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Tachyon Condensation in String Field Theory

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Abstract

It has been conjectured that at a stationary point of the tachyon potential for the D-brane of bosonic string theory, the negative energy density exactly cancels the D-brane tension. We evaluate this tachyon potential by off-shell calculations in open string field theory. Surprisingly, the condensation of the tachyon mode alone into the stationary point of its cubic potential is found to cancel about 70% of the D-brane tension. Keeping relevant scalars up to four mass levels above the tachyon, the energy density at the shifted stationary point cancels 99% of the D-brane tension.

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It has been argued on various general grounds that the classical tachyon potential on a D-p brane of the bosonic string theory has a stationary point where the total negative potential energy due to the tachyon exactly cancels the tension of the D-brane[1, 2]. At this stationary point the configuration is indistinguishible from the vacuum where there is no brane. A similar argument can be given for the tachyon potential on a D-brane anti-D-brane system or a non-BPS D-brane of type II string theories [3, 4, 5, 6, 7, 8]. There is, however, no direct proof of these relations.

In this paper we demonstrate this phenomenon directly using string field theory. We restrict ourselves to bosonic string theory and use Witten's string field theory with cubic action [9], although in principle the version of open string field theory suitable for off shell inclusion of closed strings [10] could have been used as well. The analysis could be performed for superstring theories as well using open superstring field theory[11]. The background independent features of the tachyon potential noted in [12] make earlier studies of this potential in string field theory [13, 14, 15, 16, 17, 18][3] relevant to the problem of D-brane annihilation. Indeed, in a very interesting paper, Kostelecky and Samuel [14] gave evidence that the stationary point of the cubic tachyon potential survives with controllable corrections the inclusion of higher mass scalars of the string field expansion. Further evidence to this effect was given in [18]. It is this non-perturbative vacuum that we focus on in the present paper. Our present advantage is that we have an explicit conjecture for the value of the potential at the stationary point we are looking for. Hence we can compare the results obtained from string field theory with the conjectured value. As we shall see, using a suitable approximation scheme, we can find a stationary point of string field potential where the value of the potential is about 1% away from the conjectured answer.[4]

The conjecture and the setup. Some general properties of the tachyon potential in string field theory were analysed in [12], where it was shown that the tachyon potential on a D-brane of bosonic string theory takes a universal form:

$$V(T) = M f(T)\,, \tag{1}$$

where M is the mass of the D-brane[5] and $f(T)$ is a universal function independent of

[3]For studies of the tachyon potential in the first quantized formulation, see [19, 20].

[4]Our calculations appear to be consistent with the related computations of [14, 18].

[5]We are assuming that all the directions tangential to the brane are compact so that the brane has a finite mass.

the background in which the D-brane is embedded. The tachyon field T and the function $f(T)$ are defined as follows[12]. Let $\mathcal{H}$ denote the space of states of ghost number one of the two-dimensional conformal field theory of the (b, c) ghost system and a matter system of central charge 26 on the upper half plane, and let $|0\rangle$ denote the SL(2,R) invariant vacuum of this conformal field theory. Let $\mathcal{H}_1 \subset \mathcal{H}$ denote the space of states of ghost number one obtained by acting on $|0\rangle$ with oscillators b_n, c_n, and matter Virasoro generators L_n. The subspace $\mathcal{H}_1$ of $\mathcal{H}$ is a background independent subspace containing the zero momentum tachyon state $c_1|0\rangle$ and having the property that we can consistently set the component of the string field along $\mathcal{H} - \mathcal{H}_1$ to zero in looking for a solution of the equations of motion. Since all fields in $\mathcal{H}_1$ may acquire expectation values, the real problem is finding a stationary point of the string field potential $V(T)$ associated to the string field $|T\rangle$ corresponding to a general state in $\mathcal{H}_1$. This string field $|T\rangle$, still called here the tachyon field, includes an infinite collection of variables corresponding to the coefficients of expansion of a state in $\mathcal{H}_1$ in some basis.[6]

The function $f(T)$ is given by the following string field theory expression:

$$f(T) = 2\pi^2 \Big(\frac{1}{2}\langle I \circ \mathcal{T}(0) Q_B \mathcal{T}(0)\rangle + \frac{1}{3}\langle h_1 \circ \mathcal{T}(0) h_2 \circ \mathcal{T}(0) h_3 \circ \mathcal{T}(0)\rangle \Big) . \qquad (2)$$

Here Q_B is the BRST charge, and $\mathcal{T}(z)$ denotes the two dimensional field which creates the state $|T\rangle$ from the SL(2,R) invariant vacuum: $|T\rangle = \mathcal{T}(0)|0\rangle$. I, h_1, h_2 and h_3 are a set of familiar conformal transformations[21] whose expressions were reviewed in [12]. Given any conformal transformation described by the function $h(z)$, and a vertex operator $\Phi(z)$ of the conformal field theory, $h \circ \Phi(0)$ denotes the conformal transform of $\Phi(0)$ by h. Thus for example if Φ denotes a dimension d primary field, then $h \circ \Phi(0) = (h'(0))^d \Phi(h(0))$. For non-primary fields there will be extra terms involving higher derivatives of h. Finally $\langle\ \rangle$ denotes the correlation function in the conformal field theory of matter and ghost fields, normalized so that $\langle c_{-1} c_0 c_1\rangle = 1$.[7,8]

[6]Although we shall refer to these coefficients as fields, we should keep in mind that these represent zero momentum modes of the fields corresponding to space-time independent field configurations.

[7]Note that this differs from the convention of ref.[12] by a factor of L, – the (infinite) length of the time interval. This is due to the fact that we are writing down the expression for the potential instead of the action. We can make these two notations consistent by choosing $L = 1$; in that case the potential can be identified to the negative of the action. The final results of course are independent of L.

[8]The factor of $2\pi^2$ in eq.(2) arises as follows. With $S(\Phi) = -\frac{1}{g_0^2}\Big(\frac{1}{2}\langle \Phi, Q_B \Phi\rangle + \cdots\Big)$, the D-brane mass is : $M = 1/(2\pi^2 g_0^2)$[12]. Then $V(T) = -S(T) = M \cdot (2\pi^2)\Big(\frac{1}{2}\langle T, Q_B T\rangle + \cdots\Big) = M f(T)$.

The conjecture of ref.[1] can now be restated as follows. In the space $\mathcal{H}_1$ there must be a state $|T_c\rangle$ such that $f(T)$ has a stationary point at $T = T_c$, and

$$f(T_c) = -1\,. \tag{3}$$

The total D-brane mass at $T = T_c$ vanishes: $M + V(T_c) = M(1 + f(T_c)) = 0$.

Zeroth Approximation. We proceed to verify this conjecture using a systematic approximation scheme suggested by Kostelecky and Samuel[14]. In order to explain this procedure, let us first consider setting all components of $|T\rangle$ to zero except for the coefficient of the state $c_1|0\rangle$, a state that will be said to be of *level zero*. Thus we take

$$|T\rangle = t\, c_1|0\rangle\,. \tag{4}$$

Substituting this into eq.(2) we get the zeroth approximation to the tachyon potential

$$f^{(0)}(t) = 2\pi^2\Big(-\frac{1}{2}t^2 + \frac{1}{3}\,\frac{t^3}{r^3}\Big)\,, \qquad r = \frac{4}{3\sqrt{3}}. \tag{5}$$

This has a local minimum at

$$t = t_c \equiv r^3 = \Big(\frac{4}{3\sqrt{3}}\Big)^3 \simeq 0.456\,. \tag{6}$$

At this minimum

$$f(t_c) = -2\pi^2\cdot\frac{1}{6}\cdot r^6 = -\frac{\pi^2}{3}\cdot\Big(\frac{4}{3\sqrt{3}}\Big)^6 = -\frac{4096}{59049}\,\pi^2 \simeq -0.684\,. \tag{7}$$

We found it very encouraging that this zeroth order approximation to the vacuum energy at the stationary point gives essentially 70% of the expected value! In fact, the off-shell choice of cubic string field theory (as opposed to string field theory with higher order vertices) yields at this level the best possible approximation to the expected value. Indeed, the constant r defined above is essentially the mapping radius of the disks defining the three string vertex [16], and it is maximal for the vertex of the cubic theory. Thus $|f(t_c)|$ is maximal for this choice.

Subsidiary conditions on T. In order to compute corrections to this result, we need to include the higher level fields in our analysis. The analysis can be simplified by noting that the potential (2) has a twist symmetry under which all coefficients of states at odd levels above $c_1|0\rangle$ change sign, whereas coefficients of states at even level above $c_1|0\rangle$

remain unchanged [14, 22].[9] Thus coefficients of states at odd levels above $c_1|0\rangle$ must always enter the action in pairs, and we can trivially satisfy the equations of motion of these fields by setting them to zero. Thus we look for solutions where $|T\rangle$ *contains only even level states*. With the state $c_1|0\rangle$ defined to be at level zero, the additional fields we must consider will be at levels two, four, and higher. At level two, for example, we find three states, $c_{-1}|0\rangle$, $L_{-2}c_1|0\rangle$ and $b_{-2}c_0c_1|0\rangle$.

We can further simplify the expansion by using the Feynman-Siegel gauge:

$$b_0|T\rangle = 0\,. \tag{8}$$

This gauge choice can be justified by first showing that such a gauge can be chosen at the linearized level, and then assuming that the fields are small enough so that we can continue to make this gauge choice even in the presence of interactions. The proof of validity of this gauge at the linearized level proceeds as follows. Let $|T^{(2n)}\rangle$ denote an arbitrary level $2n$ state in $\mathcal{H}_1$. Let us define $|\Lambda^{(2n)}\rangle = b_0|T^{(2n)}\rangle$. Then $|\widetilde{T}^{(2n)}\rangle \equiv |T^{(2n)}\rangle - (2n-1)^{-1}Q_B|\Lambda^{(2n)}\rangle$ satisfies the desired gauge condition $b_0|\widetilde{T}^{(2n)}\rangle = 0$.[10] This shows that for $n \geq 1$, it is possible to gauge transform a general level $2n$ state $|T^{(2n)}\rangle$ to a state $|\widetilde{T}^{(2n)}\rangle$ satisfying the Feynman-Siegel gauge.

The equations of motion in the Feynman-Siegel gauge are equivalent to the equations of motion of the gauge invariant action if there are no residual gauge transformations which act non-trivially and preserve the gauge, *i.e.* if there are no pure gauge directions inside the subspace (8). Assume there is pure gauge direction $|\eta^{(2n)}\rangle$ satisfying the gauge condition (8). Then, $Q_B|\eta^{(2n)}\rangle = 0$, and together with the gauge condition (8), gives $L_0^{tot}|\eta^{(2n)}\rangle = \{Q_B, b_0\}|\eta^{(2n)}\rangle = 0$. Since $L_0^{tot}|\eta^{(2n)}\rangle = (2n-1)|\eta^{(2n)}\rangle$, we see that $|\eta^{(2n)}\rangle$ must vanish, as we wanted to show. Hence the Feynman-Siegel gauge is a valid gauge choice for sufficiently small field configurations.

<u>Approximation with level two fields</u>. Using the Feynman-Siegel gauge we have

$$|T\rangle = tc_1|0\rangle + uc_{-1}|0\rangle + v\cdot\frac{1}{\sqrt{13}}L_{-2}c_1|0\rangle\,, \tag{9}$$

[9] The origin of this symmetry can be traced to the relations $h_1(-z) = \widetilde{I}(h_3(z))$, $h_2(-z) = \widetilde{I}(h_2(z))$, and $h_3(-z) = \widetilde{I}(h_1(z))$, where $\widetilde{I}(y) = 1/y$ is a combination of the SL(2,R) and world-sheet parity transformations. In the restricted sector $\mathcal{H}_1$ the world-sheet parity as well as SL(2,R) is a symmetry of the theory.

[10] In deriving the above equation we have used that (a) $\{Q_B, b_0\} = L_0^{tot}$, with L_n^{tot} denoting the combined Virasoro generators of the matter and the ghost sectors, (b) $b_0|\Lambda^{(2n)}\rangle = 0$ due to the relation $(b_0)^2 = 0$, and that (c) the L_0^{tot} eigenvalue of a level $2n$ state is $(2n-1)$.

which includes the two level two fields u and v. The $(1/\sqrt{13})$ factor in the last term was chosen for convenience (L_n's denote matter Virasoro generators.). At this stage we can simply substitute (9) into (2) and find $f(t,u,v)$, but there is a further approximation which is possible[14]. For this let us define the level of a given term in $f(T)$ as the sum of the levels of all the fields appearing in this term. We can now approximate the potential $f(T)$ by keeping only terms up to a certain level. Since the quadratic terms involving level two fields is already level four, it does not make sense to truncate the potential to terms below level four once we have included the level two fields in our analysis. Thus the next approximation to the potential, $f^{(4)}$, will be obtained by substituting in (2) the expansion (9) and keeping only t^3, t^2u, t^2v, tu^2, tv^2 and tuv interaction terms. Since we cannot have fields appear in interactions before their quadratic terms appear, we define *the level $2n$ approximation $f^{(2n)}$ to contain all interaction terms up to level $2n$ built from fields up to level n.* Thus at level six, we do not include any new fields (odd level fields are set to zero) but we need to include four new interactions: u^3, v^3, u^2v and uv^2.

Ref.[14] evaluated the potential up to level six and found that the stationary point of the potential persists up to this level. Their result, when translated into the normalization convention of this paper, is as follows.[11] At level four the potential is given by

$$\begin{aligned} f^{(4)}(T) &= 2\pi^2\Big(-\frac{1}{2}\,t^2 + \frac{3^3\sqrt{3}}{2^6}\,t^3 \\ &\quad -\frac{1}{2}\,u^2 \;+\; \frac{1}{2}\,v^2 \;+\frac{11\cdot 3\sqrt{3}}{2^6}\,t^2\,u - \frac{5\cdot 3\sqrt{39}}{2^6}\,t^2\,v \\ &\quad +\frac{19}{2^6\,\sqrt{3}}\,t\,u^2 + \frac{7\cdot 83}{2^6\cdot 3\sqrt{3}}\,t\,v^2 - \frac{11\cdot 5\sqrt{13}}{2^5\cdot 3\sqrt{3}}\,t\,u\,v\Big)\,. \end{aligned} \tag{10}$$

$f^{(4)}(T)$ has a stationary point[12] at T_c ($t_c \simeq 0.542$, $u_c \simeq 0.173$, $v_c \simeq 0.187$) at which $f(T_c) \simeq -0.949$. This is about 95% of the expected answer -1![13]

[11] In order to convert the result of ref.[14] to our convention, we need to set $\alpha' = 1$ and $g = 2$ in the expressions given in ref.[14], and multiply their potential by a factor of $2\pi^2$. The fields t, u and v used in our paper correspond to their fields ϕ, $-\beta_1$ and B respectively.

[12] Strictly, this point is not a minimum of the potential, nor even a local minimum. This is not necessarily problematic. The string field theory has ghost and auxiliary modes with negative mass squared that are not physical tachyons. The field u appearing in (10) is an example of this. We still expect physical stability of this stationary point.

[13] Since ref.[14] did not have a reference scale to compare with, they expressed the potential in units of the string tension and the on-shell three tachyon coupling, and concluded that the potential is quite shallow. On the other hand, using the mass of the D-brane as the reference scale, we see that the potential is in fact quite deep. Already at this level it is about 95% of the mass of the D-brane.

At level six the potential includes the level four interactions plus additional terms:

$$f^{(6)} = f^{(4)} + 2\pi^2 \Big(\frac{1}{2^6\sqrt{3}}\, u^3 \,-\, \frac{7\cdot 41\cdot 73}{3^4\cdot 2^6\sqrt{39}}\, v^3 - \frac{5\cdot 19\sqrt{13}}{2^6\, 3^3\sqrt{3}}\, u^2\, v + \frac{11\cdot 7\cdot 83}{2^6\cdot 3^4\sqrt{3}}\, u\, v^2 \Big). \quad (11)$$

Solving the equations of motion that follows from the total level six potential $f^{(6)}$, one finds that the location of T_c is shifted slightly ($t_c \simeq 0.544$, $u_c \simeq 0.190$, $v_c \simeq 0.202$) with $f(T_c) \simeq -0.959$. By including two modes in addition to the tachyon we have gone from 68% to 96% of the expected vacuum energy! This is certainly encouraging and leads us to believe that the expansion converges rapidly to the expected answer.

Approximation with level eight interactions. To establish the convergence beyond reasonable doubt, we now undertake the substantially more involved calculation of the potential to level eight. For this we need to include all the level four fields. A general tachyon field configuration in the Feynman-Siegel gauge, including fields up to level four, has the form:[14]

$$\begin{aligned} |T\rangle \;\; = \;\; & t c_1|0\rangle + u c_{-1}|0\rangle + v\cdot \frac{1}{\sqrt{13}} L_{-2}c_1|0\rangle \\ & + A L_{-4}c_1|0\rangle + B L_{-2}L_{-2}c_1|0\rangle + C c_{-3}|0\rangle \\ & + D b_{-3}c_{-1}c_1|0\rangle + E b_{-2}c_{-2}c_1|0\rangle + F L_{-2}c_{-1}|0\rangle\,. \end{aligned} \quad (12)$$

In order to construct the potential to level eight, we need to substitute (12) into (2), and evaluate the action keeping terms up to level eight. We use two different methods to compute the cubic interaction vertices. In the first approach we explicitly compute the conformal transformation of all the vertex operators associated with the state (12) under the conformal maps h_1, h_2 and h_3, and compute the three point correlation functions of the resulting operators. In the second approach we use a representation of the matter Virasoro algebra in terms of 26 free bosonic fields, and use the Neumann function method to compute the three string vertex [23]. Both approaches give the same results.

Besides the terms given in eqs.(10) and (11) (which we explicitly verify), there are four different kinds of additional terms in the computation of $f^{(8)}$. These are

1. The quadratic term involving the level 4 fields. This is a level 8 contribution to the potential, and is given by:

$$\Delta_0 f^{(8)} = 2\pi^2 \,\Big(\, 195\, A^2 + 663\, B^2 \,+\, 234\, AB \,+\, 3\, CD - \frac{3}{2}\, E^2 - \frac{39}{2}\, F^2 \Big)\,. \quad (13)$$

[14] Since we are using background independent modes, we have here less fields than in Refs.[14, 18], the last of which cites a computation of the potential using level six fields.

2. There are new level four interaction terms, originating from the coupling between two level 0 and a level 4 field. These are given by,

$$\Delta f^{(4)} = 2\pi^2 \frac{1}{\sqrt{3}} t^2 \Big(\frac{585}{32}A + \frac{3523}{96}B - \frac{5}{12}C + \frac{5}{4}D + \frac{19}{64}E - \frac{715}{192}F \Big) . \qquad (14)$$

3. There are new level six interaction terms, originating from the coupling between a level 0, a level 2 and a level 4 field. These are given by

$$\Delta f^{(6)} = \frac{2\pi^2}{\sqrt{3}} t \Big[u \Big(\frac{715}{48}A + \frac{38753}{1296}B - \frac{25}{54}C + \frac{25}{18}D + \frac{3827}{2592}E - \frac{1235}{864}F \Big) + \sqrt{13}\, v \Big(- \frac{7495}{1296}A - \frac{12101}{432}B + \frac{25}{162}C - \frac{25}{54}D - \frac{95}{864}E + \frac{6391}{2592}F \Big) \Big] . \qquad (15)$$

4. There are cubic interaction terms of level 8. These involve coupling of two level 2 fields with a level 4 field, and also the coupling of two level 4 fields with a level 0 field. The 2-2-4 interaction terms are given by,

$$\Delta_1 f^{(8)} = \frac{2\pi^2}{\sqrt{3}} \Big[u^2 \Big(\frac{1235}{864}A + \frac{66937}{23328}B - \frac{5}{108}C + \frac{5}{36}D + \frac{124849}{139968}E - \frac{65}{576}F \Big) + \sqrt{13}\, uv \Big(- \frac{82445}{34992}A - \frac{133111}{11664}B + \frac{125}{1458}C - \frac{125}{486}D - \frac{19135}{69984}E + \frac{11039}{23328}F \Big) + v^2 \Big(\frac{254381}{23328}A + \frac{1598597}{23328}B - \frac{2905}{8748}C + \frac{2905}{2916}D + \frac{11039}{46656}E - \frac{230461}{46656}F \Big) \Big] . \qquad (16)$$

On the other hand, the 0-4-4 interaction terms are given by,

$$\Delta_2 f^{(8)} = \frac{2\pi^2}{\sqrt{3}} t \Big(\frac{3539315}{11664}A^2 + \frac{9440977}{17496}AB + \frac{4367233}{3888}B^2 - \frac{325}{81}AC - \frac{17615}{2187}BC + \frac{25}{729}C^2 + \frac{325}{27}AD + \frac{17615}{729}BD + \frac{1598}{2187}CD + \frac{25}{81}D^2 + \frac{1235}{432}AE + \frac{66937}{11664}BE + \frac{665}{1458}CE - \frac{665}{486}DE - \frac{4061}{5184}E^2 - \frac{1071785}{34992}AF - \frac{1730443}{11664}BF + \frac{1625}{1458}CF - \frac{1625}{486}DF - \frac{248755}{69984}EF + \frac{143507}{46656}F^2 \Big) . \qquad (17)$$

The level 8 approximation to the full tachyon potential is obtained by combining the contributions (10), (11) and (13)-(17):

$$f^{(8)} = f^{(6)} + \Delta f^{(4)} + \Delta f^{(6)} + \Delta_0 f^{(8)} + \Delta_1 f^{(8)} + \Delta_2 f^{(8)} \,. \tag{18}$$

Given this potential we can search for a stationary point of the potential.[15] We find that the equations of motion following from this potential are satisfied for[16]:

$$\begin{aligned} t_c = 0.5482, \quad u_c = 0.2043, \quad v_c = 0.2045, \quad A_c = -0.00495, \quad B_c = -0.00056, \\ C_c = -0.0549, \quad D_c = 0.0183, \quad E_c = 0.0317, \quad F_c = -0.0066 \,. \end{aligned} \tag{19}$$

The value of the potential at this stationary point is given by

$$f(T_c) = -0.9864 \,. \tag{20}$$

This is about 1% away from the expected answer -1. Note the near equality of the values of u_c and v_c. We suspect this to be an exact equality in the complete answer. This, in turn, indicates that there might be a closed form expression for the state $|T_c\rangle$ describing the stationary point, since otherwise it will be very difficult to explain the equality of these coefficients.

Discussion. We shall end this paper by discussing the significance of our results and some further investigations which these results suggest.

1. Our result indicates that the tachyonic vacuum of the bosonic D-brane, representing its annihilation, is described by a string field dominated by the low lying modes of the theory. This is certainly surprising since total brane disappearance is a highly non-perturbative phenomenon and one could have expected non-trivial participation of all the higher string modes. As we saw, however, condensation of states up to level four account for 99% of the potential energy required to cancel the tension of the D-brane. Our results show that string field theory captures non-perturbative string dynamics. The string field $|T_c >$ appears to be well-defined.

[15] The actual computations require use of symbolic manipulation programs Maple and Mathematica.

[16] This stationary point is closest to the one given earlier for level four and level six potentials, in the sense that if we look for a numerical solution with the stationary point of the level four or level six potential as the starting point, we arrive at the solution (19). There may be other solutions to the equations of motion whose interpretation is not clear.

2. Associated with the phenomenon of tachyon condensation on the D-brane is the problem of the extra U(1) gauge field [24, 6, 25]. How does the U(1) gauge field living on the brane disappear after the tachyon condenses and the brane annihilates? Since the tachyon is neutral, the gauge field cannot acquire mass via the Higgs mechanism. Also, how do open strings with one endpoint lying on the D-brane in question disappear after tachyon condensation?

   Ref.[26] proposed that at the extrema $T = T_c$, the action of the gauge field vanishes identically. This explains the absence of a dynamical gauge field. In addition, the path integral over the gauge field now sets to zero the charged currents, thus explaining the absence of open strings with one endpoint on the brane in question.

   Since our analysis shows that the tachyonic vacuum can be studied efficiently with string field theory, one can ask if the above proposal can also be verified using string field theory. In other words, can one study the fate of action involving the gauge fields at the extremum $T = T_c$ and show that its coefficient becomes small?

3. According to the conjecture of ref.[1, 2], a D$(p-1)$ brane of the bosonic string theory can be regarded as a lump solution on a Dp brane, where far away from the core of the lump the tachyon condenses to the critical value T_c. Since the configuration $T = T_c$ seems to have a good description in string field theory, it is natural to ask whether the lump also has a good description in string field theory.

4. Another question that arises from our analysis is: is it possible to write down a closed form expression for the exact extremum $|T_c\rangle$ of the tachyon potential, and/or of the lump solutions describing lower dimensional branes? As we have already mentioned, the near equality of u_c and v_c in eq.(19) can be taken as an evidence that there is a closed form expression for $|T_c\rangle$.

5. Finally, we can wonder about the existence of a stationary point in the tachyon potential for the bosonic closed string field theory [27]. Could this vacuum, if it exists, be a state of unbroken general coordinate invariance having no dynamical graviton? While there appears to be no physical prediction for such hypothetical stationary state, the methods discussed here may improve on earlier computations [15, 16, 17] to give some new insight into this problem.

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