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Dual D-brane actions on $AdS_5 \times S^5$*

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Abstract: Utilizing coset superspace approach, dual actions of super D1-and D3-branes on $AdS_5 \times S^5$ are constructed by carrying out duality transformation of world-volume $U(1)$ gauge field. Resulting world-volume actions are shown to possess expected $SL(2, \mathbf{Z})$ properties. Crucial ingredient for deriving $SL(2, \mathbf{Z})$ transformation property of the D-brane actions is covariance of $SU(2, 2|4)$ coset superspace algebra under $SO(2)$ rotation between two ten-dimensional Type IIB Majorana-Weyl spinors.

Keywords: D-branes, Superstrings and Heterotic Strings, Superstring Vacua.

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Contents

1. Introduction

Recently, motivated by Maldacena's conjecture on AdS/CFT correspondence, closed forms of κ-symmetric superstring and D3-brane actions have been constructed using the coset superspace approach [1, 2]. The formal superspace expression of the superstring or D-brane actions on Type-IIB supergravity backgrounds have been constructed previously by various authors [3, 4, 5]. However, using the coset superspace approach, it has become possible to obtain an explicit form of the action in terms of the coordinate fields (x, θ). This development is quite exciting since, for example, it now opens up a possibility for solving large-N limit of $\mathcal{N} = 4$ super-Yang Mills theory exactly.

As is now well known, Type IIB string theory possesses $SL(2,\mathbf{Z})$-duality symmetry [6], which is believed to hold not only for flat spacetime but also for any classical background. Hence, it would be of some interest to check whether the D-brane actions on $AdS^5 \times S^5$ has indeed the right transformation property under the $SL(2,\mathbf{Z})$-duality. In particular, on the $AdS^5 \times S^5$ background, the D3-brane action should be self-dual under $SL(2,\mathbf{Z})$, and the actions for D1- and D5-branes should form a multiplet of $SL(2,\mathbf{Z})$.

In this paper, we prove explicitly the expected $SL(2,\mathbf{Z})$ transformation properties of the D1- and D3-brane actions on $AdS_5 \times S^5$, at least in the semi-classical limit. The formal expression of $AdS_5 \times S^5$ D3-brane action obtained in [1] is quite similar to the action in flat spacetime [7]. This then suggests that the proof of the $SL(2,\mathbf{Z})$ duality symmetry ought to be similar to that for the flat spacetime case [8]. An important ingredient in the proof of the $SL(2,\mathbf{Z})$ symmetry in the flat spacetime has been that the S-generator of the $SL(2,\mathbf{Z})$ corresponding to $\tau \to -\frac{1}{\tau}$ acts as a finite $SO(2)$ rotation in the space of the two fermionic coordinates [8, 9]. We show that the same continues to be true on $AdS_5 \times S^5$ background. In fact, the $\mathcal{N} = 2$ Type IIB supersymmetry algebra from which the Mauer-Cartan equations are derived is covariant under the $SO(2)$ rotation. Stated differently, if a set of the two supercharges, Q_1 and Q_2, is related to an another set, Q'_1 and Q'_2, by a $SO(2)$ rotation, then the $\mathcal{N}=2$ supersymmetry algebra turns out to retain the same form when expressed in either set of the supercharges.

In section 2, we recapitulate relevant part of the construction [1] of D3-brane action on $AdS_5\times S^5$. We then prove the self-duality of the action under $SL(2,\mathbf{Z})$ by performing the duality transformation on the world-volume fields as well as the background fields. The proof is in fact almost the same as that for the flat space case [8]. In section 3, we proceed to the D1-brane and construct the action on $AdS_5\times S^5$. We then show that the action is, as expected, $SL(2,\mathbf{Z})$ covariant.[1]

As our work is being finished, we have received related works [12]. In [12], gauge fixing of the κ-symmetry was indispensible for showing the $SL(2,\mathbf{Z})$ symmetry. Our proof is based only on the covariance of the $SU(2,2|4)$ superalgebra under the $SO(2)$-rotation and hence proves the $SL(2,\mathbf{Z})$ duality symmetry without any gauge fixing.

2. $SL(2,\mathbf{Z})$ invariance of the D3-brane on $AdS_5\times S^5$

The superstring or the D3-brane action can be constructed using the $SU(2,2|4)$ superalgebra. If we consider the coset superspace $\frac{SU(2,2|4)}{SO(1,4)\times SO(5)}$, the even part is $\frac{SO(2,4)}{SO(1,4)}\otimes\frac{SO(6)}{SO(5)}\simeq AdS^5\otimes S^5$. The corresponding algebra is described in $SO(1,4)\times SO(5)$ basis, as shown in [1]. The even generators are two pairs of translation and rotation, (P_a,J_{ab}) for AdS^5 and (P'_a,J'_{ab}) for S^5 and the odd generator are the two 10-d Majorana-Weyl spinors $Q_I^{\alpha\alpha'}$. Let γ_a and γ'_a be the 4×4 matrices generating the $SO(1,4)$ and $SO(5)$ Clifford algebra with the signature $\gamma^{(a}\gamma^{b)}=(-++++)$ and $\gamma^{(a'}\gamma^{b')}=(+++++)$. The 10-dimensional 32×32 Dirac matrices $\Gamma^{\hat{a}}$ of $SO(1,9)$ are represented as

$$\Gamma^a=\gamma^a\otimes I\otimes\tau_1,\quad \Gamma^{a'}=I\otimes\gamma^{a'}\otimes\tau_2\,, \tag{2.1}$$

where I is the 4×4 unit matrix and τ_i are the Pauli matrices. In [1], the diagonal Majorana condition is chosen

$$\bar{Q}_{\alpha\alpha' I}\equiv(Q_I^{\beta\beta'})^\dagger(\gamma^0)_\alpha^\beta\delta_{\alpha'}^{\beta'}=-Q_I^{\beta\beta'}C_{\beta\alpha}C_{\beta'\alpha'}\,, \tag{2.2}$$

where $C=(C_{\alpha\beta})$ and $C'=(C_{\alpha'\beta'})$ are the charge conjugation matrix of the $SO(4,1)$ and $SO(5)$ Clifford algebra. The part of the supersymmetry algebra containing the spinors is

$$\begin{aligned}[][Q_I,P_a]&=-\frac{i}{2}\varepsilon_{IJ}Q_J\gamma_a\,,\\ [Q_I,P'_a]&=\frac{1}{2}\varepsilon_{IJ}Q_J\gamma'_a\,,\end{aligned}$$

[1] The behavior under the $SL(2,\mathbf{Z})$ transformation in the general Type II background has been also discussed in [10, 11].

$$[Q_I, J_{ab}] = -\frac{1}{2} Q_I \gamma_{ab} \,,$$
$$[Q_I, J_{a'b'}] = -\frac{1}{2} Q_I \gamma_{a'b'} \,,$$

$$\{Q_{\alpha\alpha' I}, Q_{\beta\beta' J}\} = \delta_{IJ}(-2iC_{\alpha'\beta'}(C\gamma^a)_{\alpha\beta}P_a + 2C_{\alpha\beta}(C'\gamma^{a'})_{\alpha'\beta'}P_{a'}) + \\ +\varepsilon_{IJ}(C_{\alpha'\beta'}(C\gamma^{ab})_{\alpha\beta}J_{ab} - C_{\alpha\beta}(C'\gamma^{a'b'})_{\alpha'\beta'}J_{a'b'}) \,. \tag{2.3}$$

Note that the diagonal Majorana condition is preserved under the $SO(2)$ rotation between Q_1 and Q_2. And one can easily check that the $SU(2,2|4)$ superalgebra is covariant under the $SO(2)$ rotation, as δ_{IJ} and ε_{IJ} are invariant under $SO(2)$.

The left invariant Cartan 1-forms $L^A = dX^M L_M^A$, $X^M = (x, \theta)$ are given by

$$G^{-1}dG = L^A T_A = L_a P^a + L_{a'} P^{a'} + \frac{1}{2} L^{ab} J_{ab} + \frac{1}{2} L^{a'b'} J_{a'b'} + L^{\alpha\alpha' I} Q_{\alpha\alpha' I} \,, \tag{2.4}$$

where $G = G(x, \theta)$ is a coset representative in $SU(2,2|4)$, $L^a, L^{a'}$ are the 5-beins, $L^{\alpha\alpha' I}$ are two spinor 16-beins and $L^{ab}, L^{a'b'}$ are the Cartan connection.

The D3-brane action presented in [1] is[2]

$$S = -\int d^4\sigma \sqrt{-\det\left(G_{ij} + e^{-\frac{\phi}{2}} F_{ij} - b_{ij}\right)} + \\ +\int \left(C_4 + C_2 \wedge (e^{-\frac{\phi}{2}} F - b_2) + \frac{1}{2} C_0 F \wedge F\right) . \tag{2.5}$$

Here we have $G_{ij} = L_i^{\hat a} L_j^{\hat a} = \partial_i X^M L_M^{\hat a} \partial_j X^N L_N^{\hat a}$ where $\hat a = (a, a')$ run through the $SO(1,4)$ and $SO(5)$ indices and $L^{\hat a}(X(\sigma)) = d\sigma^i L_i^{\hat a}$. Also $b_2 \equiv \frac{1}{2} b_{ij} d\sigma^i d\sigma^j$ and $db_2 = i\bar L \tau_3 \hat L L$ with $\hat L_i = L_i^{\hat a} \Gamma^{\hat a}$. Finally C_4 and C_2 are determined by the condition

$$H \equiv d(C_4 + C_2 \wedge (e^{-\frac{\phi}{2}} F - b_2) = \frac{i}{6} \bar L \hat L^3 \tau_3 \tau_1 L + i\bar L \tau_1 (e^{-\frac{\phi}{2}} F - b_2) \hat L L + \\ +\frac{1}{30}(\varepsilon^{a_1 \ldots a_5} L^{a_1} \wedge \cdots \wedge L^{a_5} + \varepsilon^{a'_1 \ldots a'_5} L^{a'_1} \wedge \cdots \wedge L^{a'_5}) \,. \tag{2.6}$$

From this condition one obtains the following useful identity:

$$dC_4 - C_2 \wedge db_2 = \frac{i}{6} \bar L \tau_3 \tau_1 \hat L^3 L + \frac{1}{30}(\varepsilon^{a_1 \ldots a_5} L^{a_1} \wedge \cdots \wedge L^{a_5} + \varepsilon^{a'_1 \ldots a'_5} L^{a'_1} \wedge \cdots \wedge L^{a'_5}) \,. \tag{2.7}$$

In eq. (2.5), for later convenience, we have written down the dilaton dependence in the Einstein frame of the $AdS_5 \times S^5$ background. In fact, the string and Einstein frames are equivalent for D3-brane, since the dilaton takes a constant value on $AdS_5 \times S^5$. The C_0 term in eq. (2.5) is absent in [1], but it is a total derivative term (or boundary term) that can be added to the action without changing the classical equations of motion. The inclusion of the constant C_0 is consistent since the dilaton remains constant in the $AdS^5 \times S^5$ background. In fact, a constant shift of C_0 is a trivial classical symmetry of

[2] Our convention is that whenever an integral without $d^n\sigma$, it is an integral of a differential form.

the action. Inclusion of the C_0 term, however, is quite crucial for the proof since the $SO(2)$ duality symmetry is promoted to the full $SL(2,\mathbf{R})$ symmetry in the presence of both dilaton and C_0 field background.

Adding a Lagrange multiplier term $\int d^4\sigma \frac{1}{2}\tilde{H}^{ij}\left[F_{ij} - (\partial_i A_j - \partial_j A_i)\right]$ to the action eq. (2.5), the A_i equation of motion can be solved by $\tilde{H}^{ij} = \varepsilon^{ijkl}\partial_k B_l$. The duality transformation is essentially the same as in the flat case [8] and the resulting action is

$$S_D = -\int d^4\sigma \sqrt{-\det\left(G_{ij} + \frac{1}{\sqrt{1+e^{2\phi}C_0^2}}(e^{\frac{\phi}{2}}\tilde{F}_{ij} + C_{ij} + e^{\phi}C_0 b_{ij}\right)} + \int \Omega_D\,, \tag{2.8}$$

where

$$\begin{aligned}\Omega_D = {} & C_4 - b_2\wedge C_2 - \frac{1}{2}e^{\phi}C_0 b_2\wedge b_2 + b_2\wedge\left(e^{\frac{\phi}{2}}\tilde{F} + C_2 + e^{\phi}C_0 b_2\right) - \\ & -\frac{e^{\phi}C_0}{2(1+e^{2\phi}C_0^2)}\left(e^{\frac{\phi}{2}}\tilde{F} + C_2 + e^{\phi}C_0 b_2\right)\wedge\left(e^{\frac{\phi}{2}}\tilde{F} + C_2 + e^{\phi}C_0 b_2\right),\end{aligned} \tag{2.9}$$

where $\tilde{F} = dB$. If we define rotated $SO(2)$ Pauli matrices

$$\begin{aligned}\tau_1' &\equiv -\left(\tau_3 + e^{\phi}C_0\tau_1\right)/\sqrt{1+e^{2\phi}C_0^2}\,,\\ \tau_3' &\equiv +\left(\tau_1 - e^{\phi}C_0\tau_3\right)/\sqrt{1+e^{2\phi}C_0^2}\end{aligned} \tag{2.10}$$

and using eq. (2.7) and the identity $\tau_3\tau_1 = \tau_3'\tau_1'$, we were able to show that

$$\begin{aligned}d\Omega_D = {} & \frac{i}{6}\bar{L}\hat{L}^3\tau_3'\tau_1' L + i\bar{L}\tau_1'\left(\frac{e^{\frac{\phi}{2}}}{\sqrt{1+e^{2\phi}C_0^2}}\tilde{F} - b_2'\right)\hat{L}L + \\ & + \frac{1}{30}\left(\varepsilon^{a_1\ldots a_5}L^{a_1}\wedge\cdots\wedge L^{a_5} + \varepsilon^{a_1'\ldots a_5'}L^{a_1'}\wedge\cdots\wedge L^{a_5'}\right).\end{aligned} \tag{2.11}$$

Thus the dual action can be written as

$$\begin{aligned}S = {} & -\int d^4\sigma\sqrt{-\det\left(G_{ij} + \frac{e^{\frac{\phi}{2}}}{\sqrt{1+e^{2\phi}C_0^2}}\tilde{F}_{ij} - b_{ij}'\right)} + \\ & + \int\left\{C_4' + C_2'\wedge\left(\frac{e^{\frac{\phi}{2}}}{\sqrt{1+e^{2\phi}C_0^2}}\tilde{F} - b_2'\right) - \frac{e^{2\phi}C_0}{2(1+e^{2\phi}C_0^2)}\tilde{F}\wedge\tilde{F}\right\},\end{aligned} \tag{2.12}$$

where the tensor fields C_4', C_2' and b_2' are the same as C_4, C_2 and b_2 but with $SO(2)$ rotated Pauli matrices in the corresponding expression. The resulting action still has the κ-symmetry since the Mauer-Cartan equation and Fierz identity used in the derivation of the κ-symmetry are covariant under the $SO(2)$ rotation.

Also, we can check the transformation of the dilaton and the axion under the duality transformation. From the coefficient of $\tilde{F}$ in the Born–Infeld part of eq. (2.12), we obtain the transformation

$$e^{-\phi} \qquad \rightarrow \qquad +\frac{e^{\phi}}{1+e^{2\phi}C_0^2} = \frac{1}{e^{\phi} + e^{-\phi}C_0^2}\,. \tag{2.13}$$

From the coefficient of $\tilde{F} \wedge \tilde{F}$, we also have

$$C_0 \qquad \rightarrow \qquad -\frac{e^{2\phi} C_0}{1 + e^{2\phi} C_0^2} = -\frac{e^{\phi} C_0}{e^{-\phi} + e^{\phi} C_0^2} \,. \tag{2.14}$$

Combining this transformation with the symmetry under a constant shift of C_0, at the classical level, one deduces that the D3-brane action has $SL(2, \mathbf{R})$ symmetry. Once semi-classical quantum effects on the D3-brane world-volume are taken into account, one finds that magnetic monopoles should be introduced as ends of the D-strings. In this case, the constant shift of C_0 is replaced by a quantized shift, much as in QCD. Thus, the $U(1)$ gauge group is compact, and the surviving symmetry group of quantum Type IIB string theory is $SL(2, \mathbf{Z})$.

3. $SL(2, \mathbf{Z})$ covariance of the D1-brane on $AdS_5 \times S^5$

We now turn to the D1-brane on $AdS_5 \times S^5$. The action of the D1-brane with κ-symmetry is given by

$$S = -\int d^2\sigma e^{-\phi} \sqrt{-\det\left(G_{ij} + F_{ij} - b_{ij}\right)} + \int e^{-\phi} C_2 \,, \tag{3.1}$$

where C_2 is determined by $dC_2 = i\bar{L}\tau_1 \hat{L} L$. We were able to check explicitly, using Mauer-Cartan equation and Fierz identity, that dC_2 is closed. Alternatively, we can make a $SO(2)$ rotation mapping τ_1 to τ_3. The resulting 3-form is the standard closed 3-form appearing in the superstring action [2]. In Eq. (3.1), we have expressed the dilaton dependence in the string frame. The κ-symmetry transformation is given by

$$\delta_\kappa x^{\hat{a}} = 0, \;\; \delta_\kappa \theta = \kappa \,, \tag{3.2}$$

where the κ-parameter satisfies

$$\Gamma \kappa = \kappa, \;\; \Gamma^2 = 1 \tag{3.3}$$

and the matrix Γ is given by

$$\Gamma = \frac{\varepsilon^{i_1 i_2}}{\sqrt{-det(G_{ij} + F_{ij} - b_{2ij})}} \left(\frac{1}{2} \tau_1 \hat{L}_{i_1} \hat{L}_{i_2} + \frac{i\tau_2}{2} (F_{i_1 i_2} - b_{2 i_1 i_2}) \right) . \tag{3.4}$$

The corresponding variation of G_{ij} and $F_{ij} - b_{ij}$ is given by

$$\delta_\kappa G_{ij} = -2i \delta_\kappa \bar{\theta} (\hat{L}_i L_j + \hat{L}_j L_i) \,, \tag{3.5}$$

$$\delta_\kappa (F - b_2) = +2i \delta_\kappa \bar{\theta} \hat{L} \tau_3 L \,. \tag{3.6}$$

One can also add a total derivative term

$$e^{-\phi} C_2 \qquad \rightarrow \qquad e^{-\phi} C_2 - C_0 F \,, \tag{3.7}$$

where C_0 is a constant "axion" background field. Since $C_0 F$ is a total derivative, it does not affect the classical equations of motion. As in the previous section, a constant shift of C_0 is a trivial classical symmetry of the action Eq. (3.1). At quantum level, the symmetry is reduced to a discrete shift symmetry, T-symmetry of $SL(2,\mathbf{Z})$.

Now we perform the duality transformation. We first introduce a Lagrange multiplier field $\tilde{H}^{\mu\nu} = -\tilde{H}^{\nu\mu}$ as follows:

$$S' = \int d^2\sigma \Big(-e^{-\phi}\sqrt{-\det\left(G_{ij} + F_{ij} - b_{ij}\right)} + +\frac{1}{2}\tilde{H}^{ij}(F_{ij} - 2\partial_i A_j) + \\ +\frac{1}{2}e^{-\phi}\epsilon^{ij}C_{ij} - \frac{1}{2}C_0\epsilon^{ij}F_{ij}\Big) \tag{3.8}$$

and regard the field strength F_{ij} to be an independent field. Varying the action with respect to A_j yields $\partial_i\tilde{H}^{ij} = 0$, which implies that $\tilde{H}^{ij} = \epsilon^{ij}\Lambda$ with Λ constant. If we use the equation of motion for F to rewrite the action in terms of Λ instead of F, the total action can be written as

$$S_D = \sqrt{e^{-2\phi} + (\Lambda - C_0)^2}\int d^2\sigma\left(-\sqrt{-\det G_{ij}} - \frac{1}{2}\epsilon^{ij}b'_{ij}\right), \tag{3.9}$$

where $db'_2 = i\bar{L}\tau'_3\hat{L}L$ and τ'_3 is defined to be

$$e^{-\phi}\,\tau_1 - (\Lambda - C_0)\,\tau_3 \equiv \sqrt{(e^{-2\phi} + (\Lambda - C_0)^2)}\,\tau'_3\,. \tag{3.10}$$

The action eq. (3.9) is the κ-symmetric superstring action in the Nambu-Goto form with the modified tension

$$T' = \sqrt{e^{-2\phi} + (\Lambda - C_0)^2}\,. \tag{3.11}$$

The κ-transformation is given by eqs. (3.2), (3.3), (3.5), but now Γ is given by

$$\Gamma = \frac{\varepsilon^{i_1 i_2}}{\sqrt{-det(G_{ij})}}\left(\frac{1}{2}\tau_3\hat{L}_{i_1}\hat{L}_{i_2}\right). \tag{3.12}$$

This result admits the same interpretation as in the flat space. The expression eq. (3.11) can be interpreted as the tension for the $SL(2,\mathbf{Z})$ covarinat spectrum of strings provided that one identifies the integer value $\Lambda = m$ as corresponding to the $SL(2,\mathbf{Z})$ dyonic string of charge $(m, 1)$ in the $AdS_5 \times S^5$ background with constant dilaton ϕ and axion C_0. An equivalent interpretation is that eq. (3.9) describes the fundamental $(1, 0)$ string with an $SL(2,\mathbf{Z})$ transformed metric, dilaton and axion. The relevant $SL(2,\mathbf{Z})$ transformations maps $C_0 + ie^{-\phi}$ to $-(C_0 - \Lambda + ie^{-\phi})^{-1}$. Thus the coupling constant of the fundamental string after the duality transformation is given by $e^{\tilde{\phi}} = e^{-\phi} + e^{\phi}(\Lambda - C_0)^2$.

4. Discussion

In this paper, using coset superspace approach, we have investigated the $SL(2,\mathbf{Z})$ duality property of the D1-and D3-brane actions on $AdS^5 \times S^5$ background. We have proven that the D1-and D3-branes behave as in the expected way under the $SL(2, \mathbf{Z})$ transformation. Moreover, we have shown that the actual dual transformation is formally the same as in the flat space. In both cases, the $SL(2, \mathbf{Z})$ acts on the fermionic coordinates in the same way, viz, as the $SO(2)$ rotation. In particular, with the dilaton and axion field background turned on, we have shown that the full $SL(2, \mathbf{Z})$ duality symmetry follows from the $SO(2)$ duality rotation. Throughout this paper, our discussion has been mainly semi-classical. It would be interesting to extend our proof to the full-fledged quantum level.

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