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4D SUPERGRAVITY COUPLED TO AN ANTISYMMETRIC TENSOR FIELD
WITH GREEN-SCHWARZ MECHANISM

S. Ferrara and M. Villasante^{*)}

CERN - Geneva

and

Physics Department
University of California, Los Angeles

ABSTRACT

We derive the component Lagrangian for general matter coupling to 4D supergravity of a real linear multiplet, containing an antisymmetric tensor field, which transforms non-trivially under gauge and Lorentz transformations, as it comes from the Green-Schwarz anomaly cancellation mechanism in $D = 10$. It is found that local supersymmetry requires a third Chern-Simons form related to the axial gauge field of conformal supergravity.

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It is well known that $N = 1$ Einstein supergravity coupled to chiral and vector multiplets is described by three basic functions: the Kähler function Φ , the superpotential g and the gauge kinetic function f_{AB} . Supersymmetry requires g and f_{AB} to be "holomorphic" functions of the chiral multiplets.

The standard supergravity action for chiral multiplets Z_i and vector multiplets V (V is a Lie-algebra-valued superfield) takes the form¹⁾

$$\int d^8\mu \Phi(z, \bar{z}e^V) + \text{Re} \int d^6\mu g(z_i) + \text{Re} \int d^6\mu f_{AB}(z) W^A W^B \quad (1)$$

where $d^8\mu$ and $d^6\mu$ denote full and chiral measures in superspace. The Kähler potential is given by

$$K = -3 \log \left(-\frac{\Phi}{3} \right) \quad (2)$$

and the Kähler invariant $\mathcal{G}$ function is

$$\mathcal{G} = K + \log \frac{|g|^2}{4} \quad (3)$$

If some chiral multiplet out of Z_i , generically denoted by S is a gauge singlet and moreover it appears in the function Φ through the combination $S + \bar{S}$ (this means that the first component of the Φ superfield depends only on $\text{Re } S$), then it can be transformed in a new multiplet, called "linear multiplet"²⁾, containing an antisymmetric tensor field $B_{\mu\nu}$. This generalizes to supersymmetric theories the well-known fact that, in $D = 4$, massless antisymmetric tensors are equivalent to massless scalar fields³⁾.

In order for this equivalence to work it is further required that the superpotential is independent of S (up to redefinition of the supervierbein which at most would allow an exponential dependence of g in terms of S). In addition, the f_{AB} function must be at most linear in S .

The linearity in S is allowed because it can be shown that

$$\text{Re} \int d^6\mu S W^2 = \int d^8\mu (S + \bar{S}) \Omega_{3YM}(V) \quad (4)$$

in which $\Omega_{3YM}(V)$ is a gauge-variant multiplet which generalizes the Yang-Mills

Chern-Simons form. In this case, the duality transformation in superspace language amounts to add to the previous Lagrangian the term

$$- \int d^8 \mu (S + \bar{S}) L \quad (5)$$

and to consider $S + \bar{S}$ as an independent (vector) superfield U . L is a "linear" multiplet, namely it satisfies the constraints

$$(\bar{D}\bar{D} - 8R) L = 0, \quad \bar{L} = L \quad (6)$$

Variation with respect to L implies

$$U = S + \bar{S} \quad \text{with} \quad \bar{D}S = 0 \quad (7)$$

On the other hand, variation with respect to U gives

$$\hat{L} = L - \Omega_{3YM} = \Phi_S \quad \left(\phi_s = \frac{\partial \phi}{\partial S} \right) \quad (8)$$

Equation (8) can be inverted to express the original Lagrangian given by Eq. (1) in terms of the L multiplet. An example of a supergravity Lagrangian with the previous terms is obtained by a dimensional reduction of $D = 10$ supergravity coupled to Yang-Mills gauge supermultiplets⁴⁾. Since such a theory may also be regarded as a point field limit of superstrings, this kind of Lagrangian will also appear in the low-energy effective action for more general superstring compactifications preserving $N = 1$ supersymmetry in $D = 4$.

As Green and Schwarz discovered⁵⁾, the anomaly cancellation in $D = 10$ requires a modification of the $B_{\mu\nu}$ field strength by an addition of a Chern-Simons Lorentz three form $\Omega_{3L}(\omega)$. The full $B_{\mu\nu}$ field strength is therefore given by

$$\hat{H} = dB + \frac{1}{30} \Omega_{3YM} - \Omega_{3L} \quad (9)$$

In $D = 4$ dimensions the Ω_{3L} term, in analogy with the corresponding Yang-Mills term, implies a coupling of the S multiplet to the square of the Weyl multiplet $W_{\alpha\beta\gamma}^2$. Unitarity⁶⁾ further requires that $W_{\alpha\beta\gamma}^2$ comes in the form of the super-Gauss-Bonnet combination of the supergravity curvatures⁷⁾:

$$W_{\alpha\beta\gamma} , E_{\alpha\dot{\alpha}} , R . \quad (10)$$

This is the result of Cecotti et al.⁸⁾. The F-type action containing the super-Gauss-Bonnet combination is

$$\text{Re} \int d^6\mu S \left[W_{\alpha\beta\gamma}^2 + \frac{1}{2} (\bar{S} S - 8R) (E_{\alpha\dot{\alpha}}^2 - 4R\bar{R}) \right] \quad (11)$$

It can be rewritten as a D-type action:

$$\int d^8\mu (S + \bar{S}) \left[\Omega_{3L} + E_{\alpha\dot{\alpha}}^2 - 4R\bar{R} \right] \quad (12)$$

provided

$$\text{Re} \int d^6\mu S W_{\alpha\beta\gamma}^2 = \int d^8\mu (S + \bar{S}) \Omega_{3L} \quad (13)$$

where Ω_{3L} is a gauge-variant expression. Although we do not have at present a superfield proof of Eq. (13), it can be shown that it is true in the component form in the Wess-Zumino gauge.

The superfield containing the Chern-Simons Lorentz term is given in Eq. (13) since, in components, the left-hand side of this equation contains the term

$$\text{Im} S \varepsilon^{\mu\nu\rho\sigma} R_{\mu\nu}{}^{ab} R_{\rho\sigma ab} \quad (14)$$

Eq. (12) also gives rise to a term containing the Euler density

$$\text{Re} S \varepsilon^{\mu\nu\rho\sigma} \varepsilon_{abcd} R_{\mu\nu}{}^{ab} R_{\rho\sigma}{}^{cd} \quad (15)$$

as shown in Ref. 8).

The purpose of this note is to write the general component form of the Lagrangian given in Eq. (1) but with the chiral multiplet S replaced by a linear multiplet L , as obtained by a duality transformation. It is the Lagrangian in terms of L , which among its components contains the $B_{\mu\nu}$ field, which is directly related to the low-energy limit of superstring theories. We will first derive the Lagrangian and transformation rules given by Eq. (1) in terms of the linear multiplet and later extend the results with the inclusion of the Chern-Simons Lorentz term.

The component form of the superspace Lagrangian given in Eq. (1) is given by:

$$e^{-1} \mathcal{L} = e^{-1} \mathcal{L}(\phi) + e^{-1} \mathcal{L}(g) + e^{-1} \mathcal{L}(f_{AB}) \quad (16)$$

according to the conventions of Cremmer et al.¹⁾

The conditions for the ϕ and f functions are now

$$\Phi(S, S^*, z_i, z^{*i}) = \Phi(\text{Re } S, z_i, z^{*i}) \quad (17)$$

$$f_{AB}(S, z_i) = a \delta_{AB} S + f_{AB}(z_i) \quad (a: \text{numerical constant}) \quad (18)$$

Under these assumptions, the field $\text{Im } S$ only appears through its derivative $\partial_\mu \text{Im } S$ in Eq. (16), up to integration by parts. The relevant part of the Lagrangian (16) is the one containing the S multiplet. One actually needs only the terms containing S and h_s^{*} . They are given by

$$\begin{aligned} e^{-1} \mathcal{L}_S = & -\frac{1}{2} \phi_{ss} \partial_\mu S \partial^\mu S^* - \frac{1}{2} \phi_s^i D_\mu z_i \partial^\mu S^* - \frac{1}{2} \phi_{is} \partial_\mu S D^\mu z^{*i} + \\ & + \frac{1}{2} \phi_{ss} |h_s|^2 + \frac{1}{2} \phi_s^i h_i h^{*s} + \frac{1}{2} \phi_{si} h_s h^{*i} - \phi_s^{I2} \bar{\chi}_{LI} \chi_{L2} h^{*s} + \\ & + \phi_s^{I2} \bar{\chi}_{LI} \not{\partial} S \chi_R^I + \frac{1}{3} \mu^* \phi_s h_s + \frac{i}{3} A^\mu \phi_s \partial_\mu S + \\ & + \phi_s^I \bar{\psi}_{\mu L} \not{\partial} S^* \gamma^\mu \chi_{LI} - \frac{1}{8} e^{-1} \varepsilon^{\mu\nu\rho\sigma} \bar{\psi}_\mu \gamma_\nu \psi_\rho \phi_s \partial_\sigma S + \\ & + \frac{a}{2} S \left(-\frac{1}{4} F_{\mu\nu}^A F^{\mu\nu}_A - \frac{1}{2} \bar{\lambda}^A \not{D} \lambda_A + \frac{1}{2} D^A D_A + \frac{1}{4} F_{\mu\nu}^A \bar{\psi}_\rho \sigma^{\mu\nu} \gamma^\rho \lambda_A \right) + \\ & - \frac{a}{8} e^{-1} \varepsilon^{\mu\nu\rho\sigma} \Omega_{\nu\rho\sigma}^{(YM)} \partial_\mu S + \frac{a}{4} \partial_\mu S \bar{\lambda}^A \gamma_s \gamma^\mu \lambda_A - \frac{a}{4} h_s \bar{\lambda}_L^A \lambda_{LA} + h.c. \end{aligned} \quad (19)$$

where $I = (i, s)$ and the i index does not include the S field, and $\Omega_{\mu\nu\rho}^{(YM)}$ is the Yang-Mills Chern-Simons form:

$$\Omega_{\mu\nu\rho}^{(YM)} = A_{[\mu}^A F_{\nu\rho]A} - \frac{1}{3} \tilde{g} C_{ABC} A_\mu^A A_\nu^B A_\rho^C \quad (20)$$

($\tilde{g}$ is the gauge coupling constant and C_{ABC} are the structure constants of the

*) The other terms in the Lagrangian (16) only contain $\text{Re } S$ implicitly through the ϕ function and its derivatives.

Yang-Mills group). Brackets enclosing a set of indices denote total antisymmetrization with strength one. The duality transformation is obtained by the following rules³⁾:

- a) elimination of the S auxiliary field h_s from the Lagrangian through its equations of motion;
- b) replace $\partial_\mu \text{Im } S$ by L_μ ;
- c) add to the Lagrangian a term with the Lagrange multiplier $B_{\mu\nu}$

$$\mathcal{L}_c = \frac{i}{4} \varepsilon^{\mu\nu\rho\sigma} B_{\mu\nu} \partial_\rho L_\sigma \quad (21)$$

Requirement of local supersymmetry of the Lagrangian

$$\mathcal{L} + \mathcal{L}_c \quad (22)$$

fixes the transformation rules of the $B_{\mu\nu}$ tensor provided L_μ transforms as $\partial_\mu \text{Im } S$.

By varying Eq. (22) with respect to L_μ , it allows to solve L_μ as a function of $B_{\mu\nu}$ and of the other fields. The $B_{\mu\nu}$ transformation rules are

$$\delta B_{\mu\nu} = -2 \left[\phi_s \bar{\psi}_{I\mu} \gamma_\nu \varepsilon + \bar{\varepsilon} \Gamma_{\mu\nu} (\phi_{sI} \chi_R^I + \phi_s^I \chi_{LI}) + \frac{a}{2} \bar{\varepsilon} \gamma_\mu \lambda^A A_{\nu A} \right] \quad (23)$$

(here $\Gamma_{\mu\nu} = \frac{i}{2} [\gamma_\mu, \gamma_\nu] = \Gamma_{[\mu} \Gamma_{\nu]} = 2\sigma_{\mu\nu}$).

The L_μ field equation is

$$\begin{aligned} L^\mu = \frac{1}{2\phi_{ss}} & \left[-i(\phi_{si} D^\mu \bar{z}^i - \phi_s^i D^\mu z_i) + 2i \phi_{s\bar{s}}^I \bar{\chi}_{LI} \gamma^\mu \chi_R^I - \frac{2}{3} \phi_s A^\mu + \right. \\ & \left. + i \bar{\psi}_\rho \gamma^\mu \gamma^\rho (\phi_{sI} \chi_R^I - \phi_s^I \chi_{LI}) + \frac{i}{4} e^{-1} \phi_s \varepsilon^{\mu\nu\rho\sigma} \bar{\psi}_\nu \gamma_\rho \psi_\sigma + e^{-1} H^\mu \right] \end{aligned} \quad (24)$$

where

$$\begin{aligned} H^\mu = \frac{i}{4} \varepsilon^{\mu\nu\rho\sigma} & \left[\partial_\nu B_{\rho\sigma} - a \Omega_{\nu\rho\sigma}^{(YM)} - \frac{a}{3!} \bar{\lambda}^A \Gamma_{\nu\rho\sigma} \lambda_A \right] \\ (\Gamma_{\nu\rho\sigma} = \Gamma_{[\nu} \Gamma_\rho \Gamma_{\sigma]}) \end{aligned} \quad (25)$$

The h_s field equation is

$$h_s = - \frac{A}{\phi_{ss}} \quad (26)$$

where

$$A = \phi_s^i h_i - \phi_s^{I\bar{J}} \bar{\chi}_{LI} \chi_{L\bar{J}} + \frac{1}{3} \mu \phi_s - \frac{a}{4} \bar{\lambda}_R^A \lambda_{RA} \quad (27)$$

The L_μ and h_s parts of the Lagrangian are:

$$\mathcal{L}(L_\mu) = e \phi_{ss} L_\mu L^\mu \quad (28)$$

$$\mathcal{L}(h_s) = -e \frac{|A|^2}{\phi_{ss}} \quad (29)$$

After using this, one finally gets:

$$\begin{aligned} e^{-1} (\mathcal{L}_s + \mathcal{L}_c) = & -\phi_{ss} (\partial_\mu \text{Re} S)^2 - \partial_\mu \text{Re} S (\phi_{is} D^\mu \bar{z}^i + \phi_s^i D^\mu z_i) + \\ & + a \text{Re} S \left(-\frac{1}{4} F_{\mu\nu}^A F_A^{\mu\nu} - \frac{1}{2} \bar{\lambda}^A \not{D} \lambda_A + \frac{1}{2} D^A D_A + \frac{1}{4} F_{\mu\nu}^A \bar{\psi}_\rho \sigma^{\mu\nu} \gamma^\rho \lambda_A \right) + \\ & + \bar{\psi}_\mu \not{D} \text{Re} S \gamma^\mu (\phi_{sI} \chi_R^I + \phi_s^I \chi_{LI}) + \\ & - \frac{1}{\phi_{ss}} \left| \phi_s^i h_i - \phi_s^{I\bar{J}} \bar{\chi}_{LI} \chi_{L\bar{J}} + \frac{1}{3} \mu \phi_s - \frac{a}{4} \bar{\lambda}_R^A \lambda_{RA} \right|^2 + \\ & + \frac{1}{4\phi_{ss}} \left[-i (\phi_{si} D^\mu \bar{z}^i - \phi_s^i D^\mu z_i) + 2i \phi_{s\bar{J}}^I \bar{\chi}_{LI} \gamma^\mu \chi_R^{\bar{J}} - \frac{2}{3} \phi_s A^\mu + \right. \\ & \left. + i \bar{\psi}_\rho \gamma^\mu \gamma^\rho (\phi_{sI} \chi_R^I - \phi_s^I \chi_{LI}) + \frac{i}{4} e^{-1} \phi_s \epsilon^{\mu\nu\rho\sigma} \bar{\psi}_\nu \gamma_\rho \psi_\sigma + e^{-1} H^\mu \right]^2 \end{aligned} \quad (30)$$

The other transformation rules are the same as before but we have to replace $\partial_\mu \text{Im} S$ by L_μ as given by Eq. (24) and h_s by its equation of motion (26).

The antisymmetric tensor field $B_{\mu\nu}$ has a scalar partner and a fermionic partner. The scalar partner is given by the lowest component of the superfield in Eq. (8):

$$C^L = \phi_s \quad (31)$$

The fermion partner is obtained from $\delta C^L = \bar{\epsilon} \chi^L$

$$\chi^L = \phi_s^I \chi_{LI} + \phi_{sI} \chi_R^I \quad (32)$$

and its transformation rule is

$$\begin{aligned} \delta \chi^L = & \frac{1}{4} \hat{D} C^L \varepsilon - \frac{1}{6} C^L (S_{mn} - i \gamma_5 P - i A \gamma_5) \varepsilon - \frac{i}{4} e^{-1} H^\mu \gamma_\mu \gamma_5 \varepsilon \\ & + \frac{1}{16} e^{-1} \varepsilon^{\mu\nu\rho\sigma} (C^L \bar{\psi}_\nu \gamma_\rho \psi_\sigma + 2 \bar{\psi}_\nu \Gamma_{\rho\sigma} \chi^L) \gamma_\mu \gamma_5 \varepsilon + \frac{a}{8} (\bar{\lambda}_R^\Lambda \lambda_{RA} \varepsilon_L + \bar{\lambda}_L^\Lambda \lambda_{LA} \varepsilon_R) \end{aligned} \quad (33)$$

We want to discuss now the inclusion of the Lorentz-Chern Simons term as given by Eq. (13) into the previous Lagrangian. For simplicity, we will confine our discussion here to the bosonic part of the Lagrangian. The full expression will be given elsewhere⁹⁾. The component expression of Eq. (13) is¹⁰⁾

$$\begin{aligned} e^{-1} \mathcal{L} = & b \operatorname{Re} S \left[W_{\mu\nu ab} W^{\mu\nu ab} - 6 R_{\mu\nu}(A) R^{\mu\nu}(A) \right] + \\ & + i b \operatorname{Im} S \left[-\frac{1}{2} \varepsilon^{abcd} R_{\mu\nu ab} R^{\mu\nu}_{cd} + 6 R_{\mu\nu}(A) \tilde{R}^{\mu\nu}(A) \right] \end{aligned} \quad (34)$$

$$\left(R_{\mu\nu}(A) = \frac{2}{3} (\partial_\mu A_\nu - \partial_\nu A_\mu) \right) \quad (b: \text{numerical constant})$$

The terms containing $\operatorname{Im} S$ amount to adding in Eq. (25) to the Yang-Mills Chern-Simons form the new Lorentz and A_μ Chern-Simons forms:

$$H_{\text{Bos}}^\mu = \frac{i}{4} \varepsilon^{\mu\nu\rho\sigma} \left[\partial_\nu B_{\rho\sigma} - a \Omega_{\nu\rho\sigma}^{(YM)} + 4b \Omega_{\nu\rho\sigma}^{\text{LOR}} - 24b \Omega_{\nu\rho\sigma}(A) \right] \quad (35)$$

where

$$\Omega_{\mu\nu\rho}^{\text{LOR}} = \omega_{[\mu}{}^{ab} R_{\nu\rho]ab} + \frac{2}{3} \omega_{\mu ab} \omega_{\nu}{}^{bc} \omega_{\rho c}{}^a \quad (36)$$

$$\Omega_{\mu\nu\rho}(A) = \frac{2}{3} A_{[\mu} R_{\nu\rho]}(A) \quad (37)$$

The expression for h_s remains the same when we neglect fermionic terms. If we consider the super Gauss-Bonnet completion of Eq. (34), there are some extra modifications⁸⁾. We get the extra terms in H^μ :

$$H_{\text{Bos}}^{\mu'} = -\frac{16ie}{9} b \left[u^* \overleftrightarrow{\partial} u - \frac{4}{3} i A^\mu (|u|^2 - A_\nu A^\nu) + 4i A_\nu \hat{B}^{\nu\mu} \right] \quad (38)$$

where $\hat{B}_{ab}$ has been given in Ref. 11)*). Incidentally we note that the $\hat{B}_{ab}$ tensor changes the coefficient of the A_μ Chern-Simons three-form in Eq. (35). Invariance of H^μ implies non-trivial gauge and Lorentz transformations for $B_{\mu\nu}$. The h_s equation of motion gets a modification which corresponds to adding to A in Eq. (27) the terms:

$$A' = \frac{16}{9} \ell \left[u \left(\frac{1}{2} R - \frac{1}{3} (|u|^2 + 5 A_\nu A^\nu) - i D_a A^a \right) - 2 i A^a \partial_a u \right] \quad (39)$$

Finally, the (bosonic) part of the Lagrangian containing $\text{Re } S$ [see Eq. (30)] gets the extra terms:

$$e^{-1} \mathcal{L}'_{\text{bos}} = \ell \text{Re } S \left[{}^*R^*_{abcd} R^{abcd} + \frac{32}{9} D_a D_b \left(\frac{1}{2} |u|^2 \eta^{ab} + A^a A^b \right) \right] \quad (40)$$

To summarize, we have shown that supersymmetry requires a new Chern-Simons term to be added to the definition of the supercovariant field strength H_μ of the $B_{\mu\nu}$ field. However, since A_μ does not fulfil an algebraic equation, this implies that in terms of physical fields, $H_{\mu\nu\rho}$ contains an infinite sequence of bosonic terms. This is the same conclusion reached in the dual formulation in Ref. 8). This result has a drawback in $D = 10$ dimension. The supercovariantized three-form $H_{\mu\nu\rho}$ cannot be given in a closed form in terms of the massless fields of $D = 10$ supergravity but rather it must be a solution of a self-consistency equation dictated by supersymmetry.

Different formulations of $D = 10$ supergravity with the inclusion of Chern-Simons Lorentz terms do indicate that this is indeed the case¹²⁾⁻¹⁵⁾. In addition, σ -model calculations in heterotic superstrings also suggest a non-trivial modification of the H field strength as dictated by a self-consistency condition equation¹⁶⁾.

*) Our $\hat{B}_{ab}$ was called B_{ab} in Ref. 11). Here we introduce a hat to avoid confusion with the antisymmetric tensor $B_{\mu\nu}$ which comes from the linear multiplet.

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