More on the Gauge-Fixed D3-Brane Action with Dilaton-Axion Coupling from N=1 Superspace

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Abstract

The gauge-fixed action of a ‘spacetime-filling’ D3-brane with dilaton-axion coupling is formulated in N=1 superspace. We investigate its symmetries by paying special attention to a possible non-linearly realized extra supersymmetry, and emphasize the need of a linear superfield coupled to an abelian Chern-Simons superfield to represent a dilaton-axion supermultiplet in the off-shell manifestly supersymmetric approach.

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1 Introduction

The supersymmetric D-brane actions with local fermionic kappa symmetry were constructed in ref. [1]. When the kappa-symmetry is fixed, half of supersymmetry is spontaneously broken, whereas the fermionic superpartner (with respect to unbroken half of supersymmetry) of the $U(1)$ gauge field in the D-brane worldvolume can be identified with the Goldstone fermion. The most relevant part of the gauge-fixed D-brane action is given by a supersymmetric Born-Infeld (BI) action [1]. Gauge-fixing results in the D-brane actions whose all supersymmetries are non-linearly realized, i.e. non-manifest. Unbroken supersymmetries can sometimes be made manifest by using superspace [2, 3].

The electric-magnetic self-duality of the BI action can be extended to a full $SL(2,\mathbb{Z})$ duality in the case of a gauge-fixed ‘spacetime-filling’ D3-brane with axion-dilaton coupling [4]. This feature can be made manifest when considering the D3-brane action as the double dimensionally reduced M5-brane action on a 2-torus [5]. The dilaton-axion can be identified with the complex structure of the torus, while the $SL(2,\mathbb{Z})$ self-duality of a D3-brane is then nothing but the modular group of the torus [5]. In this Letter we make manifest the unbroken N=1 supersymmetry of the spacetime-filling D3-brane action with dilaton-axion coupling, and investigate its other relevant symmetries in flat N=1 superspace.

2 N=1 BI action in superspace

In this section we briefly describe the N=1 BI action is superspace, which is the prerequisite to our investigation in sect. 3. The BI action in Minkowski spacetime of signature $\eta = \text{diag}(+,-,-,-)$ is [6]

$$S_{\text{BI}} = \frac{1}{\kappa^2} \int d^4x \sqrt{-\det(\eta_{mn} + \kappa F_{mn})} ,$$

(1)

where $F_{mn} = \partial_m A_n - \partial_n A_m, \ m, n = 0, 1, 2, 3, \text{ and } \kappa$ is the dimensional coupling constant ($\kappa = 2\pi\alpha'$ in string theory). The N=1 supersymmetric extension of the action (1) can be interpreted as the Goldstone-Maxwell action associated with partial (1/2) spontaneous supersymmetry breaking, N=2 to N=1, whose Goldstone fermion is photino of a Maxwell (vector) N=1 multiplet with respect to unbroken N=1 supersymmetry [2, 3]. Manifest supersymmetry does not respect the standard form (1) of the BI action. The complex bosonic variable, having the most natural supersymmetric
extension, is given by
\[ \omega = \alpha + i\beta, \quad \alpha = \frac{1}{4} F_{mn} F_{mn}, \quad \beta = \frac{1}{4} F_{mn} \tilde{F}_{mn}, \quad \tilde{F}_{mn} = \frac{1}{2} \varepsilon^{mnpq} F_{pq}. \]

The BI Lagrangian (1) can be rewritten in terms of \( \omega \) and \( \bar{\omega} \) as
\[ \mathcal{L}_{\text{BI}}(\omega, \bar{\omega}) = \mathcal{L}_{\text{free}} + \mathcal{L}_{\text{int.}} \equiv -\frac{1}{2} (\omega + \bar{\omega}) + \kappa^2 \omega \bar{\omega} \mathcal{Y}(\omega, \bar{\omega}), \]
where the particular structure function \( \mathcal{Y}(\omega, \bar{\omega}) \) has been introduced,
\[ \mathcal{Y}(\omega, \bar{\omega}) = \frac{1}{1 + \frac{\kappa^2}{2} (\omega + \bar{\omega}) + \sqrt{1 + \kappa^2 (\omega + \bar{\omega}) + \frac{\kappa^4}{4} (\omega - \bar{\omega})^2}}. \]

A supersymmetrization of the bosonic BI theory (1) in the form (3) amounts to replacing the field strength \( F_{mn} \) by the N=1 chiral spinor superfield strength \( W_\alpha \), and \( \omega \) by the N=1 chiral scalar superfield \( \bar{K} = \frac{1}{8} \bar{D}^2 \bar{W}^2 \), viz.
\[ S_{\text{sBI}} = \frac{1}{4} \left( \int d^4 x d^2 \theta W^2 + \text{h.c.} \right) + \frac{\kappa^2}{8} \int d^4 x d^4 \theta W^2 \bar{W}^2 \mathcal{Y}(K, \bar{K}) \]
with the same structure function (4), so that the bosonic terms of eq. (5) exactly reproduce eq. (1). We use the standard notation, \( W^2 = W_\alpha W^\alpha \) and \( \bar{W}^2 = \bar{W}_{\bar{\alpha}} \bar{W}^{\bar{\alpha}} \), and similarly for the N=1 flat superspace covariant derivatives \( D^\alpha \) and \( \bar{D}^\alpha \) with \( \alpha = 1,2 \) and \( \bar{\alpha} = 1,2 \). The gauge superfield strength \( W_\alpha \) obeys the superfield Bianchi identities
\[ \bar{D}^\alpha W_\alpha = 0 \quad \text{and} \quad D^\alpha W_\alpha = \bar{D}^\bar{\alpha} \bar{W}^{\bar{\alpha}}. \]

In the chiral basis the gauge superfield strength reads
\[ W_\alpha(x, \theta) = -i \psi_\alpha(x) + \left[ \delta^\beta_\alpha D(x) - i (\sigma^{mn})^\beta_\alpha F_{mn}(x) \right] \theta_\beta + \theta^2 (\sigma^m \partial_m)_{\alpha\beta} \psi^\beta(x), \]
where \( \psi_\alpha(x) \) is the fermionic superpartner (photino) of the abelian BI vector field \( A_m \), and \( D \) is the real auxiliary field. In the N=1 super-BI theory (5) setting \( D = 0 \) is consistent with its equations of motion (this is called the ‘auxiliary freedom’ [7]).

The action (5) can be put into the simple ‘non-linear sigma-model’ form [2, 3]
\[ S_{\text{sBI}} = \int d^4 x d^2 \theta X + \text{h.c.}, \]
whose chiral superfield Lagrangian \( X \) is determined via the recursive relation [2, 3]
\[ X + \frac{\kappa^2}{8} X \bar{D}^2 \bar{X} = \frac{1}{4} W^2 W_\alpha. \]

The BI action (1) is well-known to be invariant under non-trivial electric-magnetic duality [8]. This means that treating \( F \) as a generic two-form, enforcing the Bianchi
identity, $dF = 0$, by means of a Lagrange multiplier (= dual vector potential) in the first-order action, and integrating out $F$ in favor of the Lagrange multiplier yield the dual action having the same form as eq. (1) in terms of the dual vector potential. The same is true in $N=1$ superspace for the action (5) when introducing the dual $N=1$ superfield strength as an $N=1$ Lagrange multiplier, and integrating over $W$ in the corresponding first-order action, i.e. after the $N=1$ superfield Legendre transform [3].

Another highly non-trivial property of eq. (5) is its invariance under the (non-linearly realized and spontaneously broken) second $N=1$ supersymmetry with rigid spinor parameter $\eta_\alpha$ [2],

$$\delta_\eta W_\alpha = \frac{1}{\kappa} \eta_\alpha + \frac{\kappa^2}{4} \bar{D}^2 \bar{X} \eta_\alpha + i \kappa (\sigma^m \bar{\eta})_\alpha \partial_m X . \quad (10)$$

The transformations (10) are consistent with the $N=1$ Bianchi identities (6), and they realize a supersymmetry algebra. The invariance of the action (5) under the transformations (10) follows from the remarkable fact that $\int d^2 \theta \delta_\eta X = \frac{1}{4\pi} \int d^2 \theta W^\alpha \eta_\alpha$ is a total derivative in spacetime.

To make manifest the hidden second supersymmetry of the the $N=1$ BI theory, one can reformulate it in the formalism of non-linear realizations [9]. The Goldstone superfield $\Psi$ having the standard transformation law in the chiral version of the non-linearly realized supersymmetry [10], $\delta_\eta \Psi = \frac{1}{\kappa} \eta - 2i \kappa (\Psi \sigma^m \bar{\eta}) \partial_m \Psi$, is given by

$$\Psi_\alpha = \frac{W_\alpha}{1 + \kappa^2 \bar{D}^2 \bar{X}} + \ldots , \quad (11)$$

where the dots stand for the higher-order fermionic terms [9]. The new Goldstone superfield $\Psi$ obeys the non-linear $N=1$ superspace constraints

$$\mathcal{D}_\alpha \Psi_\alpha = \mathcal{D}_\alpha \bar{\Psi}_\alpha = 0 \quad (12)$$

that are also covariant under the second non-linearly realized supersymmetry. The $N=2$ covariant derivatives in $N=1$ superspace [2]

$$\mathcal{D}_\alpha = D_\alpha + i \kappa^2 (D_\alpha \Psi \sigma^m \bar{\Psi} + D_\alpha \bar{\Psi} \bar{\sigma}^m \Psi) D_m \quad \text{and} \quad D_m = (\omega^{-1})_m \partial_n , \quad (13)$$

where $\omega_m^n = \delta_m^n - i \kappa^2 (\partial_m \Psi \sigma^n \bar{\Psi} + \partial_m \bar{\Psi} \bar{\sigma}^n \Psi)$, form a closed algebra. The action (5) may be rewritten in terms of $\Psi$ and the $N=2$ covariant derivatives (13) as

$$S_{BI} = \frac{1}{4} \int d^4 x d^2 \theta \bar{\epsilon}^{-1} \Psi^2 + \text{h.c.} , \quad (14)$$

whose $N=1$ chiral superfield $\bar{\epsilon}^{-1} = 1 + \frac{\kappa^4}{4} \bar{D}^2 \bar{X} + \ldots$, should transform as a density under the second supersymmetry, $\delta_\eta \bar{\epsilon}^{-1} = -2i \kappa \partial_m (\bar{\epsilon}^{-1} \Psi \sigma^m \bar{\eta})$. 

4
Both the electric-magnetic self-duality and the second non-linearly realized supersymmetry of the N=1 BI action may have been expected from its anticipated connection to the D3-brane action. It is just these key properties that allow one to identify the N=1 BI action with the low-energy effective action of the spacetime-filling D3-brane in the case of slowly varying fields. Any direct gauge-fixing of the kappa-symmetric D3-brane action [1] would yield highly involved supersymmetry transformations, whose precise relation to the standard N=1 superspace transformations implies complicated field redefinitions. We didn’t attempt to establish this connection explicitly.

3 N=1 BI action with dilaton-axion coupling

The bosonic BI action coupled to a background dilaton $\phi$ and axion $C$ reads

$$S_{\text{bosonic}} = \frac{1}{4\pi} \int d^4x \sqrt{-\det(\eta_{mn} + e^{-\phi/2}F_{mn})} + \frac{1}{32\pi i} \varepsilon^{mnpq} CF_{mn}F_{pq}. \quad (15)$$

The dilaton-axion background now plays the role of the effective coupling constant, so that we chose $\kappa = 1$ for simplicity. We also rescaled the BI action by a factor of $4\pi$, in order to make it invariant under the T-duality transformations, $C \rightarrow C + n$, where $n \in \mathbb{Z}$, because $C$ multiplies the topological density in eq. (15).

It is not difficult to supersymmetrize eq. (15) in N=1 superspace, by using the results of sect. 2. First, let’s define a complex scalar

$$\rho = e^{-\phi} + iC, \quad (16)$$

and assume that it belongs to an N=1 chiral superfield,

$$\Phi = \rho + \theta^\alpha \lambda_\alpha + \theta^2 F, \quad (17)$$

where we have introduced the physical dilatino $\lambda_\alpha$ and the ‘auxiliary’ field $F$. This is not quite innocent procedure in the theories with higher derivatives, because the field $F$ should be truly auxiliary or, at least, $F = 0$ should be a solution to the equations of motion (the auxiliary freedom). Equation (5) implies the N=1 supersymmetric extension of eq. (15) in the form

$$4\pi S = \frac{1}{4} \left( \int d^4xd^2\theta \Phi W^2 + h.c. \right) + \frac{1}{32} \int d^4xd^4\theta (\Phi + \bar{\Phi})^2 W^2\bar{W}^2 \mathcal{Y} \left( \frac{1}{2}(\Phi + \bar{\Phi})K, \frac{1}{2}(\Phi + \bar{\Phi})\bar{K} \right), \quad (18)$$
with the same function $\mathcal{Y}$ defined by eq. (4) at $\kappa = 1$, where we have used the identity

$$D^2W^2 - \bar{D}^2\bar{W}^2 = ie^{mpq}F_{mn}F_{pq}.$$  \hfill (19)

The N=1 Legendre transform of the action (18) with respect to the gauge superfield $W$ yields the dual N=1 superspace action that has the same form (18) in terms of the dual N=1 superfield strength and the dual coupling

$$\tilde{\Phi} = \frac{1}{\Phi}.$$  \hfill (20)

Together with imaginary shifts of $\Phi$ by integers the S-duality transformation (20) generates the full $SL(2,\mathbb{Z})$ duality, as required. In fact, the action (18) is invariant under the continuous $SL(2,\mathbb{R})$ duality, as it belongs to the class of the $SL(2,\mathbb{R})$ duality invariant actions constructed in ref. [11]. Of course, in quantum theory only $SL(2,\mathbb{Z})$ survives.

The $SL(2,\mathbb{R})$ duality invariant dilaton and axion kinetic terms to be added to eq. (18),

$$\mathcal{L}(\phi, C) = \frac{1}{2}(\partial_m \phi)^2 + \frac{1}{2}e^{2\phi}(\partial_m C)^2,$$  \hfill (21)

are given by the Kähler non-linear sigma-model with a Kähler potential

$$K(S, \bar{S}) = -\ln(S + \bar{S}).$$  \hfill (22)

The N=1 supersymmetrization of eq. (21) in superspace is straightforward,

$$S_{\text{kin.}} = -\int d^4x d^4\theta \ln(S + \bar{S}).$$  \hfill (23)

There is, however, a problem with another (non-linearly realized) supersymmetry. A variation of the leading terms in eq. (18) yields

$$\delta_\eta \mathcal{L} = \frac{1}{2} \int d^2\theta \Phi W^\alpha \eta_\alpha + \text{h.c.},$$  \hfill (24)

which is a total derivative only for a constant dilaton-axion background $\Phi$. Yet another problem is the auxiliary freedom of $F$.

The way out of both problems may be the assignment of dilaton and axion to an N=1 linear multiplet $G$, instead of the N=1 chiral multiplet $\Phi$. As regards the bosonic action (15), this means trading $C$ against a gauge two-form $B$, at the expense of giving up the manifest $U(1)$ gauge invariance, viz.

$$\int CF \wedge F = -\int dC \wedge (A \wedge F) = \int^* dB \wedge \Theta,$$  \hfill (25)
where the star denotes the Poincaré dual, \(* (dC) = dB\) and \(\Theta = A \wedge F\) is the abelian Chern-Simons three-form. In N=1 superspace a real linear superfield \(G\) is defined by the constraints

\[
D^2 G = \bar{D}^2 G = 0 .
\]  
(26)

It consists of a real scalar (dilaton), an antisymmetric tensor \((B)\) subject to the gauge transformation \(\delta B = d\xi\) with the one-form gauge parameter \(\xi\), a dilatino \(\lambda\), and no auxiliary fields. The two-form \(B\) enters the superfield \(G\) only via its field strength \(dB\).

The leading term in eq. (18) can then be rewritten to the form

\[
\frac{1}{4} \left( \int d^4 x d^2 \theta \Phi W^2 + \text{h.c.} \right) = \frac{1}{4} \int d^4 x d^4 \theta (\Phi + \bar{\Phi}) \Omega ,
\]  
(27)

where we have introduced the Chern-Simons superfield \(\Omega\) via the equations

\[
W^2 = \frac{1}{2} \bar{D}^2 \Omega , \quad \bar{W}^2 = \frac{1}{2} D^2 \Omega .
\]  
(28)

By using a solution \(W_\alpha = -\frac{1}{2} \bar{D}^2 D_\alpha V\) to the Bianchi identities (6), in terms of the real gauge scalar superfield \(V\) subject to the gauge transformations \(V \to V + i(\Lambda - \bar{\Lambda})\), with \(D_\alpha \Lambda = 0\), we easily find \(\Omega = -\frac{1}{4} (D^\alpha V) W_\alpha + \text{h.c.}\).

The full action given by a sum of eqs. (18) and (23) is now dependent upon the chiral superfields \(\Phi\) and \(\bar{\Phi}\) only through their linear combination \(\frac{1}{2} (\Phi + \bar{\Phi})\), so that it is possible to dualize this action in terms of the linear superfield \(G\) by Legendre transform.\(^{2}\) We replace in eqs. (18) and (23) the combination \(\frac{1}{2} (\Phi + \bar{\Phi})\) by a general real superfield \(U\), and add extra term

\[
\int d^4 x d^4 \theta U G
\]  
(29)

to the action (18). On the one hand side, varying eq. (26) with respect to \(G\) (in fact, with respect to a potential \(J_\alpha\) in the general solution \(G = D^\alpha \bar{D}^2 J_\alpha + \bar{D}_\alpha D^2 J_\bar{\alpha}\) to the defining constraints (26)), we get \(U = \frac{1}{2} (\Phi + \bar{\Phi})\) back. On the other hand side, varying with respect to \(U\) in the action

\[
S = \int d^4 x d^4 \theta \left[ - \ln U + U G + \frac{1}{32 \pi} U \Omega + \frac{1}{2} U W^2 \bar{W}^2 \psi(UK,UK) \right]
\]  
(30)

we find an algebraic equation on \(U:\)

\[
\frac{1}{U} = \left( G + \frac{1}{8 \pi} \Omega \right) + \frac{1}{32 \pi} W^2 \bar{W}^2 \left( 2U \psi(UK,UK) + U^2 \frac{\partial \psi(UK,UK)}{\partial U} \right) .
\]  
(31)

\(^{2}\)The possibility of such transformation was noticed in ref. [11].
Since $W_\alpha W_\beta W_\gamma = 0$ due to the anti-commutativity of $W_\alpha$, the second term on the right-hand-side of recursive relation (31) can be considered as an ‘exact’ perturbation. This leads to a complete solution to eq. (31) in the form

$$U^{-1} = G_{\text{mod}} + \frac{1}{32\pi} W^2 \bar{W}^2 \left( 2 \partial \gamma(G_{\text{mod}}^{-1} K, G_{\text{mod}}^{-1} \bar{K}) - \partial \gamma(G_{\text{mod}}^{-1} K, G_{\text{mod}}^{-1} \bar{K}) \right) ,$$

(32)

where we have introduced the ‘modified’ N=1 linear multiplet $G_{\text{mod}}$ as

$$G_{\text{mod}} = G + \frac{1}{8\pi} \Omega .$$

(33)

The appearance of the N=1 Chern-Simons superfield $\Omega$ is quite natural from the point of view of string theory and D-branes, where Chern-Simons-type couplings (in components) are known to appear in the famous Green-Schwarz anomaly cancellation mechanism and in the (dual) D-brane actions. In particular, the dilaton superfield $G$ must transform under the $U(1)$ gauge transformations as

$$\delta G = i \frac{1}{32\pi} (D^\alpha \Lambda) W_\alpha + \text{h.c.}$$

(34)

in order to make $G_{\text{mod}}$ gauge-invariant. Equations (26) and (28) lead to the manifestly gauge-invariant constraints on $G_{\text{mod}}$:  

$$\bar{D}^2 G_{\text{mod}} = \frac{1}{4\pi} W^2 , \quad D^2 G_{\text{mod}} = \frac{1}{4\pi} \bar{W}^2 .$$

(35)

Such couplings were extensively studied in superspace (see e.g., ref. [12] for a recent review), while the relevant superspace geometry appears to be closely related to a three-form N=1 multiplet introduced in ref. [13].

Substituting the solution (32) into the action (30) yields the dual action in the form

$$S = \int d^4 x d^4 \theta \left\{ \ln G_{\text{mod}} + \frac{1}{4\pi} W^2 W^2 G_{\text{mod}}^{-2} \gamma(G_{\text{mod}}^{-1} K, G_{\text{mod}}^{-1} \bar{K}) \right\} .$$

(36)

By construction this action is equivalent (dual) to the action given by a sum of eqs. (18) and (23). However, it seems to be much easier to find a supersymmetric completion of the action (36) with respect to the second (spontaneously broken) supersymmetry, since the action (36) is given by the full N=1 superspace integral, while the constraints (35) are also easy to be covariantized. The second non-linearly realized supersymmetry with the transformation law $\delta W_\alpha = \eta_\alpha + \ldots$ implies a non-trivial transformation law of $G_{\text{mod}}$ as well, because of the constraint (35),

$$\delta \eta G_{\text{mod}} = - \frac{1}{8\pi} (\eta^\alpha D_\alpha V + \bar{\eta}_\alpha \bar{D}^\alpha V) + \ldots .$$

(37)

The minimal, manifestly N=2 covariant version of the constraints (35) is given by

$$D^2 G_{\text{mod}} = \frac{1}{4\pi} \psi^2 , \quad \bar{D}^2 G_{\text{mod}} = \frac{1}{4\pi} \bar{\psi}^2 ,$$

(38)
where we have substituted the Maxwell-Goldstone N=1 superfield $W$ by the N=1 Goldstone superfield $\Psi$, and the N=1 linear (dilaton-axion) superfield $G_{\text{mod}}$ by its fully covariant counterpart $\tilde{G}_{\text{mod}}$. The superfield $W$ obeys the ‘canonical’ constraints (6) but it has the complicated transformation law (10), whereas the N=1 Goldstone superfield $\Psi$ has the ‘canonical’ transformation law under the second supersymmetry but it obeys the complicated constraints (12). The same remarks also apply to $G_{\text{mod}}$ and $\tilde{G}_{\text{mod}}$, respectively.

The defining constraints (38) on $\tilde{G}_{\text{mod}}$ are consistent with the constraints (12) because of the identities

$$D_\alpha D_\beta D_\gamma = D_\alpha D_\beta D_\gamma = 0$$

(39)

that follow from the definitions (13). The fully covariant action is thus of the form

$$S = \int d^4 x d^4 \theta E^{-1} \ln \tilde{G}_{\text{mod}},$$

(40)

where we have introduced a density $E^{-1}(\Psi, \bar{\Psi}, \tilde{G}_{\text{mod}})$ in the full N=1 superspace.

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