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Type I D-particle and its Interactions

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Abstract

In a previous paper (hep-th/9808141) we showed that the type I string theory contains a stable non-BPS D-particle carrying SO(32) spinor charge. In this paper we formulate the rules for computing the spectrum and interaction of open strings with one or both ends lying on this D-particle.

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The discovery that D-branes carry Ramond-Ramond (RR) charge[1] − and hence represent stable BPS saturated particles in string theory − has dramatically changed our understanding of string theory. The inclusion of (wrapped) D-branes in the particle spectrum has provided key tests of many duality conjectures. Most of these tests have focussed on BPS particles in supersymmetric string theories; although there are some notable exceptions where non-BPS D-branes in non-supersymmetric string theories have been used to test duality conjectures involving non-supersymmetric string theories[2].

Although comparing the spectrum of BPS states in two theories provides us with a stringent test of duality between these two theories, quite often non-BPS states may also be useful in testing a duality conjecture. Thus for example if one of the theories contains, in its perturbative spectrum, a non-BPS state which is stable due to charge conservation (being the state of minimum mass carrying a given charge quantum number), then the dual theory must also contain a stable non-BPS state in its spectrum carrying the same quantum numbers. Several examples of this kind were discussed in [3]. One such example involved the duality relating the SO(32) heterotic and the type I string theories. The heterotic theory contains massive non-BPS states in the spinor representation of SO(32) which must be stable due to charge conservation. Thus the type I theory must also contain states in the spinor representation of SO(32). It was suggested in [3] that these states correspond to stable non-BPS D-particles in type I, although no explicit construction of such D-particles was known at that time.

This question was addressed in [4] from a slightly different angle. The starting point in [4] was a D-string anti- D-string pair of type I string theory. There is a tachyon field T living on the world volume of this system. The tachyon potential $V(T)$ is invariant under $T \to -T$ and hence has a doubly degenerate minimum at $\pm T_0$. Thus one can construct a kink configuration on the world volume of this system which interpolates between the tachyonic vacua $\pm T_0$. It was shown in [4] that this kink solution carries SO(32) spinor charge and has finite mass, and therefore represents the state that we have been searching for. Apparently this description is very different from that of a D-particle in type I. However it was found that this kink configuration admits a description in terms of a solvable boundary conformal field theory; and in this description it indeed corresponds to a D-particle for which ends of open strings satisfy Dirichlet boundary condition in all nine spatial directions.

Regarding the D-particle of type I as the tachyonic kink solution on the D-string anti-

D-string pair, and following the algorithm of ref.[4], we can in principle derive the complete set of rules for the spectrum and interaction of open strings whose one or both ends lie on the D-particle. However, in this paper we shall take a slightly different approach. Instead of trying to derive the interaction rules from [4] directly, we shall try to formulate a set of rules without any reference to the kink solution, take them to be the defining relations for a D-particle of type I and check the internal consistency of these rules. Of course in formulating the rules we shall be using the insight gained from [4].

As pointed out in [4], this D-particle can also be regarded as a configuration of type IIB string theory, although it has a tachyonic mode and hence is unstable. We shall find it more convenient to first describe the rules in the context of type IIB string theory, and specify the action of the world-sheet parity operator on the various open string modes. The open string spectrum in type I string theory is then obtained by projecting onto Ω-invariant states (vertex operators). The interaction rules involving open strings ending on these D-particles can be obtained by following the usual procedure of including unoriented world-sheet diagrams and extra combinatoric factors.

The D-particle of type IIB string theory, constructed this way, is in fact identical to the non-BPS D-particle discussed in [5].[2] This paper dealt with type IIB string theory in the presence of an orbifold plane with a D-particle stuck to the orbifold plane, and constructed the boundary state describing the D-particle[3]. The tachyon in the spectrum was removed by the orbifold projection instead of Ω projection as in the present case. However, since the boundary state describing a D-brane does not contain full information about the Chan Paton factors, it was not obvious that the D-particle described in this manner admits a consistent set of interaction rules. The present paper shows that it is indeed possible to construct such interaction rules.

We shall now specify the rules for computing the spectrum and interactions of open strings which end on the D-particle. Since besides the non-BPS D-particle, the type IIB string theory also contains BPS Dp-branes for odd p, we need to specify the rules for the spectrum and interaction of open strings in the 0-0, 0-p and p-0 sector.

0-0 sector: In this sector we have Dirichlet-Dirichlet (DD) boundary condition on all the space-like coordinates $X^1, \ldots X^9$ and Neumann-Neumann (NN) boundary condition

[2]This was in fact originally motivated by an attempt to construct a D-particle state in type I string theory[6].

[3]This state also has an alternate but equivalent description as a soliton on the D-string anti- D-string pair[7].

on the time coordinate X^0. There is no GSO projection on these open string states so that before the Ω projection the open string spectrum includes all the Fock space states. We associate a 2×2 Chan-Paton factor with each vertex operator in this sector, with the restriction that for states which are even under $(-1)^F$, the Chan Paton factor is the identity matrix, whereas for states which are odd under $(-1)^F$, the Chan Paton factor is given by

$$\sigma_2 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix} . \tag{1}$$

As usual, we take the Fock vacuum to be odd under $(-1)^F$.

In order to find the spectrum of 0-0 open strings in the type I theory we need to know the action of Ω on different open string states. Ω acts by transposition on the Chan Paton factors and has the usual action (see, for example [8]) on the various oscillators:

$$\alpha^\mu_r \to \pm e^{i\pi r} \alpha^\mu_r , \qquad \psi^\mu_r \to \pm e^{i\pi r} \psi^\mu_r , \tag{2}$$

where the upper (lower) sign holds for NN (DD) boundary condition, and α^μ_r and ψ^μ_r denote the oscillators associated with the bosonic and fermionic fields X^μ and ψ^μ respectively. Finally, the action of Ω on the Fock vacuum $|0\rangle$ is taken to be identity in computing its action on $(-1)^F$ odd states, and $-i$ for computing its action on the $(-1)^F$ even states. Thus for example, according to these rules, the Fock space states:

$$|0\rangle, \qquad \text{and} \qquad \psi^m_{-\frac{1}{2}}|0\rangle, \quad \text{for} \quad 1 \le m \le 9 , \tag{3}$$

are Ω even. However the state $|0\rangle$, being $(-1)^F$ odd, has a Chan Paton factor σ_2 associated with it. Since the action of Ω changes the sign of σ_2, the state corresponding to $|0\rangle$ is actually odd under Ω and hence is projected out in the type I theory. This rule can be simplified by ignoring the action of Ω on the Chan Paton factors, and taking the Ω charge of the Fock vacuum to be -1 and $-i$ for $(-1)^F$ odd and even states respectively.[4]

0-p and p-0 sectors: Let us take the Dp-brane to extend along $X^1, \ldots X^p$ direction. Then in the Neveu-Schwarz (NS) sector, the oscillators associated with the fermionic fields $\psi^1, \ldots \psi^p$ are integer moded and the oscillators associated with the fermionic fields $\psi^0, \psi^{p+1}, \ldots \psi^9$ are half integer moded. In the R sector the situation is reversed, $-$ the oscillators associated with the fields $\psi^1, \ldots \psi^p$ are half integer moded and the oscillators

[4]Although these rules seem completely ad hoc, they are the ones which emerge out of the analysis of ref.[4]. Also, as we shall show, they are internally consistent.

associated with the fields $\psi^0, \psi^{p+1}, \ldots \psi^9$ are integer moded. The allowed vertex operators (including the Chan Paton factors) in the 0-p sector are taken to be of the form:

$$\begin{pmatrix} 1 \\ 0 \end{pmatrix} \otimes V_s + i\epsilon \begin{pmatrix} 0 \\ 1 \end{pmatrix} \otimes V_{s'} , \tag{4}$$

where V_s is a $(-1)^F$ even vertex operator, and $V_{s'}$ is a $(-1)^F$ odd vertex operator, related to V_s through the relation:

$$|s'\rangle = U|s\rangle , \tag{5}$$

where,

$$U = (-1)^{F'} \psi_0^1 \ldots \psi_0^p , \tag{6}$$

in the NS sector, and

$$U = i(-1)^{F'} \psi_0^0 \psi_0^{p+1} \ldots \psi_0^9 , \tag{7}$$

in the R sector. Here $(-1)^{F'}$ changes the sign of all the fermionic oscillators except the zero modes, and $|s\rangle$ and $|s'\rangle$ are the Fock space states corresponding to the vertex operators V_s and $V_{s'}$. ϵ is a phase factor which is taken to be 1 if $U^2 = 1$ and i if $U^2 = -1$. From the structure of (4) we see that the physical states in the 0-p sector are in one to one correspondence to the set of $(-1)^F$ even vertex operators, although the actual construction of the vertex operator involves both $(-1)^F$ even and $(-1)^F$ odd operators.

In the p-0 sector the physical vertex operators are taken to be of the form:

$$\begin{pmatrix} 1 & 0 \end{pmatrix} \otimes V_s - i\epsilon^* \begin{pmatrix} 0 & 1 \end{pmatrix} \otimes V_{s'} , \tag{8}$$

where V_s is a $(-1)^F$ even vertex operator, and $V_{s'}$ and V_s are again related as in eq.(5). Since Ω relates the 0-p sector to the p-0 sector, in type I string theory we do not have independent contribution to the spectrum from the 0-p and the p-0 sectors.

Rules for computing amplitudes: Let us now specify the details of the combinatoric factors associated with a given amplitude. The most convenient way to specify these rules is to give a comparison with the amplitude associated with a D-particle of type IIA string theory. Compared to the various combinatoric factors in that theory, we have the following extra factors:

- For every hole with its boundary lying on the type IIB D-particle, we associate a factor of $\frac{1}{\sqrt{2}} \times Tr(I) = \sqrt{2}$.

- Around every hole whose boundary lies on a D-particle, we only impose periodic boundary condition on the fermion fields, instead of summing over periodic and anti-periodic boundary conditions.[5]

- For a world sheet boundary with insertion of open string vertex operators, different segments of the boundary can lie on different D-branes; hence we cannot specify separate rules for separate D-branes. The rule to be followed here is that we should sum over both periodic and anti-periodic boundary conditions on the fermions. We can regard the world sheet diagram with anti-periodic boundary condition as having a cut associated with the $(-1)^F$ operator ending on the boundary, If the cut ends on a segment lying on a p-brane for $p \neq 0$ we use the usual rules of taking trace over the product of Chan Paton factors. On the other hand *if the cut lands on a D-particle, then we insert an extra Chan Paton matrix σ_2 on the segment where the cut ends*. Of course, the final answer should not depend on how we choose the locations of the various cuts on the world-sheet.

This concludes our description of the rules for computing the spectrum and interactions of the open strings with one or both ends on the D-particle. Our next task is to check the consistency of this prescription. For this we first need to check the closure of the open string operator algebra. In particular, we need to make sure that in an arbitrary product of various operators, we do not generate any operator outside the set that we have included in the spectrum. We see this as follows:

1. First we shall check that if we take the product of a 0-0 and a 0-p sector vertex operator, we again get back a vertex operator of the form (4). We shall carry out the analysis in the case where the 0-0 sector vertex operator is in the NS sector. If this vertex operator is of the form $\mathcal{O} \otimes I$ with $\mathcal{O}$ being an $(-1)^F$ even operator, then the operator product has the form:

$$\begin{pmatrix} 1 \\ 0 \end{pmatrix} \otimes V_t + i\epsilon \begin{pmatrix} 0 \\ 1 \end{pmatrix} \otimes V_{t'} \,, \tag{9}$$

   where the states $|t\rangle$, $|t'\rangle$ are given by:

$$|t\rangle = \mathcal{O}|s\rangle, \qquad |t'\rangle = \mathcal{O}|s'\rangle \,. \tag{10}$$

[5]Specifying the boundary condition on the fermions requires a choice of coordinate system. We are using the coordinate system in which the hole corresponds to a circle around the origin of the complex plane.

Since p is odd, in the NS (R) sector the operator $(-1)^{F'}\psi_0^1\ldots\psi_0^p$ $(i(-1)^{F'}\psi_0^0\psi_0^{p+1}\ldots\psi_0^9)$ commutes with all oscillators and hence $\mathcal{O}$. This shows that $|t\rangle$ and $|t'\rangle$ are related in the same way as $|s\rangle$ and $|s'\rangle$ are related in eq.(5). Thus the resulting vertex operator has the form of (4).

If instead the 0-0 sector vertex operator has the form $\mathcal{O}'\otimes\sigma_2$, with $\mathcal{O}'$ a $(-1)^F$ odd operator, then the operator product has the form:

$$\begin{pmatrix}1\\0\end{pmatrix}\otimes V_u + i\epsilon\begin{pmatrix}0\\1\end{pmatrix}\otimes V_{u'}\,, \tag{11}$$

where the states $|u\rangle$, $|u'\rangle$ are given by:

$$|u\rangle = \epsilon\mathcal{O}'|s'\rangle, \qquad |u'\rangle = \epsilon^{-1}\mathcal{O}'|s\rangle\,. \tag{12}$$

Again it is easy to verify that $|u\rangle$ and $|u'\rangle$ are related in the same way as $|s\rangle$ and $|s'\rangle$ are in eq.(5). Thus the product has the structure of (4).

2. In the same way one can verify that the product of a p-0 and 0-0 sector vertex operator has the form of (8).

3. Now consider the product of a p-0 and 0-q vertex operator. From (4) and (8) we see that this product has the structure

$$V_sV_t + \epsilon_1\epsilon_2^* V_{s'}V_{t'}\,, \tag{13}$$

where ϵ_1 and ϵ_2 are the ϵ-factors associated with the p-0 and 0-q sectors respectively. Here V_s and V_t are $(-1)^F$ even, and $V_{s'}$ and $V_{t'}$ are $(-1)^F$ odd. Thus the product is a $(-1)^F$ even operator and represents an allowed vertex operator in the p-q sector.

4. Finally we consider the product of a 0-p and a p-0 vertex operator. For simplicity we shall carry out the analysis for the case where both vertex operators are in the NS sector or both are in the R sector, but the result is valid even when one is in the NS sector and the other is in the R sector. From (4) and (8) we see that this has the form:

$$\begin{aligned}&\begin{pmatrix}1\\0\end{pmatrix}(\,1\quad 0\,)\otimes V_s(x)V_t(y) + \begin{pmatrix}0\\1\end{pmatrix}(\,0\quad 1\,)\otimes V_{s'}(x)V_{t'}(y)\\ &-i\epsilon^*\begin{pmatrix}1\\0\end{pmatrix}(\,0\quad 1\,)\otimes V_s(x)V_{t'}(y) + i\epsilon\begin{pmatrix}0\\1\end{pmatrix}(\,1\quad 0\,)\otimes V_{s'}(x)V_t(y)\,.\end{aligned} \tag{14}$$

From the relation (5) between $|s\rangle$, $|s'\rangle$, and using the definition of ϵ, we can easily verify that

$$\langle s|\mathcal{O}|t\rangle = \langle s'|\mathcal{O}|t'\rangle, \qquad \langle s|\mathcal{O}|t'\rangle = \epsilon^2\langle s'|\mathcal{O}|t\rangle, \tag{15}$$

for all operators $\mathcal{O}$. This gives the relations:

$$V_s(x)V_t(y) = V_{s'}(x)V_{t'}(y), \qquad V_s(x)V_{t'}(y) = \epsilon^2 V_{s'}(x)V_t(y)\,. \tag{16}$$

Using eqs.(16) we can rewrite (14) as

$$I \otimes V_s(x)V_t(y) + \epsilon^* \sigma_2 \otimes V_s(x)V_{t'}(y)\,. \tag{17}$$

Both of these are allowed operators in the 0-0 sector. Thus we see that the product of vertex operators in the 0-p and p-0 sectors closes on the allowed vertex operators in the 0-0 sector.

This establishes the closure of the operator algebra in the open string sector. Next we need to verify the closure of the open-closed operator algebra. In particular, since the spectrum of open strings contains both $(-1)^F$ odd and $(-1)^F$ even states, one might worry if the interaction of such open strings could produce closed string states carrying odd world-sheet fermion number. We shall now show that this does not happen. For this let us consider a disk amplitude with open string vertex operators inserted at the boundary and a closed string vertex operator inserted in the interior of the disk. We shall only consider the cases of NSNS and RR sector closed string vertex operators. Using the closure of the open string operator algebra, we can reduce this to a diagram where only the 0-0 open strings are inserted at the boundary. In this sector, the $(-1)^F$ even states are accompanied by the Chan Paton factor I whereas the $(-1)^F$ odd states are accompanied by the Chan Paton factor σ_2. First let us take the closed string vertex operator to be in the NS-NS sector so that the fermions satisfy periodic boundary condition along the boundary of the disk and hence there is no extra Chan Paton factor besides those coming from the open string vertex operators. In this case an amplitude containing odd number of $(-1)^F$ odd open string states will have a net Chan factor factor σ_2 inserted at the boundary and will vanish upon taking the trace. Hence the only non-vanishing amplitudes are those with even number of $(-1)^F$ odd open string states. This, in turn guarantees that the closed string vertex operator inserted in the interior of the disk must have the correct GSO projection in order to get a non-zero amplitude. On the other hand if the closed string vertex operator is in the RR sector, then there will be an extra Chan Paton factor of σ_2 at the boundary, and hence now non-vanishing trace over the Chan Paton factor will require an odd number of $(-1)^F$ odd open string vertex operators. This would indicate that the closed string vertex operator must be $(-1)^F$ odd. However note that in the language of type IIA string theory (which we are using as the reference for comparison) the allowed RR-sector vertex operators in the type IIB theory are indeed odd under $(-1)^F$. Thus the interaction rules again guarantee that only those closed string states which satisfy the correct GSO projection condition of type IIB string theory are produced in the scattering of open strings.

The above analysis also explains the reason for the extra Chan Paton factor of σ_2 that needs to be put in if a cut lands on a segment of the boundary lying on the D-particle. Let us consider a 0-p vertex inserted on a boundary, with a cut passing through the segment lying on the p-brane. As the cut moves from the p-brane to the 0-brane, we naturally shift from the IIB to the IIA convention for defining the $(-1)^F$ quantum numbers of spin fields.[6] This means that the definition of $(-1)^F$ odd and $(-1)^F$ even spin fields get reversed. Since the operator ϵU defined in (6), (7) anti-commutes with $(-1)^F$, this shift in the definition of $(-1)^F$ is generated by multiplying the 0-p vertex operators by ϵU. On the other hand, from (4) we can easily verify that multiplying a 0-p vertex operator by ϵU is equivalent to multiplying it by σ_2. This shows that as the cut moves from the p-brane to the 0-brane, it is accompanied by an extra Chan Paton factor of σ_2.

Let us now turn to the consistency of the annulus diagram. In particular we need to show that the annulus diagram admits a double interpretation, $-$ as an open string loop diagram and as a closed string tree diagram. We shall only discuss the case where both boundaries of the annulus lie on the D-particle, but the consistency of other diagrams with one boundary on the 0-brane and the other boundary on a p-brane can also be checked. Since each boundary gives a contribution proportional to $Tr(I)/\sqrt{2} = \sqrt{2}$, we get a net factor of 2 compared to the annulus diagram of the type IIA D-particle. At the same time, we put only periodic boundary condition on the fermions along the non-trivial cycle, instead of summing over periodic and anti-periodic boundary conditions. Since the annulus diagram for the D-particle in type IIA string theory is known to be given by an

[6] This is related to the fact that the RR sector of a boundary state describing a D-particle can only be defined for the type IIA string theory, and not for type IIB string theory.

integral of

$$Tr\Big(\frac{1+(-1)^F}{2}q^{L_0}\Big)\,, \tag{18}$$

we see that the annulus diagram for the D-particle in IIB has the form of an integral over

$$Tr(q^{L_0})\,. \tag{19}$$

Here q is a modular parameter to be integrated over. The absence of the $(-1)^F$ term in (19) is a reflection of the absence of anti-periodic boundary condition on the fermions. The extra multiplicative factor of 2 is due to the contribution from the Chan Paton factors. (19) can be clearly interpreted as the partition function of the open string states without any GSO projection, and hence can be identified as the partition function of an open string with both ends on the type IIB D-particle.

We can also reinterprete the annulus diagram in the closed string channel. In the case of type IIA D-particle, the $\frac{1}{2}Tr(q^{L_0})$ term may be interpreted as the result of interaction between the D-particles due to NSNS sector closed string exchange, and the $\frac{1}{2}Tr((-1)^F q^{L_0})$ term may be interpreted as the result of interaction due to the exchange of RR sector closed string exchange. Comparing this with (19) we see that in the present case there is no contribution from the exchange of RR sector closed strings, whereas the contribution from the NSNS sector exchange doubles. Thus the D-particle of type IIB does not carry any RR charge, and is $\sqrt{2}$ times heavier than the D-particle of type IIA. This gives the type IIB D-particle mass to be

$$\frac{\sqrt{2}}{g}(\alpha')^{-\frac{1}{2}}\,, \tag{20}$$

where g is the string coupling constant and $(2\pi\alpha')^{-1}$ is the string tension.

This extra $\sqrt{2}$ factor in the mass can also be seen from the fact that the disk tadpole for the graviton has an extra factor of $\sqrt{2}$ for the D0-brane of IIB compared to that of IIA, due to the extra contribution of $\frac{1}{\sqrt{2}}Tr(I)$ from the Chan Paton factor. On the other hand, the disk tadpole for an RR sector state vanishes in this theory due to periodic bounday condition on the fermions along the boundary of the disk.

This result can be stated in a different way by constructing the boundary state for the D-particle of IIB. This has been given in [5]. Let us denote the boundary state describing the D-particle in type IIA by

$$\frac{1}{\sqrt{2}}(|NSNS\rangle + |RR\rangle)\,, \tag{21}$$

where $|NSNS\rangle$ and $|RR\rangle$ denote the contribution from the NSNS and RR sectors respectively. Then the boundary state describing the D-particle of type IIB is given by

$$|B\rangle = |NSNS\rangle \,. \tag{22}$$

By construction, the inner product of this boundary state with itself produces the partition function of the 0-0 string. But as a further consistency check, one can verify that the inner product of this boundary state with the boundary state of a p-brane in type IIB correctly reproduces the partition function of the 0-p string. Furthermore, in type I string theory we also need to take the inner product of this with the crosscap boundary state $|C\rangle$[9, 10]. This combined with $\langle B|B\rangle$ can be shown to give the correct Ω projected open string partition function in the 0-0 sector.

This finishes our analysis of the consistency checks for the interaction rules given at the beginning of this paper. We shall now use these rules to study some properties of the D-particle. First we note that even though the D-particle in type IIB is unstable due to the presence of the tachyonic ground state in the NS sector, this tachyon is projected out in type I string theory (as discussed below eq.(3)). Thus the D-particle is stable in type I string theory. In order to find its quantum numbers, we need to find the spectrum of massless fermionic open string states living on the D-particle and quantize them. Again the easiest way to proceed is by comparing the present situation to that of a D-particle of type IIA string theory. The Ramond sector ground state in the 0-0 sector gives rise to massless fermions as usual, but due to the absence of GSO projection, the number of such fermionic states living on the type IIB D-particle is double of that living on the type IIA D-particle. Embedding this in type I halves the number of fermion zero modes. Thus the number of fermion zero modes living on the type I D-particle is *equal* to the number of such zero modes living on the type IIA D-particle. Since the type IIA D-particle is 256-fold degenerate, we see that the type I D-particle also acquires a 256-fold degeneracy. This corresponds to a long multiplet of the N=1 supersymmetry in ten dimensions, exactly as expected of a non-BPS state.

But the type I D-particle has further degeneracy from the 0-9 sector strings. Since the type I vacuum contains 32 nine branes, we get 32 extra massless fermions from the Ramond sector ground state of the 32 0-9 strings. These zero modes form a 32 dimensional Clifford algebra, and their quantization gives rise to a state in the spinor representation of SO(32).[7] Thus we see that the type I D-particle describes a stable, non-BPS state in the spinor representation of SO(32). This is precisely the state we have been looking for.

The next question we shall address is the stability of the D-particle in the presence of various other branes. First let us consider the case of two such D-particles on top of each other. In this case the Ω projection is no longer sufficient to project out the tachyonic mode that originates from an open string whose two ends lie on the two different D-particles. Thus this system becomes unstable. This is not surprising, since a pair of such D-particles has SO(32) quantum number in the scalar conjugacy class. Thus such a system can roll down to the vacuum (possibly with other massless particles). The tachyonic mode simply represents the possibility of such a transition.

In general, when we bring the D-particle on top of a Dp-brane, the NS sector ground state of the open string with one end on the D0-brane and the other end on the Dp-brane corresponds to a particle with squared mass

$$m^2 = -\frac{1}{2}(1-\frac{p}{4})(\alpha')^{-\frac{1}{2}}\,, \tag{23}$$

Thus we see that the D0-D1 system has a tachyonic mode whereas the D0-D5 and D0-D9 system has no tachyonic mode. The tachyonic instability of the D0-D1 system again reflects the existence of a lower energy configuration carrying the same quantum number. To see this let us consider the case of type I string theory compactified on a circle, with the D-string wrapped on the circle. In this case the spinor charge of the D0-brane can be absorbed into the D-string by putting a Z_2 Wilson line along the circle at no cost of energy[11, 12]. The tachyonic instability of the D0-D1 system simply reflects the possibility of rolling down to this lower energy configuration.

Thus we see that the D-particle of type I described in this paper not only satisfies the mathematical conditions required for consistent open-closed interaction, but also satisfies all the physical properties expected of such a particle.

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[7] Actually we get states in spinor as well as conjugate spinor representation of the gauge group. However, there is a residual Z_2 gauge symmetry on the world-volume of the D-particle, as in [11], under which the 0-9 strings are odd, and requiring invariance under this gauge transformation projects out states in the conjugate spinor representation. I wish to thank O. Bergman and M. Gaberdiel for discussion on this point.

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