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A note on D-brane–anti-D-brane interactions in plane wave backgrounds

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Abstract: We study aspects of the interaction between a D-brane and an anti-D-brane in the maximally supersymmetric plane wave background of type IIB superstring theory, which is equipped with a mass parameter μ. An early such study in flat spacetime ($\mu = 0$) served to sharpen intuition about D-brane interactions, showing in particular the key role of the "stringy halo" that surrounds a D-brane. The halo marks the edge of the region within which tachyon condensation occurs, opening a gateway to new non-trivial vacua of the theory. It seems pertinent to study the fate of the halo for non-zero μ. We focus on the simplest cases of a lorentzian brane with $p = 1$ and an euclidean brane with $p = -1$, the D-instanton. For the lorentzian brane, we observe that the halo is unaffected by the presence of non-zero μ. This most likely extends to other (lorentzian) p. For the euclidean brane, we find that the halo is affected by non-zero μ. As this is related to subtleties in defining the exchange amplitude between euclidean branes in the open string sector, we expect this to extend to all euclidean branes in this background.

Keywords: D-branes, Superstring Vacua, Tachyon Condensation, Penrose limit and pp-wave background.

Contents

1. Introduction

A D-brane and its "anti-particle", an anti-D-brane, upon approaching each other, will annihilate. The generic product of this annihilation process is expected to be a state of closed strings, which carry no net R-R charge. This expectation is supported by field theory intuition and knowledge of which objects are the carriers of the available conserved charges in perturbative string theory. From experience with field theory one expects to be able to see the beginnings of the process of annihilation *via* the opening up of new decay channels at coincidence. These can be seen by studying the amplitude for exchange of quanta between the two branes, which gives a potential. At small separations, the behaviour of the interaction potential can signal new physics. Basically, a divergence in the amplitude as the objects are brought together can signal the opening up of a new channel (or new channels) not included in the computation of the amplitude away from the divergent regime.

In field theory, for a separation X of the two objects, the divergence follows simply from the fact that the amplitude for exchange is controlled by the position space propagator $\Delta(X)$ which (for more than two transverse directions) is divergent at $X = 0$. This is where the new channels can open up, which can include the processes for complete annihilation into a new sector, if permitted by the symmetries of the theory.

For D-branes in superstring theory, such a divergence does indeed show up, but there is an important new feature [1]. The divergence occurs when the D-branes are finitely separated, by an amount set by $X_H^2 = 2\pi^2\alpha'$, where α' is the characteristic length scale set by a fundamental string's tension. This is interpreted as the fact that in addition to the many special features of D-branes, they have a "stringy halo" originating in the fact that the bulk of the open strings which (by definition) end on them can reach out in the transverse directions, forming a region of potential activity of size set by X_H. This halo means that the D-branes can interact with each other before zero separation, as there is an enhancement of the physics of interaction by new light states formed by the entanglement of the halos, and the crossover into the annihilation channel begins before the branes are coincident.

Recall that the amplitude of exchange can be thought of using two equivalent pictures: Either as tree level exchange of closed string quanta between the branes, or (after a modular transformation) as the one-loop vacuum diagram for open strings stretched between the two D-branes. In the open string description, at separation X_H, the lightest open string becomes massless, and for any closer separation it becomes tachyonic, signalling that the entire vacuum configuration is unstable and wishes to roll to another vacuum. It is this tachyon which produces the divergence in the amplitude, converting a decaying exponential into a growing one, spoiling the convergence of the amplitude in the infra-red (IR) region.

The D-branes annihilate *via* conversion to closed strings in the generic situation, but the tachyon picture can be exploited in a beautiful way to produce more structure [2]–[5]. For the $G = \mathrm{U}(N) \times \mathrm{U}(N)$ gauge theory on the $(p+1)$-dimensional world-volume on N Dp-branes and N anti-Dp-branes, the tachyon field, transforming as the $(\mathbf{N}, \bar{\mathbf{N}})$, can be put into a configuration endowed with non-trivial topological charge, and the tachyon potential need not yield a runaway to a sector containing only closed strings. Having such topological vacuum solutions in the tachyon sector allows for the possibility of a stable remnant — interpreted as a D-brane of lower dimension — of the annihilation process after the debris that is the closed string products has cleared. It turns out that the spectrum of hypermultiplets in the $\mathrm{U}(N) \times \mathrm{U}(N)$ world-volume theory supplies a set of variables which is isomorphic to those needed to perform a K-theoretic analysis of the topology of G-vector bundles over the world-volume, and so the classification of all D-branes which can appear on a spacetime is apparently elegantly and economically by using the results of the appropriate K-theory of the spacetime which the Dp-branes and anti-Dp-branes fill [6, 7, 8]. The case of $p = 9$ for Minkowski spacetime yields the entire classification of D-branes in the most familiar symmetric vacuum of type IIB superstring theory.

This is all well understood for the case of flat ten dimensional spacetime. So when one encounters another background which enjoys the same maximal supersymmetry as flat spacetime — a plane wave with R-R flux [9]:

$$
\begin{aligned}
ds^2 &= 2dx^+dx^- - \mu^2 x^2 (dx^+)^2 + \sum_{i=1}^{4} dx^i dx^i + \sum_{i=5}^{8} dx^i dx^i \,, \\
F_{+1234} &= F_{+5678} = 2\mu \,, \qquad x^2 = \sum_{i=1}^{8} x^i x^i \,, \qquad x^{\pm} = \frac{1}{\sqrt{2}}(x^9 \pm x^0) \,,
\end{aligned}
\tag{1.1}
$$

which also yields an exactly solvable string model [10] (in light-cone gauge defined by relating worldsheet time τ to x^+ *via* $x^+ = 2\pi\alpha' p^+ \tau$, where p^+ is the $+$ component of spacetime momentum):

$$
\mathcal{L} = \frac{1}{4\pi\alpha'}(\partial_+ x^i \partial_- x^i - M^2 x^2) + \frac{i}{2\pi\alpha'}(S^a \partial_+ S^a + \tilde{S}^a \partial_- \tilde{S}^a - 2M S^a \Pi_{ab} \tilde{S}^a) \,, \tag{1.2}
$$

with a mass parameter $M = 2\pi\alpha' p^+ \mu$ — it is inevitable that questions about the key lessons which were learned about D-branes will spring to mind.[1] Is the picture of D-branes

[1] There has been a number of papers studying D-branes in plane wave and pp-wave backgrounds. Some of them are refs. [11]–[38].

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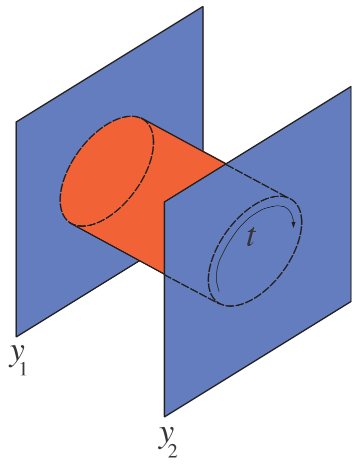

Figure 1: Cylinder diagram for computing the amplitude of interaction between two branes. The parameter t is open string propagation time, and is the modulus of the cylinder.

as Dirichlet open string boundary conditions as powerful in this context as it has been in flat spacetime? In particular, do the dynamics hidden within a halo's breadth of the branes bear any similarity to the flat spacetime case? Are all D-branes classified by K-theory, now of the new background?

In this paper we note that the properties of the halo — the fact that it exists, and also its location and size — are unaffected by non-zero μ for all branes that have a lorentzian definition, i.e., are at a definite position in space, but not time. So this particular (and important) property of D-branes in this non-trivial R-R background is very much like that in flat space. This bodes well for an attempt to classify such D-branes in this background using tachyon condensation and K-theory. Howver, for branes with a euclidean definition, such as the $p = -1$ brane, we find that the halo — or at least its analogue in this context — is deformed by non-zero μ.

2. The interaction

It is convenient [13, 18] to label D-branes in the plane wave background given in equation (1.1) as (r, s), if they are euclidean, where r denotes the spatial extent in directions $i = 1, 2, 3, 4$ and s denotes the spatial extent in directions $i = 5, 6, 7, 8$. A Dp-brane would then have $r + s = p + 1$. If the D-branes are lorentzian, then their worldvolume extends in the x^+ and x^- direction, and the notation is $(+, -, r, s)$. In that case, a Dp-brane has $r + s = p - 1$.

The string theory diagram of interest is a cylinder, representing either the tree level exchange of closed string quanta between two D-branes, or the one-loop vacuum process involving the circulation of open strings with ends on either D-brane. See figure 1.

We will focus on the results for the simplest branes in the euclidean and lorentzian classes. These are the D(-1)-branes (or $(0, 0)$-branes), and the D1-branes (or $(+, -, 0, 0)$-branes), discussed in ref. [18]. The former requires the time direction, in which the branes are also pointlike, to be euclidean.

The results are reasonably simple for these cases, compared to other (r,s) with $r\neq s\neq 0$, and it would be interesting to explore those other cases in detail. We expect that the key observations made in this paper for these $r=0=s$ cases will be quite generic, although there may be additional features to be deduced from studying other cases in detail.

2.1 The amplitude and potential

We consider a Dp-brane and its antiparticle for $p=\pm1$. If $p=-1$, it is an instanton, (a $(0,0)$-brane) and we consider it to be pointlike in euclidean time. If $p=+1$ it is a string, (a $(+,-,0,0)$-brane) and the theory is lorentzian.

So we place a Dp-brane at position y_1^i in the x^i directions ($i=1,\dots,8$), and a $\overline{\mathrm{D}p}$-brane (anti-brane) at position y_2^i, with a separation $X^\pm$ in the $x^\pm$ directions if $p=-1$. The cylinder amplitude A is [18]:

$$A=\int_0^\infty \frac{dt}{2t}t^{-\left(\frac{p+1}{2}\right)}e^{it\frac{X^+X^-}{2\pi\alpha'}}\hat{h}_0(t;y_1,y_2)\frac{\hat{g}_4^{(m)}(t)^4}{f_1^{(m)}(t)^8}\,. \tag{2.1}$$

where we have performed a double Wick rotation: $\tau\to it$, $x^+\to ix^+$. For $p=+1$, the factor $\exp(itX^+X^-/2\pi\alpha')$ is not present. For higher p, (which we will not be considering here) there are no additional powers of t in the integrand. These are normally due to integration over continuous zero modes in the flat spacetime case. The plane wave background has no such modes for the directions x^i, (the zero modes are instead themselves harmonic oscillators [39, 40, 41, 42]) and so no such t^{-1} factors beyond those appearing here are present. See below equation (2.5) for some further discussion of how to read this expression.

In the above, we have the functions:

$$\begin{aligned}
f_1^{(m)}(t) &= q^{-\Delta_m}(1-q^m)^{\frac12}\prod_{n=1}^\infty(1-q^{\omega_n})\,,\\
\hat{h}_0(t;y_1,y_2) &= \exp\Big(-\frac{mt}{2\alpha'\sinh(\pi m)}[\cosh(\pi m)(y_1^2+y_2^2)-2y_1\cdot y_2]\Big)\,,\\
\hat{g}_4^{(m)}(t) &= q^{-\hat{\Delta}_{(m)}}\prod_{l\in\mathcal{P}_+}\left(1-q^{|\omega_l|}\right)^{\frac12}\prod_{l\in\mathcal{P}_-}\left(1-q^{|\omega_l|}\right)^{\frac12}\,,\\
\Delta_m &= -\frac{1}{(2\pi)^2}\sum_{p=1}^\infty\int_0^\infty ds\; e^{-p^2s-\frac{\pi^2m^2}{s}}\,,\\
\hat{\Delta}_m &= -\frac{1}{(2\pi)^2}\sum_{p=1}^\infty(-1)^p\sum_{r=0}^\infty c_r^p m\frac{\partial^r}{(\partial m^2)^r}\frac{1}{m}\int_0^\infty ds\left(\frac{-s}{\pi^2}\right)^r e^{-p^2s-\frac{\pi^2m^2}{s}}\,,
\end{aligned} \tag{2.2}$$

and the parameter q and the deformed harmonic oscillator frequencies are defined as:

$$q=e^{-2\pi t}\,,\qquad \omega_n=\mathrm{sign}(n)\sqrt{n^2+m^2}\,. \tag{2.3}$$

Note that Δ_m and $\hat{\Delta}_m$ are zero-point energies which arise naturally in the closed and open string sectors, respectively. The coefficients c_r^p in $\hat{\Delta}_m$ are the coefficients of a specific Taylor expansion:

$$\left(\frac{x+1}{x-1}\right)^p + \left(\frac{x-1}{x+1}\right)^p = \sum_{r=0}^{\infty} c_r^p x^{2r} . \tag{2.4}$$

The sets $\mathcal{P}_-$ and $\mathcal{P}_+$ are given as solutions of the equations

$$l \in \mathcal{P}_- : \quad \frac{l-im}{l+im} + e^{2\pi il} = 0 , \qquad l \in \mathcal{P}+ : \quad \frac{l+im}{l-im} + e^{2\pi il} = 0 . \tag{2.5}$$

The details of the derivation of these amplitudes can be found in ref. [18]. We will not need them all here, and refer the reader there for more information. Some comments are in order however. For the case $p = 1$, the computation was done directly in terms of the open string channel, with open string light cone gauge $x^+ = 2\pi\alpha' p^+ \tau$, so we have

$$t = \frac{X^+}{2\pi\alpha' p^+} , \qquad m = 2\pi\alpha' p^+ \mu . \tag{2.6}$$

For the case $p = -1$, however, things are more subtle. A Dirichlet condition is needed in the time direction, but this is incompatible with the standard light-cone gauge choice. The amplitude is defined by appealing to the open-closed string duality instead. The amplitude is defined in the closed string sector by propagating for a distance X^+ between two boundary states. The propagation time is

$$\tilde{t} = \frac{X^+}{2\pi\alpha' p^+} , \tag{2.7}$$

since light-cone gauge in the closed string sector is $x^+ = 2\pi\alpha' p^+ \tau$, with mass parameter

$$M = 2\pi\alpha' p^+ \mu . \tag{2.8}$$

Open-closed string duality is then invoked to define the amplitude given in equation (2.1), where modular transformation gives the expression above, with

$$t = \frac{1}{\tilde{t}} = \frac{2\pi\alpha' p^+}{X^+} , \quad \text{and} \quad m = \mu X^+ = M t^{-1} . \tag{2.9}$$

This will be very important later.

2.2 Divergences, tachyons, and the halo

What is important for our discussion is the structure of the full amplitude for the cylinder diagram, given above in equation (2.1) as an integral over the modulus t. It can be written as:

$$A = \int_0^{\infty} \frac{dt}{2t} t^{-\left(\frac{p+1}{2}\right)} \exp\Big\{ -2\pi t Z(m, y_1, y_2) \Big\} G(t) , \tag{2.10}$$

where the the exponent $Z(m, y_1, y_2))$ is defined as (delete the X^+X^- term to get the D1-brane result):

$$Z(m, y_1, y_2) = \frac{m\pi}{4\pi^2\alpha' \sinh(m\pi)} \left[\cosh(m\pi)(y_1^2 + y_2^2) - 2y_1 \cdot y_2\right] - 4(\hat{\Delta}_m - 2\Delta_m) - i\frac{X^+X^-}{4\pi^2\alpha'} . \tag{2.11}$$

and the function $G(t)$ is defined as:

$$G(t) = \frac{\prod_{l\in\mathcal{P}_+}(1-q^{|\omega_l|})^2 \prod_{l\in\mathcal{P}_-}(1-q^{|\omega_l|})^2}{(1-q^m)^4 \prod_{n=1}^{\infty}(1-q^{\omega_n})^8}$$
$$= \prod_{n=1}^{\infty}(1-q^{\omega_n})^{-8} \prod_{l\in\mathcal{P}_+, l>0}(1-q^{\omega_l})^4 \prod_{l\in\mathcal{P}_-, l>0}(1-q^{\omega_l})^4 . \tag{2.12}$$

For our discussion, the only important fact about the function $G(t)$ is that its behaviour at large and small t is such that generically, the amplitude is convergent. That A is finite as $t \to 0$ follows from the fact that small t is the closed string IR limit, where this amplitude should reproduce simple low energy field theory results for massless exchange at tree level. The $t \to \infty$ limit is also well behaved generically, since this is the open string IR limit, which is fine — away from special circumstances which will not show up in the oscillator contributions since their energies are higher than the lowest lying states. In fact, it is clear that $G(t) \to 1$ as $t \to \infty$, and so whether A is finite as $t \to \infty$ depends on the sign of the exponent Z, which controls those lowest lying states.

The divergence for negative Z is related to the lowest lying states becoming tachyonic at this point, as is most familiar in the RNS formulation in the flat spacetime background. Then the worldsheet hamiltonian is given as $H = L_0 = \alpha' p^2 + N + a_{R(NS)}$, where the constant $a_{R(NS)}$ is the zero point energy and N is the total number operator. The z.p.e. is $a_R = 0$ in the Ramond sector, and $a_{NS} = -1/2$ in NS sector.

For strings stretched between two D-branes, we have $p^m = x^m/2\pi\alpha'$ for transverse (to the branes) directions x^m. So, splitting transverse (labelled m) and parallel (labelled i) directions we can write

$$L_0 = \alpha' p^i p_i + N + \frac{z^2}{4\pi^2\alpha'} + a_{R(NS)} . \tag{2.13}$$

This gives a mass spectrum

$$M^2 = -p^i p_i = \frac{1}{\alpha'}\left(N + a_{R(NS)} + \frac{z^2}{4\pi^2\alpha'}\right) . \tag{2.14}$$

The NS ground state ($N = 0$, $a_{NS} = -\frac{1}{2}$) has mass squared

$$M_0^2 = \frac{1}{2\alpha'}\left(\frac{z^2}{2\pi^2\alpha'} - 1\right) . \tag{2.15}$$

This is a tachyon if $z^2 < 2\pi^2\alpha'$.

In the usual case this ground state is eliminated by the GSO projection $P = \frac{1+(-1)^F}{2}$ in superstring theory. When we consider a brane-anti-brane system, we are effectively reversing the GSO projection in the partition function, giving $P = \frac{1-(-1)^F}{2}$, since anti-branes come with a minus sign. This means that the NS ground state ($N = 0$) will now survive, and the possible tachyon above is present in the spectrum. So for $z^2 < 2\pi^2\alpha'$ there is a tachyon, and so there is a 1-1 correspondence between the tachyon's appearance and divergence of the integral. (For the case when all of the directions are transverse, as is the case for D-instantons, the tachyon interpretation follows from continuation and T-duality.)

Let us write everything in terms z^i, the separation between the branes in the eight directions x^i, defined by $y_2^i = y_1^i + z^i$. The expression for Z then becomes

$$Z(m, y_1, z) = \frac{1}{4\pi^2\alpha'}\frac{m\pi}{\tanh(m\pi)}\left[(z+a)^2 - i\frac{\tanh(m\pi)}{m\pi}X^+X^- - b^2\right], \tag{2.16}$$

where we have defined:

$$a = \frac{\cosh(m\pi)-1}{\cosh(m\pi)}y_1\,, \qquad b = \tanh(m\pi)\sqrt{y_*^2 - y_1^2}\,, \qquad y_*^2 = \frac{16\pi^2\alpha'(\hat{\Delta}_m - 2\Delta_m)}{m\pi\tanh(m\pi)}\,. \tag{2.17}$$

For the lorentzian $p = 1$ case, these parameters simplify further in the $t \to \infty$ limit of interest. Since for fixed X^+ the large t region corresponds to small p^+ (this follows from equation (2.6), or on general grounds from the operator definition of the amplitude), we see that $m \to 0$ in all of these expressions, and so we obtain:

$$Z \longrightarrow \frac{1}{4\pi^2\alpha'}(z^2 - 2\pi^2\alpha')\,, \tag{2.18}$$

This is in fact the same expression one would obtain from the equivalent flat space computation, which simply has $m = 0$ throughout, and so we recover the well known [1] divergence at separation given by $X_H^2 = 2\pi^2\alpha'$. In fact, the result ought to be present for all lorentzian branes, as the relevant amplitude can be defined directly in the open string light cone gauge. Intuitively, we are looking for a result in the open string IR limit $t \to \infty$, which (from equation (2.6)) corresponds to $p^+ \to 0$. But the parameter upon which any new physics can depend is $m = 2\pi\alpha' p^+\mu$, which vanishes in the limit. So there is no new physics.

For the euclidean $p = -1$ case, the situation is very different. Now, for a given separation X^+, the $t \to \infty$ limit corresponds (due to equation (2.9)) to $p^+ \to \infty$ (this is the *closed string* momentum) and so things get quite reversed. In fact, the natural mass parameter seen by the open string physics is $m = \mu X^+$. In fact, there is quite a complicated dependence on m, as is evident from the equation (2.16). Looking (without loss of generality, since the spacetime is homogeneous[2]) at the case where we put one brane at the origin in the transverse directions, and so $y_1^i = 0$ and $z^i = y_2^i$, then the vanishing of Z can be written:

$$z^2 - i\frac{\tanh(\pi\mu X^+)}{\pi\mu X^+}X^+X^- = 2\pi^2\alpha'\mathcal{D}(\mu X^+)\frac{\tanh(\pi\mu X^+)}{\pi\mu X^+}\,, \tag{2.19}$$

where

$$\mathcal{D}(\mu X^+) = 8\left(\hat{\Delta}_m - 2\Delta_m\right)\,. \tag{2.20}$$

Recall that $\hat{\Delta}_m$ and Δ_m tend to $1/12$ and $-1/48$, respectively, when $m = \mu X^+$ tends to zero. The quantity $\mathcal{D}(\mu X^+)$ decreases from unity and asymptotes to zero as μX^+ increases. Of course, when μ (and hence m) vanishes, this gives the expected result:

$$z^2 - iX^+X^- = 2\pi^2\alpha' \equiv X_H^2\,. \tag{2.21}$$

[2]We thank Simon Ross for reminding us of the significance of this for D-instantons.

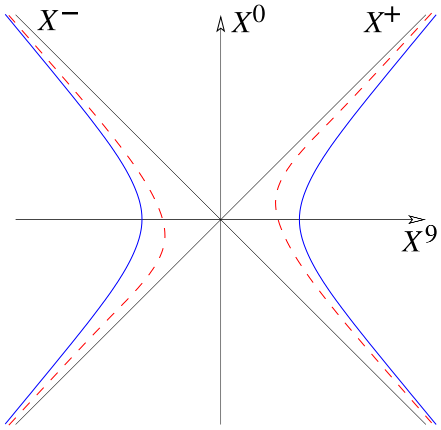

Figure 2: The hyperbola (solid curve) represents the edge of the "halo" for D-instantons in flat space, $\mu = 0$. For the $\mu \neq 0$ case, it is deformed to the dashed curve. The field theory result is the pair of lines $X^+X^- = 0$.

Note here that the unusual factor of $-i$ in this expression is as a result of the Wick rotation, which results in the (complexified) metric

$$ds^2 = -2idx^+dx^- + \mu^2 x^2 (dx^+)^2 + \sum_{i=1}^{8} dx^i dx^i \tag{2.22}$$

For non-zero μ it is hard to interpret the result cleanly, but there is certainly a non-trivial dependence of the location of the "halo" on μ, in contrast to the lorentzian case.

As a simple special case, one can place the branes at the same transverse position, and hence $z^i = 0$. Then we have the equation:

$$-iX^+X^- = 16\pi^2\alpha' \left(\hat{\Delta}_m - 2\Delta_m\right). \tag{2.23}$$

For orientation, let us consider the flat space case $\mu = 0$. We can continue to a more familiar lorentzian picture by choosing $X^- \to iX^-$. This gives a hyperbola in the plane, with equation

$$X^+X^- = 2\pi^2\alpha' \equiv X_H^2 . \tag{2.24}$$

Contrast this to the case of field theory, where the right hand side would be zero, giving us the light-cone. This is as expected for point like behaviour. The flat space string theory result gives us a hyperbola. This is the manifestation of the halo which broadens out the available region of contact by widening the light-cone into a sort of "light-funnel". For the $\mu \neq 0$ case, the hyperbola is deformed, since X^- decreases more rapidly with increasing X^+ than before due to the behaviour of the function $\mathcal{D}(X^+)$ discussed below equation (2.20). See figure 2.

For the interpretation of the shape of the halo for non-zero μ once the transverse positions of the branes are different from each other, more work is needed. This is because the metric is no longer flat, and furthermore, one has to take seriously the matter of the euclidean continuation of the metric implied in the computation of the amplitude. The choices made mean that the metric is no longer real (see equation (2.22)), and this presents difficulties of interpretation which must be explored further.

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3. Discussion

We have found that the structure of the halo for lorentzian branes in the plane wave background is independent of μ, giving the same physics as for D-branes in flat space. This is because the mass parameter induced by non-zero μ in the effective world-volume theory vanishes in the open string IR limit, the regime where the halo is to be found. We observed that this is not the case for the D-instanton (and presumably all euclidean branes), since their being pointlike in the $X^{\pm}$ directions requires the relevant amplitudes to be defined by starting with the closed string light cone gauge and then arriving at the open string physics by duality. The resulting open string physics sees a mass parameter which does not vanish in the IR limit, and hence the physics of the halo is not the same as in flat space. The significance of this non-trivial μ dependence of the structure of the halo of the D-instanton (and by extension, all euclidean branes defined by starting with the closed string amplitude) is not clear to us at present. However, it may have some significance, since D-instantons contribute to type IIB string theory processes non-perturbatively (see e.g., ref. [43]).

Note on earlier version of this manuscript. In an earlier version of this manuscript, we noted that there were μ-dependent effects for lorentzian branes as well. That was a mistake, and we apologise for any confusion caused. We misinterpreted the structure of the amplitudes in refs. [17, 18], and treated the effective mass parameter, m, in the open string sector as a fixed parameter in both the lorentzian and euclidean cases. This led us to that erroneous conclusion.

We note that Oren Bergman, Shinji Hirano and Dan Brace, and additionally Sakura Schafer-Nameki and Aninda Sinha informed us that they independently observed that our conclusions in the previous version were faulty and we thank them for comments and discussions.

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