$SL(2, Z)$ S-duality of Super D-string
in
Type IIB Supergravity Background

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Abstract

It is shown in a quantum-mechanically exact manner that a supersymmetric and $\kappa$-symmetric D-string action in a general type IIB supergravity background is transformed to a form of the type IIB Green-Schwarz superstring action with the $SL(2, Z)$ covariant tension through an S-duality transformation. This result precisely proves a conjecture mentioned previously that the $SL(2, Z)$ S-duality of a super D-string action in a flat background is also valid even in a curved IIB background geometry. We point further out the validity of the more generalized conjecture that various duality relations of super D-brane and M-brane actions originally found in a flat background also hold true in general ten dimensional type II supergravity and eleven dimensional supergravity background geometries by applying the present formalism to those cases.

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1 Introduction

Symmetries have always played a central role in theoretical physics. For instance, gauge symmetries are used not only to determine dynamics, the gauge interactions, but also to prove unitarity and renormalizability of a theory. Thus the importance of a knowledge of symmetries in modern theoretical physics is recognized as one of key ingredients in establishing a theory.

However, in superstring theory and M-theory, symmetries have remained elusive so that many vital issues have had to be left to await the advent of an understanding of symmetries. Such an understanding has begun to emerge recently owing to developments of some non-perturbative techniques although we are still quite a way (it seems) from addressing what ultimate symmetries in superstring theory and M-theory are.

In this respect, over recent years a large amount of evidence has accumulated in support of various duality relations of superstring theory and M-theory. Following the evidence, these theories have a huge group of global discrete symmetries. A particularly interesting subgroup is the $SL(2,\mathbb{Z})$ S-duality [1] of type IIB superstring theory [2] which exchanges an infinite family of solitons and bound states. It was already known that the IIB supergravity [3, 4], which is the low energy effective theory of the type IIB superstring theory, possesses a discrete global symmetry, the $SL(2,\mathbb{R})$ symmetry, but this symmetry was regarded as an artifact of the low energy approximation to underlying renormalizable theory and not taken seriously in those days. But nowadays the situation has radically changed since the discovery of the still rather mysterious more fundamental theory where $SL(2,\mathbb{Z})$ subgroup of $SL(2,\mathbb{R})$ is expected to be an exact symmetry [5]. Hence it is an important and urgent issue to obtain a fuller understanding of the $SL(2,\mathbb{Z})$ symmetry when one tries to clarify aspects of symmetries behind the more fundamental underlying theory.

In this paper, we wish to consider a derivation of the $SL(2,\mathbb{Z})$ S-duality [1] of type IIB superstring theory in a general type IIB supergravity background [6, 7]. The conjecture that the $SL(2,\mathbb{Z})$ S-duality [1] of type IIB superstring theory and other dualities of D-brane and M-brane actions, which were found originally in a flat background may be true even in a general curved on-shell background, was already stated in the final section of the papers [8, 9], but only recently was shown in the case of super D-string and super D3-brane actions on a specific $AdS_5 \times S^5$ background where it was strongly suggested that this conjecture would be valid even in a general on-shell supergravity background [10, 11] (See also an interesting work [12] where the gauge-fixing of $\kappa$-symmetry was not used).

The main purpose of this paper is to report that this conjecture is indeed the case. Because of the character as a short article, we confine ourselves to only the case of a super D-string action in a general type IIB supergravity background and present essential ideas and techniques to some extent, but it is quite straightforward to extend the ideas and techniques presented in this paper to the broader situations, namely a proof of various duality properties of super D-brane actions in ten dimensional type IIA and IIB supergravity background, and of M2 and M5-brane actions in eleven dimensional supergravity background such as $SL(2,\mathbb{Z})$ self-duality of D3-brane, D2 vs. M2-brane duality and D4 vs. M5-brane duality. These problems will be reported in a longer paper in detail [13].
In this paper, as method of derivation of the $SL(2, Z)$ S-duality in a super D-string action, we rely on path integral of the first-order Hamiltonian form, which was used in [14] for a bosonic string and in [13] for a superstring in a flat Minkowskian background, and also in [10] for a superstring on $AdS_5 \times S^5$.

The contents of this article are organized as follows. In Section 2 we shall review a super D-string action in a general IIB supergravity background [6, 7]. In Section 3 it will be shown that the super D-string action on this background is transformed to the type IIB Green-Schwarz superstring action with the $SL(2, Z)$ covariant tension in the same background through an S-duality transformation. The final section will be devoted to discussions.

2 Super D-string action in a general IIB background

We start by reviewing a super D-string action in a general IIB supergravity background [6, 7]. It is well known nowadays that super D-brane actions consist of two terms, those are, the Dirac-Born-Infeld action and the Wess-Zumino action. The former includes the NS-NS two-form and dilaton in addition to world-volume metric while the latter action contains the coupling of the D-brane to the R-R fields. In particular, a super D-string action in a general IIB on-shell supergravity background is given by

$$S = S_{DBI} + S_{WZ},$$

with

$$S_{DBI} = - \int_{M_2 = \partial M_3} d^2 \sigma \sqrt{- \det (G_{ij} + \mathcal{F}_{ij})},$$

$$S_{WZ} = \int_{M_2 = \partial M_3} C_2 = \int_{M_3} \Omega_3,$$

where we have defined as

$$\mathcal{F} = F - b_2,$$

$$F = dA,$$

$$\Omega_3 = dC_2.$$  \hfill (3)

Moreover, following the paper [3], we define the NS-NS 3-form superfield $H_3$ of the NS-NS 2-form $b_2$ and the R-R $n$-form superfield $R$ as

$$H_3 = db_2,$$

$$R = e^{b_2} \wedge d(e^{-b_2} \wedge C) = \bigoplus_{n=1}^{10} R_{(n)},$$

$$C = \bigoplus_{n=0}^{9} C_{(n)},$$  \hfill (4)
but in the case of string, it is easy to show that the R-R 3-form field strength superfield $R^{(3)}$, which is relevant below, coincides with the Wess-Zumino form $\Omega_3$. Of course, from the definitions (4) these field strengths obey the following Bianchi identities

$$dH_3 = 0,$$
$$e^{b_2} \wedge d(e^{-b_2} \wedge R) = dR - R \wedge H_3 = 0.$$  \hfill (5)

In order to reduce the enormous unconstrained field content included in the superfields to the field content of the on-shell type IIB supergravity theory, one has to impose the constraints on the field strengths by hand, which make various Bianchi identities to coincide with the equations of motion of supergravity [4]. In this section, for the sake of clarity, we shall limit ourselves to be the case of the vanishing dilaton superfield. Later we will consider the constant dilaton background when we derive the $SL(2,\mathbb{Z})$ covariant tension. The case of a general dilaton superfield will be discussed at the end of next section. Under the assumption of the vanishing (or constant) dilaton, the nontrivial constraints imposed on the field strength components [3] reduce to simple forms

$$T^{c}_{\alpha\beta} = 2i\gamma_{\alpha\beta},$$
$$H_{\alpha\beta} = -2i(\mathcal{K}\gamma_\alpha)_{\alpha\beta},$$
$$\Omega_{\alpha\beta} = 2i(\mathcal{I}\gamma_\alpha)_{\alpha\beta},$$ \hfill (6)

where $T^{c}_{\alpha\beta}$ indicates a component of the torsion superfield $T_{ABC}$. And $\mathcal{E}$, $\mathcal{I}$, and $\mathcal{K}$ describing the $SO(2)$ matrices are defined in terms of the conventional Pauli matrices

$$\mathcal{E} = i\sigma_2 = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}, \quad \mathcal{I} = \sigma_1 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \quad \mathcal{K} = \sigma_3 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}.$$  \hfill (7)

From these constraints, it is easy to obtain the NS-NS 3-form superfield $H_3$ and the 3-form Wess-Zumino term $\Omega_3$ whose concrete expressions are given by

$$H_3 = i\tilde{E} \wedge \hat{E} \wedge K E,$$
$$\Omega_3 = -i\tilde{E} \wedge \hat{E} \wedge I E,$$ \hfill (8)

where $E$ implies an appropriate component of the vielbein one-form $E^A = dZ^A E^A_M$. Writing down in an explicit manner, $\tilde{E}$, $\hat{E}$ and $E$ in (8) indicate the Dirac adjoint of $E^{\dagger}_\alpha$, $E^\alpha \gamma_\alpha$ and $E^{\dagger}\alpha$, respectively. Here we have defined the curved super-index as $M = (m, \mu)$ and the flat super-index $A = (a, \alpha)$, and $I$ as the $N = 2$ index taking 1 and 2.

For later convenience, let us recapitulate the results obtained in this section. The $\kappa$-symmetric and world-sheet reparametrization invariant super D-string action in a IIB on-shell supergravity background is given by

$$S = S_{DBI} + S_{WZ},$$
$$S_{DBI} = -\int_{M_2} d^2\sigma \sqrt{-\det(G_{ij} + \mathcal{F}_{ij})},$$
\[
S_{WZ} = \int_{M_2=\partial M_3} C_2 = \int_{M_3} \Omega_3, \\
\mathcal{F} = F - b_2, \\
H_3 = db_2 = i\hat{E} \wedge \hat{E} \wedge KE, \\
\Omega_3 = dC_2 = -i\hat{E} \wedge \hat{E} \wedge \Omega E.
\]

At this stage, it is interesting to notice that these equations are quite similar to those of a super D-string action on \(AdS_5 \times S^5\) whose observation actually made it possible to achieve the present study.

### 3 \(SL(2,\mathbb{Z})\) S-duality

Now let us turn our attention to a proof of \(SL(2,\mathbb{Z})\) S-duality of a super D-string action in a general ten dimensional IIB supergravity background. As mentioned just above, since the action is written to be a form similar to that on \(AdS_5 \times S^5\), we can follow a similar path of derivation, but we will expose the procedure to some extent for completeness. The main focus is paid on the point of how one performs the \(SO(2)\) spinor rotation in moving from the super D-string action to the Green-Schwarz action through the orthogonalization of the \(SO(2)\) matrix.

Now we are ready to show how the super D-string action (9) becomes a fundamental Green-Schwarz superstring action with the \(SL(2,\mathbb{Z})\) covariant tension by using the path integral of the first-order Hamiltonian form [14, 15, 10].

According to the Hamiltonian formalism, let us start by introducing the canonical conjugate momenta \(\pi^i\) corresponding to the gauge field \(A_i\) defined as

\[
\pi^i = \frac{\partial S}{\partial \dot{A}_i} = \frac{\partial S_{DBI}}{\partial \dot{A}_i},
\]

where we used the fact that the Wess-Zumino term is independent of the gauge potential, which holds only in the case of string theory. Then the canonical conjugate momenta \(\pi^i\) are calculated to be

\[
\pi^0 = 0, \quad \pi^1 = \frac{\mathcal{F}_{01}}{\sqrt{-\det(G_{ij} + \mathcal{F}_{ij})}},
\]

where the former equation just shows the existence of the \(U(1)\) gauge invariance. From these equations we will see that the Hamiltonian density takes the form

\[
\mathcal{H} = \sqrt{1 + (\pi^1)^2} \sqrt{-\det G_{ij} - A_0 \partial_1 \pi^1 + \partial_1 (A_0 \pi^1) + \pi^1 b_{01} - C_{01}},
\]

Now the partition function is defined by the first-order Hamiltonian form with respect to only the gauge field as follows:

\[
Z = \frac{1}{\mathcal{D}\pi^0} \int \mathcal{D}\pi^0 \mathcal{D}\pi^1 \mathcal{D}A_0 \mathcal{D}A_1 \exp i \int d^2\sigma (\pi^1 \partial_0 A_1 - \mathcal{H})
\]
\[ \int \mathcal{D}A_1 \mathcal{D}A_0 \exp i \int d^2 \sigma \left[ -A_1 \partial_0 \pi^1 + A_0 \partial_1 \pi^1 - \sqrt{1 + (\pi^1)^2} - \det G_{ij} \pi^1 b_{01} + C_{01} - \partial_1 (A_0 \pi^1) \right]. \] (13)

Provided that we take the boundary conditions for \( A_0 \) such that the last surface term in the exponential identically vanishes, then we can carry out the integrations over \( A_i \) explicitly, which gives rise to \( \delta \) functions

\[ Z = \int \mathcal{D} \pi^1 \delta(\partial_0 \pi^1) \delta(\partial_1 \pi^1) \exp i \int d^2 \sigma \left[ -\sqrt{1 + (\pi^1)^2} \right. \]
\[ \left. \sqrt{- \det G_{ij} + C_{01} - \pi^1 b_{01}} \right]. \] (14)

The existence of the \( \delta \) functions reduces the integral over \( \pi^1 \) to the one over only its zero-modes. If we require that one space component is compactified on a circle, these zero-modes are quantized to be integers \( \{16\} \). As a consequence, the partition function becomes

\[ Z = \sum_{m \in \mathbb{Z}} \exp i \int d^2 \sigma \left[ -\sqrt{1 + m^2} \right. \]
\[ \left. \sqrt{- \det G_{ij} + C_{01} - mb_{01}} \right], \] (15)

from which we can read off the effective action

\[ S = \int d^2 \sigma \left( -\sqrt{1 + m^2} \right. \]
\[ \left. \sqrt{- \det G_{ij} - i \int_{M_3} \hat{E} \wedge \hat{K} \wedge \mathcal{E}} \right). \] (16)

Moreover, recalling the relation stemming from \( \{4\} \)

\[ \int_{M_3 = \partial M_3} d^2 \sigma (C_{01} - mb_{01}) = \int_{M_3} (\Omega_3 - mH_3) = -i \int_{M_3} E \wedge \hat{E} \wedge (m \mathcal{K} + \mathcal{I})E, \] (17)

and then carrying out an orthogonal transformation

\[ U^T (m \mathcal{K} + \mathcal{I}) U = -\sqrt{1 + m^2} \mathcal{K}, \] (18)

with an orthogonal matrix \( U = \frac{1}{\sqrt{1 + (m - \sqrt{1 + m^2})^2}} [(m - \sqrt{1 + m^2})1 - \mathcal{E}] \), one finally obtains the action

\[ S = -\sqrt{1 + m^2} \left( \int_{M_2} d^2 \sigma \sqrt{- \det G_{ij} - i \int_{M_3} \hat{E} \wedge \hat{K} \wedge \mathcal{E}} \right). \] (19)

This is nothing but type IIB Green-Schwarz superstring action with the modified tension \( \sqrt{1 + m^2} \) in a type IIB supergravity background \( \{17\} \). (Note that in \( \{17\} \) a complex formalism is used while in this article a real formalism is used, but by comparison of the two formalisms we can convince ourselves that the action \( \{19\} \) corresponds to the Green-Schwarz action in the complex formalism.)

It is worthwhile to notice that the result obtained above agrees with the tension formula for the \( SL(2, \mathbb{Z}) \) S-duality spectrum of strings in the flat background \( \{1\} \) provided that we identify the integer value \( \pi^1 = m \) as corresponding to the \((m, 1)\) string. To show more clearly
that the tension at hand is the $SL(2, Z)$ covariant tension, it would be more convenient to start with the following classical action

$$S = -n \int_{M^2} d^2 \sigma \left[ e^{-\phi} \sqrt{\det(G_{ij} + F_{ij}) - C_2} + \frac{1}{2} \epsilon^{ij} \chi F_{ij} \right],$$

where $n$ is an integer, and we have introduced the constant dilaton $\phi$ and the constant axion $\chi$. Then following the same path of derivation as above, we can obtain the manifestly $SL(2, Z)$ covariant tension

$$T = \sqrt{(m + n\chi)^2 + n^2 e^{-2\phi}}.$$

Here we would like to comment two important points. One point is that we have shown that there exists $SL(2, Z)$ S-duality in type IIB superstring theory even in a general type IIB supergravity background without reference to any approximation. Thus this relation is quantum-mechanically exact.

The other point is the problem of whether one can interpret the orthogonal transformation (18) as the $SO(2)$ rotation of the $N = 2$ spinor coordinates. In our previous paper [10, 11] this problem was emphasized too much, but on reflection it turns out that this problem is rather trivial by the following reason. Notice that the torsion constraint in (6) is obviously invariant under this rotation. Moreover, since we require that the original super D-string action and the fundamental Green-Schwarz action reduce to the well-known forms of the corresponding flat space actions in the flat space limit, $E^I_\alpha$ with the $SO(2)$ index $I$ and $E^a_\alpha$ must take the following forms at the lowest order expansion with respect to the spinor coordinates $\theta$

$$E^I_\alpha = \partial_\alpha \theta^I + \ldots,$$

$$E^a_\alpha = \partial_\alpha X^a - i \bar{\theta}^I \gamma_\alpha^a \partial_\alpha \theta^I + \ldots,$$

where the dots indicate the higher order terms reflecting the curved nature of the background metric. These facts mean that $E^I$ transforms as the adjoint representation of the $SO(2)$ group, on the other hand, $E^a$ must be invariant under an $SO(2)$ rotation. Accordingly, we can understand that the orthogonal transformation (18) is indeed performed by an $SO(2)$ rotation of the $N = 2$ spinor coordinates. In this way, we have succeeded in deriving the $SL(2, Z)$ S-duality of type IIB superstring theory in type IIB on-shell supergravity background at least within the present context.

To close this section, let us discuss the case of a general non-constant dilaton superfield. We will see below that even if a formulation becomes a little complicated compared with the case of the vanishing (or constant) dilaton, the essential point remains unchanged. In this case, the relevant constraints are given by [12]

$$T^c_{\alpha \beta} = 2i \gamma^c_{\alpha \beta},$$

$$H_{a a \alpha} = -2ie^{\frac{i}{2} \phi}(K_{\gamma_a})_{\alpha \beta},$$

$$H_{a b \alpha} = e^{\frac{i}{2} \phi}(\gamma_{a b} K \Lambda)_{\alpha},$$
where \( \Lambda_\alpha = \frac{1}{2} \partial_\alpha \phi \). These constraints give rise to the NS-NS 3-form superfield \( H_3 \) and the 3-form Wess-Zumino term \( \Omega_3 \) whose concrete expressions are of form
\[
H_3 = db_2 = i e^{\frac{1}{2} \phi} \hat{E} \land \hat{K} E + \frac{1}{2} e^{\frac{1}{2} \phi} \hat{E} \land \gamma_{ab} \Lambda E^b \land E^a,
\]
\[
\Omega_3 = d(C_2 + C_0 \mathcal{F}) = -i e^{-\frac{1}{2} \phi} \hat{E} \land \hat{K} E - \frac{1}{2} e^{-\frac{1}{2} \phi} \hat{E} \land \gamma_{ab} \Lambda E^b \land E^a + 2 e^{-\phi} \hat{E} \mathcal{E} \Lambda \land \mathcal{F}.
\]
Here notice that the constraints (23) and the 3-form superfields (24) exactly reduce to the previous expressions (6) and (8), respectively when the dilaton superfield \( \phi \) is vanishing.

If we start with the classical action [9]
\[
S = S_{DBI} + S_{WZ},
\]
\[
S_{DBI} = -\int_{M_2} d^2 \sigma e^{-\frac{1}{2} \phi} \sqrt{-\det(G_{ij} + e^{-\frac{1}{2} \phi} \mathcal{F}_{ij})},
\]
\[
S_{WZ} = \int_{M_2=\partial M_3} (C_2 + C_0 \mathcal{F}) = \int_{M_3} \Omega_3,
\]
then the same argument as in the previous section leads to a dual action
\[
S = -\int_{M_2} d^2 \sigma e^{-\frac{1}{2} \phi} \sqrt{1 + (\lambda + C_0)^2 e^{2 \phi}} \sqrt{-\det G_{ij}}
\]
\[
+ \int_{M_3} e^{-\frac{1}{2} \phi} \sqrt{1 + (\lambda + C_0)^2 e^{2 \phi}} \left( i \hat{E} \land \hat{K} E + \frac{1}{2} \hat{E} \gamma_{ab} \Lambda \land E^b \land E^a \right),
\]
where \( \lambda \) is a constant scalar superfield. In this way we can also show the \( SL(2, Z) \) S-duality in the general non-constant dilaton superfield.

4 Discussions

In this paper, we have studied the property of the \( SL(2, Z) \) duality of a supersymmetric and \( \kappa \)-symmetric D-string action in a general type IIB on-shell supergravity background. It has clearly shown that the \( SL(2, Z) \) duality of type IIB superstring theory found originally in a flat space-time background holds true even in on-shell supergravity background, which was anticipated previously but has not proved so far. This fact is quite illuminating from the viewpoint that the \( SL(2, Z) \) duality is expected to be an exact symmetry of the underlying fundamental theory [3] as mentioned in the introduction. Maybe, one of the most challenging studies in future would be to promote this global discrete symmetry to the local gauge
symmetry, from which we could understand the relation between the strong coupling phase and the weak coupling phase, and the Kaluza-Klein compactification e.t.c.

Moreover, we have spelled out the problem of the $SO(2)$ rotation of the $N = 2$ spinor coordinates. Our proof utilizes only an invariance of the constraints and the boundary condition in the flat background limit so that it can be applied to the other situations in a straightforward way.

In a coming longer paper $[13]$, we will prove various duality relations in detail where D2 vs. M2 duality, the self-duality of D3-brane and D4 vs. M5 duality are discussed in a general on-shell supergravity background. As a final goal, it would be wonderful to prove the existence of duality transformations in the background independent matrix models in future $[18]$.

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