D-brane decay and Hawking Radiation*

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Tree level decay amplitudes of near-BPS D-brane configurations are known to exactly reproduce Hawking radiation rates from corresponding black holes at low energies even though the brane configurations describe semiclassical black holes only when the open string couplings are large. We show that a large class of one (open string) loop corrections to emission processes from D-branes vanish at low energies and nonvanishing loop contributions have an energy dependence consistent with black hole answers, thus providing a justification for the agreement of the tree level results with semiclassical answers.

1. Introduction

Recently the idea that massive string states become black holes when the coupling is large [1–3] has been very successful. In particular, five dimensional extremal black holes with large horizons are described by bound states of D-branes whose degeneracies exactly reproduce the Beckenstein-Hawking entropy [4] This result was extended to four dimensional extremal black holes [5] and to spinning black holes [6].

The entropy and hawking temperature continue to agree in the near extremal limit [7,8]. It was also found that the lowest order decay rate of slightly nonextremal D-brane configurations is proportional to the horizon area [7] which is consistent with the semiclassical Hawking radiation from such holes, as shown in [9]. Rather surprisingly the two emission (and absorption) rates for neutral scalars were in fact found to agree *exactly* [10]. Exact agreements were also found for neutral and charged scalars in five and four dimensional black holes [11]. Even more remarkably, the grey body factors which describe a nontrivial energy dependence of the absorption cross-section at higher energies are also in exact agreement [12], a result verified in the four dimensional case as well [13].A general analysis of classical absorption by such black holes and the possibility of agreement of classical and D-brane greybody factors, has been carried out in [14].

A more detailed test of these ideas is provided by the emission and absorption of certain "fixed" scalars by five dimensional black holes [15], where agreement between the D-brane and classical calculations were demonstrated in [16]. The grey body factors at higher energies agree as well [17].

A related example where there seem to be exact agreement of D-brane and general relativity results is the absorption of $l = 0$ and $l = 1$ waves by extremal 3-branes with no momentum [18,19]. The cross-sections for higher partial waves agree upto numerical factors [19]. In this case the near-extremal entropy also differs by a numerical factor [20].

For systems arbitrarily far from extremality (like the Schwarzschild black hole) it has been argued that the microscopic and semiclassical answers should match only at a special value of the coupling at which the horizon curvature is of the string scale [1,21]. It was shown in [21] that in all known cases the stringy and semiclassical entropies indeed match at this point, upto numerical factors.

2. The puzzle

In a sense these spectacular results are puzzling. In the standard semiclassical description of Hawking radiation the decay rate into particles of energy-momentum (ω, k) is given by

$$\Gamma(\omega) = \frac{\sigma(\omega)}{e^{\beta_H \omega} \pm 1} \frac{d^d k}{(2\pi)^d} \quad (1)$$

where d denotes the number of spatial dimensions and $\sigma(\omega)$ denotes the absorption cross-section of waves of frequency ω by the black hole. $\sigma(\omega)$ thus encodes the space-time structure of the black hole.

On the other hand in D-brane perturbation theory Hawking radiation appears as a result of transitions between solitonic states in *flat* space-time, usually in the lowest order in string perturbation theory. Thermality of the radiation is due to a large degeneracy of initial states of the brane system.

These are quite different pictures. Yet at low energies the answers for $\Gamma(\omega)$ agree in the cases mentioned above.

In fact the semiclassical black hole picture and the perturbative D-brane picture are descriptions of the same object in two quite different regimes. D-brane states are expected to describe semiclassical black holes when (gQ) is large, where g is the string coupling and Q is a typical charge of the hole. This product (gQ) is in fact the open string coupling. The D-brane calculations are, however, performed at weak open string coupling. For extremal BPS states there are well-known nonrenormalization theorems which ensure that the degeneracy of states do not change as we increase the coupling. But for non-BPS states there are no such obvious theorems.

To appreciate the point consider the metric for the five dimensional near extremal black hole which is described by bound states of 1D-branes and 5D-branes with some momentum flowing along the 1D brane. Non extremality is introduced by allowing both left and right momenta on the 1D brane and the classical solution of the low energy effective action has a ten-dimensional string metric

$$\begin{aligned} ds^2 &= (1+\frac{r_1^2}{r^2})^{-1/2}(1+\frac{r_5^2}{r^2})^{-1/2}[dt^2 - dx_5^2 \\ &\quad -\frac{r_0^2}{r^2}(\cosh\sigma\ dt + \sinh\sigma\ dx_5)^2] \\ &\quad -(1+\frac{r_1^2}{r^2})^{1/2}(1+\frac{r_5^2}{r^2})^{-1/2}[dx_1^2 + dx_2^2 \\ &\quad +dx_3^2 + dx_4^2] \\ &\quad -(1+\frac{r_1^2}{r^2})^{1/2}(1+\frac{r_5^2}{r^2})^{1/2}[(1-\frac{r_0^2}{r^2})^{-1}dr^2 \\ &\quad +r^2 d\Omega_3^2] \end{aligned} \quad (2)$$

The various length scales are given in terms of the charges by

$$\begin{aligned} r_1^2 &= \frac{16\pi^4 \alpha'^3 (gQ_1)}{V} \\ r_5^2 &= \alpha'(gQ_5) \\ \frac{1}{2} r_0^2 \sinh 2\sigma &= \frac{16\pi^4 \alpha'^4 (g^2 N)}{R^2 V} \\ r_N^2 &= r_0^2 \sinh^2\sigma \end{aligned} \quad (3)$$

Here $\alpha' = 1/(2\pi T)$ where T is the elementary string tension, g is the string coupling. The brane configuration lies on a $T^4 \times S^1$ with the one brane along the S^1. The radius of this S^1 is R, while the volume of the T^4 is V. The integers Q_1, Q_5, N are the 1-brane RR charge, 5-brane RR charge and the total momentum. The extremal limit is $r_0 \to 0$ and $\sigma \to \infty$ with N held fixed.

It is clear from the classical solution that the classical limit of the string theory corresponds to $g \to 0$ with $gQ_1, gQ_5, g^2 N$ held fixed [12]. In fact we have large black holes (compared to string scale) when $gQ_1, gQ_5, g^2 N > 1$ and small holes when $gQ_1, gQ_5, g^2 N < 1$. It is in the latter regime that the D-brane description in terms of a bound state of 1D and 5D branes with some momentum along the 1D brane is reliable.

In the dilute gas regime $r_N << r_1, r_5$ the size of the black hole is controlled by (gQ_1) and (gQ_5) which are the effective open string coupling constants. The full classical solution can be obtained by summing over an infinite number of string diagrams which does not contain any closed string loop, but contains all terms with closed strings terminating on an aribtrary number of branes. Each such insertion carries a factor gQ which has to be held finite. In other words we have to sum over all open string loops. Closed string loops do not contain any factor of the charge Q and are therefore suppressed. This perturbation expansion is a description of the black hole expanded around flat space-time with the curvature emerging as a result of summing over open string loops.

Another example which we will consider in the following is the self-dual 3-brane in Type IIB theory, first considered in this context in [20,18,19]. The extremal solution has a zero horizon area, but is completely *nonsingular* and the dilaton is a constant in the classical solution. The extremal string metric is given by

$$ds^2 = A^{-1/2}(-dt^2 + dx_1^2 + dx_2^2 + dx_3^2) + A^{1/2}(dr^2 + r^2 d\Omega_5^2) \quad (4)$$

with

$$A(r) = 1 + \frac{4\pi\alpha'^2(gN)}{r^4} \quad (5)$$

where N is the RR charge. The curvature at the horizon $r = 0$ is $\sim 1/[\alpha'\sqrt{gN}]$. Thus when gN is large, the curvatures are small and one may trust the supergravity limit. For $gN << 1$ this system is well described by parallel 3D branes. A great advantage of using extremal 3-branes is that the weak coupling description is well known in terms of Dirichlet open string theory, while for the five or four dimensional black holes not much is known about the properties of the bound states.

The question is : why is it that the absorption or emission cross sections calculated in tree level open string theory agree in detail with the semi-classical black hole answers. Why is it that open string loops do not alter the result.

In this talk I will summarize some results which seem to offer an answer to this question. Some of these results are described in [22]. The results of Section 7 are however new and not published anywhere else. For further results concerning loop corrections see the contribution of Klebanov to these proceedings [23].

3. The issue of loop corrections

A little thought shows that the situation is not as puzzling as it first appears. Consider for example absorption by extremal black holes which have a single length scale in the problem. Examples are fat black holes with $r_1 = r_5 = r_N = R$ or extremal parallel branes, like 3-branes. Let us call this length scale l. In the classical solution the string coupling can enter only through this length scale l which is typically given by the form

$$l^{(d-3)} \sim gQ\alpha'^{(d-3/2)} \quad (6)$$

where d denotes the number of non-compact dimensions. It is then clear that the classical absorption cross-section has to be of the form

$$\sigma_{class} \sim l^{d-2} F(\omega^{(d-3)} gQ\alpha'^{(d-3/2)}) \quad (7)$$

On general grounds we expect that this classical answer should agree with the D-brane answer when gQ is large. However the above expression shows that for sufficiently small ω one may have the factor $\omega^{(d-3)}gQ\alpha'^{(d-3/2)}$ small even if gQ is large so that one may imagine performing a Taylor expansion of the function F, which then becomes a power series expansion in the string coupling g as well [18]. The spectacular success of the tree level D-brane calculations of the absorption cross-section then means that the *lowest order* term in this expansion has been shown to agree with the *lowest order* term in D-brane open string perturbation theory.

The puzzle regarding this agreement of D-brane and classical calculations may be now restated as follows : In the classical limit a higher power of the string coupling comes with a higher power of the energy in a specific way dictated by (7). On the other hand, on the D-brane side these higher powers of coupling are to be obtained in open string perturbation theory and there is no *a priori* reason why this should also involve higher powers of energy in precisely the same way.

This implies that the following two kinds of loop corrections must be absent for the correspondence to work. (1) For a given absorption or emission process open string loop diagrams with the same external states must be suppressed at low energies compared to the tree diagram. (2) Suppose we concentrate on emission of some given closed string state and let the leading order tree process give a cross-section $\sigma \sim g^{\alpha}$ with some energy dependence. For α large enough it is possible that there is a string loop process with the same dependence on the coupling [2]. This must, of course, involve external states which are *different* from the tree process. For processes where

[2] The importance of this was emphasized to me by I. Klebanov.

the D-brane tree level calculation gave the correct answer such loop corrections must be suppressed at low energies.

In fact the black hole correspondence demands more. Nonzero higher loop effects must appear in the precise combination of the string coupling and the energy displayed in (7).

Note that in the D-brane calculations higher powers of ω may appear with the same power of g through grey body factors arising from the thermal distributions accompanying the initial states. Nevertheless the black hole answer tells us how the D-brane perturbation series should look like.

Additional evidence for this picture appears from an application of the correspondence principle of [21] to near-extremal black holes made of parallel $p-D$-branes with no other charge. For such black holes the entropy due to the gas of massless modes on the p-brane worldvolume agrees, upto a numerical factor, with the black hole entropy at the correspondence point defined as the value of the string coupling where the string metric curvature at the horizon becomes of the order of the string scale. Now, it is known that for any spherically symmetric static black hole in any number of dimensions the absorption cross-section of massless minimally coupled scalars is exactly equal to the horizon area in the low energy limit [24]. However, it turns out that with fairly mild assumptions about the nature of worldvoulme interactions the pD-brane cross-sections fail to reproduce this lowest order answer except for 1-branes and 3-branes [25]. This result may be interpreted to imply that the agreement of absorption/emission cross-sections should be understood only in the sense of a coupling-energy expansion.

Maldacena [26] has derived non-renormalization properties of the low energy Yang-Mills field theory on the brane to explain agreement of the non-extremal entropy with the black hole answers. Similar arguments apply to some of the absorption/emission processes. We will, however, need properties of some higher order terms in the Born-Infeld action of branes, like terms involving products of four open string field derivatives. Such terms have a specific stringy origin and cannot be investigated in the Yang-Mills approximation. Furthermore, as mentioned above, we not only need to show that certain higher loop terms vanish but also need to show that the nonvanishing higher loop contributions are consistent with the black hole answer.

Consequently we look at properties of open string loop diagrams explicitly. Strictly speaking, this restricts us to a study of parallel branes - since in that case we know the microscopic theory accurately in terms of Chan-Paton factors. Our results would be rigourously valid for the extremal 3-brane. For the five and four dimensional black holes enough is not known about the bound states involved to enable a reliable loop calculation. However we do expect some consequences of our results for these black holes as well - this is based on the success of the effective long string picture.

4. Two open string processses

Consider the process of emission of a massless closed string state from two open strings annahilating on the worldvolume of N parallel p-D branes. One example is the emission of scalars in the five dimensional black hole which arise from components of the ten dimensional graviton h_{IJ} along the T^4 direction orthogonal to the 1D brane. The relevant interaction term in the tree level effective lagrangian density is given by

$$\sqrt{2}\kappa h_{IJ}(x^I, x^\alpha, x^i)\partial_\alpha X^I \partial^\alpha X^J \qquad (8)$$

Here I, J denote directions on the T^4 orthogonal to the 1D brane, α, β denote either time or the 1D brane direction and i, j denote directions transverse to $T^4 \times S^1$ on which the branes are wrapped. Another example is the absorption of scalars arising from the longitudinal components of the ten dimensional gravition by a collection of 3-branes. The interaction term in the low energy effective action is then given by

$$h_{\alpha\beta}\mathrm{Tr}[F^\gamma_\alpha F_{\gamma\beta} - \frac{1}{2}F^2 + \partial_\gamma X^i \partial^\gamma X^i + \cdots] \qquad (9)$$

where α, β, γ denote worldvolume indices and i, j transverse indices and F is the gauge field strength. The elipses denote the fermionic terms.

For emission into S-waves the closed string fields h_{IJ} in (8) and $h_{\alpha\beta}$ in (9) are independent of the transverse coordinates. The lowest energy contribution comes from $h_{IJ}, h_{\alpha\beta}$ which are dependent only on the worldvolume coordinates. These terms can be in fact read off from the effective action *in the absence of closed string fields*. This follows from the principle of equivalence. For example the flat space term for the 1D-5D system is

$$\delta_{IJ}\partial X^I \partial X^J \tag{10}$$

In the presence of a nontrivial metric δ_{IJ} has to be replaced by a tensor in the space transverse to the 1D brane but longitudinal to the 5D brane. Thus one may have

$$\delta_{IJ} \to (G_{IJ} + R_{IJ} + \cdots) \tag{11}$$

where G_{IJ} is the metric and R_{IJ} is the Ricci etc. In (12) the only term which does not involve derivatives of h_{IJ} is the metric G_{IJ} itself. All the other terms are therefore suppressed at low energies. This is however the term which gives rise to (8). Terms responsible for emission of higher partial waves cannot be read off from the flat space term in this fashion, since in this case one has to consider dependence on the transverse coordinates [18,20].

We want to compute the one (open string) loop correction to these terms. Thus we have to compute an annulus diagram with two open string massless vertices on the boundary and a massless closed string vertex in the interior of the annulus. The coordinate fields transverse to the brane satisfy Dirichlet conditions on the boundaries while the longitudinal ones satisfy Neumann conditions. As we have just argued for the lowest energy contribution it is sufficient to evaluate these terms *without* the closed string insertion. Since we have an *oriented* open string theory on the brane we have both a planar contribution $\mathcal{A}_p$ as well as a non planar contribution $\mathcal{A}_{np}$ which are given by

$$\begin{aligned} \mathcal{A}_p &= K_{(2,0)}\mathrm{Tr}[V(k)\Delta V(-k)\Delta] \\ \mathcal{A}_{np} &= K_{(1,1)}\mathrm{Tr}[V(k)\Omega\Delta V(-k)\Omega\Delta] \end{aligned} \tag{12}$$

where $K_{(2,0)}$ and $K_{(1,1)}$ are Chan Paton factors, Δ is the open string propagator and Ω is the twist operator. The net contribution is

$$\mathcal{A} = \mathcal{A}_p + \mathcal{A}_{np} \tag{13}$$

However both the terms in (13) vanish individually due to supersymmetry of the background. This may be seen for example in the Green-Schwarz formalism in the light cone gauge. To apply this formalism to D-branes one has to perform a double Wick rotation as explained in [27].The boundary conditions for the Green-Schwarz fermions S^a are obtained by requiring that half of the supersymmetries are preserved and become

$$S^a = M^a_b S^b \tag{14}$$

where M is a matrix satisfying the conditions [28, 29]

$$\begin{aligned} M^T M &= I \\ M^T\gamma^I M &= -\gamma^I \qquad I : \text{Dirichlet} \\ M^T\gamma^\alpha M &= \gamma^\alpha \qquad \alpha : \text{Neumann} \end{aligned} \tag{15}$$

The vertex operators for the open string massless states with polarizations in the Dirichlet and Neumann directions are given by (respectively)

$$\begin{aligned} V_D &= \zeta_I(k)(\partial_\sigma X^I - S_+\gamma^{I\alpha}S_+k_\alpha)e^{ikX} \\ V_N &= \zeta_\beta(k)(\partial_\tau X^\beta - S_+\gamma^{\beta\alpha}S_+k_\alpha)e^{ikX} \end{aligned} \tag{16}$$

The trace involves a trace of the zero modes of the fermions S^a_0 and we require at least eight fermionic fields to yield a nonzero trace. However the annulus diagram with only two open string insertions contains at most four fermionic fields and therefore they vanish. This means that there is no $O(g^2\omega^2)$ contribution to the amplitude.

The first nonzero contribution comes when the closed string insertion is taken into account since each closed string vertex contains four fermion fields (two from the left sector and two from the right sector). So with two open strings and one closed string there are the required eight fermionic fields. Clearly, to evaluate this term we may ignore the terms in the vertex operators which contain the bosonic fields. Furthermore in the lowest energy contribution we may ignore the e^{ikX} parts of the vertex operators since they will result in factors of the momentum. Thus any nonzero contribution to the amplitude from such a term would be of order $O(g^2\omega^4)$. We will return to the implication of such a nonzero term later.

The above result may be directly applied to the absorption by extremal 3-branes with zero momentum. It is clear from (9) that the tree level amplitude is of order $O(g\omega^2)$. Converting this to a transition rate and summing over initial states leads to a cross-section for S-waves

$$\sigma_{3D}^{l=0,tree} \sim (gN)^2\omega^3 \tag{17}$$

and as shown in [18] the coefficient is in exact agreement with the semiclassical cross-section. If there was a nonzero contribution to the amplitude of order $O(g^2\omega^2)$ one would have a $(gN)^4\omega^3$ term in the cross-section in contradiction with the classical cross-section. Thus our result "explains" why the perturbative D-brane calculation yields the low energy classical result.

5. Three open string processes

Processes involving three open strings are relevant for the absorption of $l = 1$ modes by the extremal 3-brane through terms like

$$(\partial_i\phi)\mathrm{Tr}[X^i F^2] \tag{18}$$

Using arguments similar to that in the previous section it may be easily seen that the lowest order one loop term could be of the order $O(g^{5/2}\omega^4)$. The resulting cross-section may be seen to be suppressed in energy compared to the tree level contribution.

6. Four open string processes

Processes involving four open strings display for the first time nontrivial effects of the sum over planar and non-planar diagrams. These processes are relevant for emission of "fixed scalars" in the five dimensional black hole [16,17]. In the effective long string model, where the 1D brane has to be considered multiply wound [30] along the lines of [31,32], expansion of the Born-Infeld action in the static gauge with flat worldsheet yields an action

$$\begin{aligned} S \quad \sim \quad & \int dx^0 dx^5 e^{-\phi}[1 + \frac{1}{2}e^{\phi/2}G^E_{IJ}\partial_+X^I\partial_-X^J \\ & -\frac{1}{8}e^{\phi}G^E_{IJ}G^E_{KL}\partial_+X^I\partial_+X^J\partial_-X^K\partial_-X^L \\ & +\cdots] \end{aligned} \tag{19}$$

where G^E denotes the Einstein frame metric. An example of a fixed scalar is the size of the T^4, $G_{IJ} = e^{2\nu}\delta_{IJ}$. If we require that the five dimensional dilaton $\phi_5 = \phi - 2\nu$ is not emitted, then it is seen that the lowest order term in (19) is quartic in X^I. When no h_{55} or h_{5I} are emitted and $Q_1 = Q_5$ the tree level cross-section agrees with the classical answer [16]. For $Q \neq Q_5$ one cannot consistently set h_{55} to zero and there is no agreement [33], while in the presence of h_{5I} agreement can be obtained by modifying the effeactive action [34]. We will restrict our attention to the case of $h_{55} = 0$ and $Q_1 = Q_5$. In this case the lowest energy interaction may be read off from the effective action in the absence of any closed string field along the lines of the previous section.

The one loop contribution is from an annulus diagram with four open string vertices at the boundaries. Now we have the required number of fermion zero modes to give a nonzero answer since each open string vertex operator contains two fermions and there are eight fermions in all. Thus the lowest order nonzero contribution comes from the term in which we can ignore the fermions in the open string propagators, ignore the bosonic parts of the vertex operators and ignore the e^{ikX} parts as explained above. Furthermore we may replace the fermionic fields by their zero modes so that there are no oscillators. If $\mathcal{A}_{m,n}$ denotes the amplitude with m vertices on one boundary and n vertices on the other boundary, the various contributions are

$$\begin{aligned} \mathcal{A}_{(4,0)} &= 2K_{(4,0)}\mathrm{Tr}[V(k_1)\Delta V(k_2)\Delta V(k_3) \\ & \quad \Delta V(k_4)\Delta] \\ \mathcal{A}_{(3,1)} &= 8K_{(3,1)}\mathrm{Tr}[V(k_1)\Omega\Delta V(k_2) \\ & \quad \Omega\Delta V(k_3)\Delta V(k_4)\Delta] \\ \mathcal{A}_{(2,2)} &= 6K_{(2,2)}\mathrm{Tr}[V(k_1)\Omega\Delta V(k_2) \\ & \quad \Delta V(k_3)\Omega\Delta V(k_4)\Delta] \end{aligned} \tag{20}$$

The effect of a twist operator is to change signs of nonzero oscillators in the following way

$$\alpha_n \to (-1)^n\alpha_n \qquad S_n \to (-1)^n S_n \tag{21}$$

However to this order there are no oscillators ! Thus the traces above are all the same - the only difference being in the Chan Paton factors.

A remarkable cancellation happens for single branes. Now the gauge group is $U(1)$ and

$$K_{(4,0)} = -K_{(3,1)} = K_{(2,2)} = e^4 \tag{22}$$

so that the sum of the three contributions in (20) vanishes. This cancellation is similar to the cancellation between different topologies for the one loop contribution to F^4 term in the Type I superstring [35] and is possibly related to similar non-renormalizations required in M(atrix) theory [36]. To the extent we can trust the single long effective string model, this may be taken to be an explanation of why the fixed scalar cross-section is correctly reproduced. The cases for nonvanishing h_{55} or h_{5I} cannot be treated in this manner since these interaction terms cannot be read off from the flat space effective action using principle of equivalence.

7. How big are the loop effects ?

As we have discussed the D-brane-black hole correspondence requires more than vanishing of certain terms in the open string loop corrections to absorption or emission processes : the form of the classical cross-section gives the precise form for these loop corrections for the case where there is only a single length scale in the problem. In this section we investigate this issue for the case of absorption by 3-branes where the only length scale is given by $l \sim (gN)^{1/4}\sqrt{\alpha}'$.

Since the tree level result is $\sim (gN)^2\omega^3$ the expansion of the classical cross-section following from (7) then shows that the one loop correction should be of the form

$$\sigma_{3D-1loop} \sim (gN)^4\omega^{11} \tag{23}$$

In Section 4. we found that, on the basis of zero mode counting, there is a possibility that there is a one-loop amplitude involving two open string states of the form $g^2\omega^4$. This would lead to a cross-section

$$\begin{aligned} \sigma &\sim \frac{1}{\omega}\int \prod_{i=1}^{2} \frac{d^3p^{(i)}}{|p^{(i)}|}\delta^4(p^{(1)}+p^{(2)}-k)(g^4\omega^8) \\ &\sim g^4\omega^7 \end{aligned} \tag{24}$$

where $(p^{(1)}, p^{(2)})$ denote the momenta of the open strings which are in the Neumann directions and k denotes the momentum of the closed string massless state which may be in any direction. If such a term was indeed present we would have a direct contradiction of the correspondence of 3-branes with seven dimensional extremal black holes. To examine this we need to calculate the coefficient of the term.

We will take the polarizations of the open string states η_1^μ, η_2^μ to be in any of the directions (the longitudinal polarizations coreespond to worldvolume gauge fields while the transverse polarizations are the scalars). The polarization of the massless closed string state $\epsilon_{\alpha\beta}$ is purely longitudinal and traceless - these are the minimally coupled scalars in the seven dimensional theory whose cross-sections have been computed in [18,19]. The absorbed scalar is also neutral which means that the spatial components of k^μ are entirely in the transverse directions.

As discussed above, for the lowest energy contribution, the vertex operators may be simplified. The open string operators may be replaced by

$$(S_0\gamma^{\mu\alpha}S_0)p_\alpha^{(m)}\eta_\mu^{(m)} \tag{25}$$

where the index $m = 1, 2$ labels the open string. The indices μ, ν may run over all the ten directions, while the indices α, β runs over the four worldvolume directions. The closed string vertex is similarly replaced by

$$(S_0\gamma^{\alpha\mu}S_0)(S_0\gamma^{\beta\nu}S_0)\epsilon_{\alpha\beta}k_\mu k_\nu \tag{26}$$

The final answer for the planar diagram is proportional to the well known trace [37]

$$t^{ijklmnpq} = \mathrm{Tr}[R_0^{ij}R_0^{kl}R_0^{mn}R_0^{pq}] \tag{27}$$

where

$$R_0^{ij} = S_0\gamma^{ij}S_0 \tag{28}$$

The trace is evaluated e.g. in [37]. Using the restrictions on the polarizations and momenta components described above it may be seen that there could be terms like

$$(p_\alpha^{(1)}\epsilon_{\alpha\beta}p_\beta^{(2)})(k\cdot\eta^{(1)})(k\cdot\eta^{(2)}) \tag{29}$$

which would lead, in position space to a term like

$$(\partial_\mu\partial_\nu h_{\alpha\beta})(\partial_\alpha X^\mu\partial_\beta X^\mu) \tag{30}$$

in the effective action. This would contribute to $l = 0$ and $l = 2$ partial wave absorptions. A detailed calculation shows that such terms cancel between two groups of terms [38] ! This result may be also seen from the covariantized form of the expression for the one loop effective action for four open string states given in [35].

The nonplanar diagram with two open string states involve only the $U(1)$ piece. This cancels the planar diagram by the mechanism described for diagrams with four open string states.

The lowest order contrbution in fact comes when the factors e^{ikX} are taken into account. This gives rise to an amplitude of the form $g^2\omega^6$. The corresponding cross-section is then of the form $(gN)^4\omega^{11}$ - *exactly what the black hole ordered* !

8. Conclusions

The black hole-D brane correspondence requires that the open string perturabtion theory of D-branes has a rather specific structure with higher powers of couplings coming with specific powers of the energy. We have found evidence for this structure by looking at one loop processes explicitly. The results presented above are, however, at best indicative of some deeper structure of the theory.

A deficiency of our method is that strictly speaking we can investigate only parallel p-branes. This is because of a lack of detailed understanding of the microscopic model for bound states of branes which describe five or four dimensional black holes in terms of string diagrams. So far in these cases the effective string model has been reasonably successful, though there are troublesome exceptions like fixed scalars with $h_{55} \neq 0$ [33] (for other discussions of the validilty of the effective string model see [39,40]) Furthermore higher angular momentum emission seems to require an effective string tension T_{eff} whose origin is not understood [41,42]. Recent work on the moduli space of these models in the string theory [43] as well as in the M(atrix) theory [44–46] may be useful in this regard.

In fact, black holes may teach us important features of string theory. Central to this is the issue of non-renormalization of certain quantities which make the string theoretic models of black hole work. An important example is the non-supersymmetric but extremal holes [47,48] where the agreeemnt of the entropy may be explained by mass renormalization of these states which are generically non-vanishing but vanish in the limit of large mass [49].

We need a systematic understanding of the loop effects discussed in this talk. After all it are these open string loop effects which give rise to space-time structure and a good understanding is certainly required to understand the physics of the horizon and issues of information loss in terms of string physics.

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REFERENCES

1. L. Susskind, hep-th/9309145; J. Russo and L. Susskind, Nucl. Phys. B 437 (1995) 611.
2. M. Duff and J. Rahmfeld, Phys. Lett. B345 (1995) 441, hep-th/9406105.
3. A. Sen, Nucl. Phys. B440 (1995) 421 , hep-th/9411187 and Mod. Phys. Lett. A10 (1995) 2081.
4. A. Strominger and C. Vafa, *Phys. Lett.* B379 (1996) 99, hep-th/9601029.
5. J. Maldacena and A. Strominger, Phys. Rev. Lett. 77 (1996) 428, hep-th/9603060; C. Johnson, R. Khuri and R. Myers, Phys. Lett. B378 (1996) 78, hep-th/9603061.
6. J. Breckenridge, R. Myers, A. Peet and C. Vafa, hep-th/9602065; J. Breckenridge, D. Lowe, R. Myers, A. Peet, A. Strominger and C. Vafa, Phys. Lett. B381 (1996) 423

7. C.G. Callan and J. Maldacena, Nucl. Phys. B475 (1996) 645, hep-th/9602043.
8. G. Horowitz and A. Strominger, Phys. Rev. Lett. 77 (1996) 2368, hep-th/9602051.
9. A. Dhar, G. Mandal and S.R. Wadia, Phys. Lett. B388 (1996) 51, hep-th/9605234.
10. S.R. Das and S.D. Mathur, Nucl. Phys. B478 (1996) 561, hep-th/9606185; Nucl. Phys. B482 (1996) 153, hep-th/9607149.
11. S. Gubser and I. Klebanov, Nucl. Phys. B482 (1996) 173, hep-th/9608108.
12. J. Maldacena and A. Strominger, Phys. Rev. D55 (1997) 861, hep-th/9609026.
13. S. Gubser and I. Klebanov, Phys. Rev. Lett. 77 (1996) 4491, hep-th/9609076.
14. S.P. de Alwis and K.Sato, Phys Rev. D55 (1997) 6581, hep-th/9611189.
15. S. Ferrara, R. Kallosh and A. Strominger, Phys. Rev. D52 (1995) 5412, hep-th/9508072.; S. Ferrara and R. Kallosh, hep-th/9602136; hep-th/9603090; G. Gibbons, R. Kallosh and B. Kol, hep-th/9607108; B. Kol and A. Rajaraman, hep-th/9608126.
16. C.Callan, S. Gubser, I. Klebanov and A. Tseytlin, hep-th/9610172.
17. I. Klebanov and M. Krasnitz, Phys. Rev. D55 (1997) 3250, hep-th/9612051.
18. I. Klebanov, Nucl. Phys. B496 (1997) 231, hep-th/9702076.
19. S. Gubser, I. Klebanov and A. Tseytlin, Nucl. Phys. B499 (1997) 217, hep-th/9703040
20. S. Gubser, I. Klebanov and A. Peet, Phys. Rev. D54 (1996) 3915, hep-th/9602135; I. Klebanov and A. Tseytlin, Nucl. Phys. B475 (1996) 165, hep-th/9604089.
21. G. Horowitz and J. Polchinski, Phys. Rev. D55 (1997) 6189,hep-th/9612146; hep-th/9707170.
22. S.R. Das, Phys. Rev. D56 (1997) 3582, hep-th/9703146.
23. I. Klebanov, hep-th/9709160.
24. S.R. Das, G. Gibbons and S.D. Mathur, Phys. Rev. Lett. 78 (1997) 417, hep-th/9609052.
25. S.R. Das, hep-th/9705165 -*to appear in Phys. Lett. B*
26. J. Maldacena, Phys. Rev. D55 (1997) 7645, hep-th/9611125.
27. M. Green and M. Gutperle, Phys. Lett. B377 (1996) 28, hep-th/9602077; Nucl. Phys. B476 (1996) 9604091; K. Hamada, hep-th/9612234.
28. S. Gubser, A. Hashimoto, I. Klabanov and J. Maldacena, Nucl. Phys. B472 (1996) 231, hep-th/9601057.
29. M. Garousi and R. Myers, Nucl. Phys. B475 (1996) 193, hep-th/9603194.
30. J. Maldacena and L. Susskind, Nucl. Phys B475 (1996) 679, hep-th/9604042.
31. S.R. Das and S.D. Mathur, Phys. Lett. B375 (1996) 103, hep-th/9601152.
32. S.D. Mathur, hep-th/9609053.
33. I. Klebanov and M. Krasnitz, Phys. Rev. D56 (1997) 2173, hep-th/9703216.
34. I. Klebanov, A. Rajaraman and A. Tseytlin, hep-th/9704112.
35. A. Tseytlin, Phys. Lett. B367 (1996) 84; Nucl. Phys. B467 (1996) 383.
36. T. Banks, W. Fischler, S. Shenker and L. Susskind, hep-th/9610043.
37. M. Green, J. Schwarz and E. Witten, *Superstring Theory, Vol. 2* (Cambridge University Press, 1986).
38. S.R. Das - *unpublished*
39. I. Klebanov and S.D. Mathur, hep-th/9701187.
40. F. Dowker, D. Kastor and J. Traschen, hep-th/9702109; D. kastor and J. Traschen, hep-th/9707157.
41. S.D. Mathur, hep-th/9704156.
42. S. Gubser, hep-th/9704195.
43. S.F. Hassan and S.R. Wadia, Phys. Lett. B402 (1997) 43, hep-th/9703163.
44. R. Dijkgraff, E. Verlinde and H. Verlinde, hep-th/9704018.
45. M. Li and E. Martinec, hep-th/9703211; hep-th/9704134.
46. E. Halyo, hep-th/9705107.
47. M. Duff and J. Rahmfeld, Nucl. Phys. B481 (1996) 332, hep-th/9605085; M. Duff, J. Liu and J. Rahmfeld, Nucl. Phys. B494 (1997) 161.
48. A. Dabholkar, Phys. Lett. B402 (1997) 53, hep-th/9702050.
49. A. Dabholkar, G. Mandal and P. Ramadevi, hep-th/9705239.