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Notes on D-branes of type IIB string on $AdS_5 \times S^5$

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Abstract

We promote a study of D-branes of type IIB string on the $AdS_5 \times S^5$ background. The possible D-branes preserving half of supersymmetries were classified up to and including the fourth order of fermionic variable θ in our previous work [hep-th/0310228]. In this work we show that our classification is still valid even at the full order of θ. This proof supplements our previous results and completes the classification of D-branes in the type IIB string theory on the $AdS_5 \times S^5$.

Keywords: D-branes; AdS-string; Penrose limit; pp-wave

1. Introduction

D-brane is an important key ingredient in studies of non-perturbative aspects of superstring theories [1]. A recent interest is to investigate D-branes on general curved backgrounds, motivated by recent developments in studies of pp-wave backgrounds. To begin with, the maximally supersymmetric type IIB pp-wave background was found [2]. Then the Green–Schwarz type IIB string theory on this pp-wave was shown to be exactly solvable in the light-cone gauge [3,4]. After that D-branes on the pp-wave were intensively investigated [5–10] since one can study directly them by solving classical equations of motion and quantizing the theory.

The covariant studies of D-branes in type IIB and IIA strings on pp-waves were discussed in [11] and [12], respectively, by using the method of Lambert and West [13]. Motivated by these developments, we have carried out the covariant analysis for D-branes of type IIB string on the $AdS_5 \times S^5$ background [14]. The allowed $1/2$ supersymmetric (SUSY) D-brane configurations have been classified. Our result is also consistent to that of brane probe analysis done in [7]. In addition, Penrose limit [15,16][1] of D-branes in the $AdS_5 \times S^5$ has been discussed and we have seen that the result after this limit agrees with the possible D-brane configurations in the type IIB pp-wave background. On the other hand, by combining the methods proposed in [18,19], the covariant method is

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[1] Penrose limit of superalgebra in the $AdS_5 \times S^5$ background was discussed in [17].

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also applicable to open supermembrane theory on the pp-wave [20,21] and $AdS_{4/7} \times S^{7/4}$ [22] backgrounds. These results are related via Penrose limit and are also consistent to the brane probe analysis in eleven dimensions [23].

In this Letter, we continue to study D-branes of type IIB string on the $AdS_5 \times S^5$ background [14]. The previous analysis at the fourth order of θ is extended to the full order. We show that the higher order terms with respect to θ do not affect the classification of $1/2$ SUSY D-branes at the fourth order under the conditions obtained in the fourth order analysis. This proof completes our classification of D-branes.

The organization of the present Letter is as follows. In Section 2, we introduce our previous result on the classification of D-branes in type IIB string theory on the $AdS_5 \times S^5$ background, based on the analysis up to and including the fourth order of θ. In Section 3, we show that the classification is still valid even at the full order of θ. Section 4 is devoted to a conclusion and discussions.

2. Classification of 1/2 SUSY D-branes

Here we will briefly review our classification result of $1/2$ SUSY D-branes of type IIB string theory on the $AdS_5 \times S^5$ background [14]. We work in the notation and convention used in [14].

The open-string world-sheet Σ has the one-dimensional boundary $\partial\Sigma$, and we can impose the Neumann and Dirichlet boundary conditions on $\partial\Sigma$. These conditions are represented by

$$\partial_\sigma X^{\overline{A}} \equiv \partial_\sigma X^M e_M^{\overline{A}} = 0 \quad \text{(Neumann condition)}, \tag{2.1}$$

$$\partial_\tau X^{\underline{A}} \equiv \partial_\tau X^M e_M^{\underline{A}} = 0 \quad \text{(Dirichlet condition)}, \tag{2.2}$$

where we have used the overline as $\overline{A}_i$ $(i = 0, \ldots, p)$ for the indices of Neumann coordinates and the underline as $\underline{A}_j$ $(j = p+1, \ldots, 9)$ for the indices of Dirichlet coordinates.

By using the projection operators $P^{\pm}$, boundary conditions are imposed on the fermionic variable θ as

$$P^{\pm}\theta = \theta, \qquad P^{\pm} = \frac{1}{2}(1 \pm M). \tag{2.3}$$

The gluing matrix M is described as follows:

$$M = \begin{cases} m \otimes i\sigma_2, & d = 2\,(\mathrm{mod}\,4), \quad p = -1, 3, 7, \\ m \otimes \rho, & d = 4\,(\mathrm{mod}\,4), \quad p = 1, 5, 9, \end{cases}$$

$$m = s\Gamma^{\underline{A}_1} \cdots \Gamma^{\underline{A}_d}, \qquad s = \begin{cases} 1, & \text{for } X^0\text{: Neumann}, \\ i, & \text{for } X^0\text{: Dirichlet}, \end{cases} \qquad \rho = \begin{cases} \sigma_1, & \text{when } \sigma = \sigma_3, \\ \sigma_3, & \text{when } \sigma = \sigma_1. \end{cases} \tag{2.4}$$

In [14], a classification of $1/2$ SUSY D-branes in the $AdS_5 \times S^5$ was given by considering the vanishing conditions of the κ-variation surface terms up to and including the fourth order in θ.

For the $d = 2\,(\mathrm{mod}\,4)$ case, the possible configurations of D-branes need to satisfy the following condition:

- The number of Dirichlet directions in the AdS_5 coordinates $(X^0, \ldots, X^4)$ is even, and the same condition is also satisfied for the S^5 coordinates $(X^5, \ldots, X^9)$.

For the $d = 4\,(\mathrm{mod}\,4)$ case, D-branes satisfying the following condition are allowed:

- The number of Dirichlet directions in the AdS_5 coordinates $(X^0, \ldots, X^4)$ is odd, and the same condition is also satisfied for the S^5 coordinates $(X^5, \ldots, X^9)$.

The D-branes in the $AdS_5 \times S^5$ are restricted with respect to the directions to which a brane world-volume can extend. All of possible D-brane configurations at the origin are summarized in Table 1. When we consider the D-branes sitting outside the origin, only a D-instanton is allowed.

Table 1
The possible $1/2$ supersymmetric D-branes in $AdS_5 \times S^5$ sitting at the origin

D-instanton	D-string	D3-brane	D5-brane	D7-brane	D9-brane
(0, 0)	(0, 2), (2, 0)	(1, 3), (3, 1)	(2, 4), (4, 2)	(3, 5), (5, 3)	absent

In the next section we will discuss that the higher order terms with respect to θ do not modify the above classification of D-branes sitting at and outside the origin.

3. Validity of the classification at full order of θ

Now let us show that our classification at the fourth order of θ is still valid at the full order of θ. We shall start from the covariant Wess–Zumino term [24]:

$$\mathcal{L}_{\mathrm{WZ}} = -2i \int_0^1 dt\, \hat{E}^A \bar{\theta} \Gamma_A \sigma \hat{E}, \tag{3.1}$$

where $\hat{E}^A(X,\theta) \equiv E^A(X,t\theta)$ and $\hat{E}^\alpha(X,\theta) \equiv E^\alpha(X,t\theta)$.

We notice that the surface term coming from the κ-variation is represented by[2]

$$E_\tau^A (\bar{\theta}\Gamma_A)_\alpha \sigma \delta_\kappa Z^{\hat{M}} E_{\hat{M}}^\alpha, \tag{3.2}$$

where E_τ^A denotes the τ-component of E_i^A. Here we should remark that the Nambu–Goto or Dirac–Born–Infeld part of the action does not produce any surface term under the κ-variation. Then, in order to check our classification of D-branes at the full order of θ, it is sufficient to show the key relations:

$$E_\tau^{\underline{A}} = \partial_\tau Z^{\hat{M}} E_{\hat{M}}^{\underline{A}} = 0, \tag{3.3}$$

$$P^\mp \delta_\kappa Z^{\hat{M}} E_{\hat{M}}^\alpha = 0 \quad \text{for } \theta = P^\pm \theta \tag{3.4}$$

under the conditions denoted in Section 2. When the relations (3.3) and (3.4) are proven, we can easily see that the surface term (3.2) should vanish as

$$E_\tau^{\bar{A}} \bar{\theta} P^\pm \Gamma_{\bar{A}} \sigma \delta_\kappa Z^{\hat{M}} E_{\hat{M}} = E_\tau^{\bar{A}} \bar{\theta} \Gamma_{\bar{A}} \sigma P^\mp \delta_\kappa Z^{\hat{M}} E_{\hat{M}} = 0. \tag{3.5}$$

Therefore, all we have to do in order to show the validity at the full order is to prove two relations (3.3) and (3.4). We will prove (3.3) and (3.4) below by using the $1/2$ SUSY conditions obtained at the fourth order analysis. Before going to the detail analysis, we should remark about the D-instanton case. This case has no Neumann coordinates, and so it is sufficient to see (3.3) only in order for the surface term to vanish.

3.1. Proof of (3.3)

Here let us show the relation (3.3). For this purpose, we consider the term $\partial_\tau X^M E_M^{\underline{A}}$ and rewrite it as

$$\partial_\tau X^M E_M^{\underline{A}} = \partial_\tau X^M \left[e_M^{\underline{A}} + i\bar{\theta}\Gamma^{\underline{A}} \left(\frac{\sinh(\mathcal{M}/2)}{\mathcal{M}/2} \right)^2 [D\theta]_M \right] = \partial_\tau X^M \left[i\bar{\theta}\Gamma^{\underline{A}} P^\mp \left(\frac{\sinh(\mathcal{M}/2)}{\mathcal{M}/2} \right)^2 P^\mp [D\theta]_M \right]$$

[2] We will omit the symbol "hat" of E^A and E^α because the shift $\theta \to t\theta$ does not affect our discussion.

$$= i\frac{\lambda}{2}\partial_\tau X^M e_M^{\overline{B}}\bar\theta\Gamma^{\underline{A}}P^{\mp}\left(\frac{\sinh(\mathcal{M}/2)}{\mathcal{M}/2}\right)^2 P^{\mp}\hat\Gamma_{\overline{B}}i\sigma_2\theta$$
$$+\frac{i}{4}\partial_\tau X^M e_M^{\overline{D}}\bar\theta\Gamma^{\underline{A}}P^{\mp}\left(\frac{\sinh(\mathcal{M}/2)}{\mathcal{M}/2}\right)^2 P^{\mp}\Gamma_{\overline{B}\underline{C}}\theta\omega_{\overline{D}}^{\overline{B}\underline{C}}, \tag{3.6}$$

where we have introduced the following notation:

$$[D\theta]_M \equiv \frac{\lambda}{2}e_M^B\hat\Gamma_B i\sigma_2\theta + \frac{1}{4}\omega_M^{AB}\Gamma_{AB}\theta.$$

In the second equality in (3.6), we have used the definition of the Dirichlet boundary condition

$$\partial_\tau X^M e_M^{\underline{A}} = 0, \tag{3.7}$$

and the identities

$$P^+\mathcal{M}^{2n}P^- = P^-\mathcal{M}^{2n}P^+ = 0, \tag{3.8}$$

which hold under the conditions obtained in the fourth order analysis.

The first term in the most right-hand side of (3.6) vanishes under the conditions: the even number of Dirichlet directions are contained in the case of $d = 2\,(\mathrm{mod}\,4)$, or the odd number of Dirichlet ones are in the $d = 4\,(\mathrm{mod}\,4)$ case. The second term vanishes at the origin because the spin connection vanishes at the origin. That is, the last line of (3.6) vanishes at the origin under the conditions obtained in the fourth order analysis. Therefore, we obtain

$$\partial_\tau X^M E_M^{\underline{A}} = 0. \tag{3.9}$$

In addition, we can easily show the relation

$$\partial_\tau\theta^{\bar\alpha}E_{\bar\alpha}^{\underline{A}} = 0, \tag{3.10}$$

under the fourth order conditions. Thus, we have shown that the relation (3.3) should be satisfied under the conditions found at the fourth order.

We should note that (3.6) is trivially zero for the D-instanton because it contains no Neumann directions. The condition (3.6) $= 0$ is satisfied at and outside the origin. Namely, we have seen that $1/2$ SUSY D-instanton has no modification from higher order terms of θ. And the consideration for $1/2$ SUSY D-branes sitting outside the origin has been completed.

Finally, we would like to comment on the physical interpretation of the condition (3.3). The condition (3.3) implies that the configurations of D-branes are static, because it represents that the momenta for Dirichlet directions are zero. It would be also helpful to consider a flat limit ($\lambda \to 0$) and to see

$$E_\tau^{\underline{A}} = \dot X^{\underline{A}} - i\bar\theta\Gamma^{\underline{A}}\dot\theta = 0. \tag{3.11}$$

This interpretation should be plausible because D-branes moving in the target space would be less supersymmetric rather than $1/2$ supersymmetric ones. It is well known that the κ-symmetry in open string theories governs the dynamics of D-branes [25]. This fact may be partially realized in our case as the condition (3.3).

3.2. *Proof of* (3.4)

We shall prove the relation (3.4) below. In order to show it, we need to prove the relation:

$$\delta_\kappa X^M e_M^{\underline{A}} = 0. \tag{3.12}$$

By definition of the κ-transformation, we have

$$\delta_\kappa E^A = \delta_\kappa X^M E_M^A + \delta_\kappa\theta^{\bar\alpha}E_{\bar\alpha}^A = 0. \tag{3.13}$$

From this equation, we obtain

$$\delta_\kappa X^M e_M^A = -i\bar\theta\Gamma^A\left(\frac{\sinh(\mathcal{M}/2)}{\mathcal{M}/2}\right)^2[D\theta]_N e_B^N e_M^B\delta_\kappa X^M + E_{\bar\alpha}^A\delta_\kappa\theta^{\bar\alpha} \equiv H^A{}_B\delta_\kappa X^M e_M^B + E_{\bar\alpha}^A\delta_\kappa\theta^{\bar\alpha}. \tag{3.14}$$

Here, in order to make the structure clear, let us introduce the following abbreviations:

$$\delta_\kappa X^M e_M^A \equiv \delta x^A, \qquad \delta_\kappa\theta E^A \equiv \delta\theta^A, \tag{3.15}$$

and then (3.14) is written as

$$\delta x^A = H^A{}_B\delta x^B + \delta\theta^A. \tag{3.16}$$

By using this equation recursively, we can derive the following expression:

$$\delta x^A = (H^{15} + \cdots + 1)^A{}_B\delta\theta^B. \tag{3.17}$$

Now let us evaluate each of terms in the r.h.s. of (3.17) and show that all of them are zero in the case that A is a Dirichlet direction. Noting that

$$\delta\theta^B = E_{\bar\alpha}^B\delta_\kappa\theta^{\bar\alpha} = i\bar\theta\Gamma^B\left(\frac{\sinh(\mathcal{M}/2)}{\mathcal{M}/2}\right)^2\delta_\kappa\theta\begin{cases}\neq 0, & B\text{: Neumann},\\ = 0, & B\text{: Dirichlet},\end{cases} \tag{3.18}$$

Eq. (3.12) becomes

$$\delta x^{\underline{A}} = (H^{15} + \cdots + 1)^{\underline{A}}{}_{\overline{B}}\delta\theta^{\overline{B}} = 0. \tag{3.19}$$

The zero-th order term with respect to H is obviously zero because $\delta^{\underline{A}}{}_{\overline{B}} = 0$. The first order term $H^{\underline{A}}{}_{\overline{B}}$ is written as

$$H^{\underline{A}}{}_{\overline{B}} = -i\bar\theta\Gamma^{\underline{A}}\left(\frac{\sinh(\mathcal{M}/2)}{\mathcal{M}/2}\right)^2[D\theta]_M e_{\overline{B}}^M = -i\bar\theta\Gamma^{\underline{A}}P^{\mp}\left(\frac{\sinh(\mathcal{M}/2)}{\mathcal{M}/2}\right)^2 P^{\mp}[D\theta]_M e_{\overline{B}}^M. \tag{3.20}$$

We can rewrite furthermore $P^{\mp}[D\theta]_M e_{\overline{B}}^M$ as

$$P^{\mp}[D\theta]_M e_{\overline{B}}^M = P^{\mp}\left(\frac{\lambda}{2}e_M^C\hat\Gamma_C i\sigma_2\theta + \frac{1}{4}\omega_M^{CD}\Gamma_{CD}\theta\right)e_{\overline{B}}^M = \frac{\lambda}{2}\hat\Gamma_{\overline{B}}i\sigma_2 P^{\mp}\theta + \frac{1}{4}\omega_{\overline{B}}^{\overline{C}\underline{D}}\Gamma_{\overline{C}\underline{D}}P^{\pm}\theta. \tag{3.21}$$

The first term in the extreme right-hand side vanishes for $P^{\pm}\theta = \theta$. Because the components of spin connection vanish at the origin, i.e., $\omega_{\overline{B}}^{\overline{C}\underline{D}} = 0$, the term $H^{\underline{A}}{}_{\overline{B}}$ should vanish at the origin. Using this fact, we find that

$$H^{\underline{A}}{}_{C_1}H^{C_1}{}_{C_2}\cdots H^{C_n}{}_{\overline{B}} = H^{\underline{A}}{}_{\underline{C}_1}H^{\underline{C}_1}{}_{\underline{C}_2}\cdots H^{\underline{C}_n}{}_{\overline{B}} = 0. \tag{3.22}$$

Thus, we have shown the useful identity (3.19), and so (3.12).

Now let us return to the proof of (3.4) and consider

$$P^{\mp}\delta_\kappa Z^{\hat M}E_{\hat M}^\alpha = P^{\mp}\big(\delta_\kappa X^M E_M^\alpha + \delta_\kappa\theta^{\bar\beta}E_{\bar\beta}^\alpha\big), \tag{3.23}$$

for the boundary conditions $\theta = P^{\pm}\theta$. First, we shall consider the first term in r.h.s. of (3.23), which can be rewritten as

$$\begin{aligned}P^{\mp}\frac{\sinh\mathcal{M}}{\mathcal{M}}[D\theta]_M\delta_\kappa X^M &= P^{\mp}\frac{\sinh\mathcal{M}}{\mathcal{M}}P^{\mp}[D\theta]_{\overline{A}}e_M^{\overline{A}}\delta_\kappa X^M\\ &= P^{\mp}\frac{\sinh\mathcal{M}}{\mathcal{M}}P^{\mp}\left(\frac{\lambda}{2}\hat\Gamma_{\overline{A}}i\sigma_2\theta + \frac{1}{4}\omega_{\overline{A}}^{BC}\Gamma_{BC}\theta\right)e_M^{\overline{A}}\delta_\kappa X^M.\end{aligned} \tag{3.24}$$

In the first line of (3.24), we have used the relations, (3.8) and (3.12). The first term in the last line of (3.24) always vanishes for the conditions $\theta = P^{\pm}\theta$. The second term is also zero at the origin by using $\theta = P^{\pm}\theta$. Hence we have seen that the first term in (3.23) vanishes. Furthermore, we can see that the second term in (3.23) also vanishes from the relation:

$$P^{\mp}\frac{\sinh\mathcal{M}}{\mathcal{M}}\delta_\kappa\theta = P^{\mp}\frac{\sinh\mathcal{M}}{\mathcal{M}}P^{\mp}\delta_\kappa\theta = 0, \tag{3.25}$$

because of $\theta = P^{\pm}\theta$. Thus, we have shown that the second key relation (3.4).

Finally, we would like to comment on the physical implication of the condition (3.4). When we consider the flat limit ($\lambda \to 0$), (3.4) is reduced to the $1/2$ SUSY condition:

$$P^{\mp}\delta_\kappa\theta = 0. \tag{3.26}$$

Hence (3.4) may be a generalization of projection condition to the $AdS_5 \times S^5$ case.

4. Conclusion and discussion

We have shown that higher order surface terms with respect to θ, which come from the κ-variation, do not affect the classification obtained in the fourth order analysis. This proof completes our previous classification at the fourth order. As a matter of course, our proof is obviously applicable to the full order analysis of D-branes in the type IIB string on the pp-wave background. But the validity of D-brane classification in the pp-wave should be obvious via the Penrose limit [15–17] of D-branes on the $AdS_5 \times S^5$. Hence, we may say that the fourth order analysis in the covariant formulation is sufficient to classify the possible configurations of D-branes.

In this work, we have considered D-branes of open F- and D-strings, and presented a simple prescription for the vanishing conditions of the κ-variation surface terms. Our scenario can be extended to D-branes of an open Dp-brane in an obvious way. It is also interesting to consider intersecting D-branes (for those on pp-waves, see [26]). We reserve these issues for the next publication. Also, it is not quite trivial in the case of open supermembrane on the pp-wave and $AdS_{4/7} \times S^{7/4}$ backgrounds because of the dimensionality. We will report on these cases in another place soon [27].

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