on the world-sheet only. As a result (4.4) reduces to the BRST operator for massless particle.

In the second limit \( \alpha_{eff}' \to \infty \) the term proportional to \( \frac{1}{\alpha_{eff}'} X'^i X'^j \) goes to zero.

We must also rescale \( 2\pi \sqrt{\alpha_{eff}'} P_i = P_i \) to ensure that the energy density is finite. As a result, we obtain the following shifted BRST operator
\[ Q_{shifted} = \frac{1}{\pi} \int_0^\pi d\sigma \left\{ c^0(\sigma) \frac{1}{2} P_i P_j \eta^{ij} \right\}. \]  
(4.7)

which is the BRST operator for an infinite collection of independent massless point-particles as we expect in for tensionless strings [12, 13, 14, 15, 16, 17]. We must also mention that the tensionless limit \( \alpha' \to \infty \) was also studied in very interesting paper [35] with the similar results.

5. Conclusion and open questions

As we have seen we can easily find solutions of the string field theory equation of motion that depends on one single parameter that can be interpreted as the string effective length scale. In our solution this length scale can vary continuously from 0 that corresponds to the zero slope limit in ordinary string theory to \( \alpha_{eff}' = \infty \) corresponding to the tensionless limit of the string theory. Generally, this effective string tension depends on the orientation of the solution, which is a similar effect to that for the effective tension in NCOS string as discussed in [18, 19, 20, 21, 22, 23, 24, 25, 26, 27, 28, 29, 30]. We hope that our result could be helpful for the study the relation between tensile and tensionless strings [12, 13, 14, 15, 16, 17]. It would be...
very interesting to study the conformal field theory on the string world-sheet action with an-isotropic tension. We hope to return to this problem in future.

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References

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

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

[3] 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.

[4] W. Siegel, “Introduction To String Field Theory,” arXiv:hep-th/0107094.

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

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

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

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

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

[10] 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.

[11] I. Kishimoto and K. Ohmori, “CFT description of identity string field: Toward derivation of the VSFT action,” arXiv:hep-th/0112169.

[12] A. Karlhede and U. Lindstrom, Class. Quant. Grav. 3 (1986) L73.

[13] U. Lindstrom, B. Sundborg and G. Theodoridis, “The Zero Tension Limit Of The Superstring,” Phys. Lett. B 253 (1991) 319.

[14] J. Isberg, U. Lindstrom and B. Sundborg, “Space-time symmetries of quantized tensionless strings,” Phys. Lett. B 293, 321 (1992) arXiv:hep-th/9207005.
[15] J. Isberg, U. Lindstrom, B. Sundborg and G. Theodoridis, “Classical and quantized tensionless strings,” Nucl. Phys. B 411, 122 (1994) [arXiv:hep-th/9307108].

[16] B. Sundborg, “Stringy gravity, interacting tensionless strings and massless higher spins,” Nucl. Phys. Proc. Suppl. 102, 113 (2001) [arXiv:hep-th/0103247].

[17] P. Haggi-Mani and B. Sundborg, “Free large N supersymmetric Yang-Mills theory as a string theory,” JHEP 0004, 031 (2000) [arXiv:hep-th/0002189].

[18] N. Seiberg, L. Susskind and N. Toumbas, “Strings in background electric field, space/time noncommutativity and a new noncritical string theory,” JHEP 0006 (2000) 021 [arXiv:hep-th/0005040].

[19] R. Gopakumar, J. M. Maldacena, S. Minwalla and A. Strominger, “S-duality and noncommutative gauge theory,” JHEP 0006 (2000) 036 [arXiv:hep-th/0005048].

[20] J. Gomis and T. Mehen, “Space-time noncommutative field theories and unitarity,” Nucl. Phys. B 591, 265 (2000) [arXiv:hep-th/0005129].

[21] O. Aharony, J. Gomis and T. Mehen, “On theories with light-like noncommutativity,” JHEP 0009, 023 (2000) [arXiv:hep-th/0006236].

[22] R. Gopakumar, S. Minwalla, N. Seiberg and A. Strominger, “OM theory in diverse dimensions,” JHEP 0008, 008 (2000) [arXiv:hep-th/0006062].

[23] E. Bergshoeff, D. S. Berman, J. P. van der Schaar and P. Sundell, “Critical fields on the M5-brane and noncommutative open strings,” Phys. Lett. B 492, 193 (2000) [arXiv:hep-th/0006112].

[24] J. Gomis and H. Ooguri, “Non-relativistic closed string theory,” J. Math. Phys. 42 (2001) 3127 [arXiv:hep-th/0009181].

[25] U. H. Danielsson, A. Guijosa and M. Kruczenski, “IIA/B, wound and wrapped,” JHEP 0010 (2000) 020 [arXiv:hep-th/0009182].

[26] I. R. Klebanov and J. M. Maldacena, “1+1 dimensional NCOS and its U(N) gauge theory dual,” Int. J. Mod. Phys. A 16, 922 (2001) [Adv. Theor. Math. Phys. 4, 283 (2000)] [arXiv:hep-th/0006085].

[27] J. X. Lu, “(1+p)-dimensional open D(p-2) brane theories,” JHEP 0108 (2001) 049 [arXiv:hep-th/0102050].

[28] U. H. Danielsson, A. Guijosa and M. Kruczenski, “Newtonian gravitons and D-brane collective coordinates in wound string theory,” JHEP 0103 (2001) 041 [arXiv:hep-th/0012183].

[29] F. Kristiansson and P. Rajan, “Wound string scattering in NCOS theory,” Phys. Lett. B 502 (2001) 235 [arXiv:hep-th/0011054].
[30] J. A. Garcia, A. Guijosa and J. D. Vergara, “A membrane action for OM theory,” Nucl. Phys. B 630 (2002) 178 [arXiv:hep-th/0201140].

[31] A. A. Tseytlin, “On limits of superstring in AdS(5) x S**5,” [arXiv:hep-th/0201112].

[32] H. Nastase and W. Siegel, “A new AdS/CFT correspondence,” JHEP 0010, 040 (2000) [arXiv:hep-th/0010106].

[33] U. Lindstrom and H. G. Svendsen, “A pedestrian approach to high energy limits of branes and other gravitational systems,” Int. J. Mod. Phys. A 16, 1347 (2001) [arXiv:hep-th/0007101].

[34] H. G. Svendsen, “High-energy limits of various actions,” [arXiv:hep-th/9911038].

[35] C. S. Chu, P. M. Ho and F. L. Lin, “Cubic string field theory in pp-wave background and background independent Moyal structure,” [arXiv:hep-th/0205218].