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AN OFF-SHELL SUPERSPACE R^2 -ACTION IN SIX DIMENSIONSE. BERGSHOEFF¹ and M. RAKOWSKI*International Centre for Theoretical Physics, I-34100 Trieste, Italy*

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The superspace constraints corresponding to the two inequivalent versions of the $40+40$ off-shell Weyl multiplet in six dimensions are given. A supersymmetric R^2 -action (including quartic fermion terms) for one of the two Weyl multiplets is constructed. A crucial role in this construction is played by the observation that there exists a formal relationship between the Weyl and Yang–Mills multiplets. Both the supersymmetric R^2 -action and the supersymmetric Yang–Mills F^2 -action coupled to the Weyl multiplet are reformulated as superspace actions.

1. Introduction. Over the last two years there has been a revival of interest in R^2 -supergravity theories¹¹ due to the fact that they arise naturally in the low energy limit of superstring theories [2], and moreover, they play an important role in their compactification [3]. There have been a few attempts towards the supersymmetrization of R^2 -actions in $d=10$ [4] and $d=6$ [5]. So far, these attempts have not led to a complete result. One of the difficulties involves the mixing of R and R^2 terms in the action, and the necessity of modifying the transformation rules, both of which are due to the fact that the theory is on-shell. In $d=4$ an off-shell formulation of $N=1$ supergravity is known. This has facilitated the fully supersymmetric extension of the ghost-free [6] Gauss–Bonnet combination $\phi(R^2_{\mu\nu\alpha\beta} - 4R^2_{\mu\alpha} + R^2)$ [7]. The only other supersymmetric R^2 -actions that we know of are the $d=4$, $N=1$ [8] and $N=2$ [9] superconformal actions which contain the (Weyl tensor)².

In this note, we give an example of a supersymmetric R^2 -action (including quartic fermion terms) in a higher dimension, by constructing the off-shell superconformal extension of (Riemann tensor)² in six dimensions. The action is given both in components as well as in superspace, thereby generaliz-

ing our previous result [10], where the same action was given up to quartics, and only in components. We have considered $d=6$ because (a) matter and Yang–Mills coupled to $N=2$, $d=6$ supergravity [11] is known off-shell [12]. Therefore, without changing the supersymmetry transformation rules, we can construct the R and R^2 -invariants separately, (b) the theory resembles the $N=1$, $d=10$ Yang–Mills coupled supergravity a great deal [13], and (c) the $N=2$, $d=6$ theory may itself be the low energy limit of a heterotic string theory [14].

Before constructing the supersymmetric R^2 -action, we will first give a superspace formulation of the $d=6$ conformal supergravity theory, whose component formulation has been given in ref. [12]. We will then proceed by first constructing the supersymmetric R^2 -action in components. A crucial role in this construction is played by the observation that there exists a formal relationship between the Weyl and Yang–Mills multiplets. This observation enables us to immediately derive a supersymmetric R^2 -action from a supersymmetric F^2 -action. We then formulate both the F^2 and R^2 component actions as superspace actions. Finally, we comment on the implication of these results for the construction of supersymmetric R^2 -actions in ten dimensions.

2. The off-shell supergeometry. A component formulation of the $N=2$, $d=6$ conformal supergravity theory has already been given [12] (see also ref.

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¹¹ For some recent discussions, see ref. [1].

[15]). Here, we use the same notation and conventions. The $d=6$ Weyl multiplet contains as components, a sechsbein e^a_μ , a gravitino ψ^i_μ , an $SU(2)$ gauge field V^{ij}_μ , plus additional matter fields. In ref. [12] it was shown that there exist two distinct sets of matter fields. One set consists of an antisymmetric tensor T_{abc} of negative duality, an $SU(2)$ -Majorana-Weyl spinor χ^i of positive chirality, and a real scalar D . The other set contains an antisymmetric tensor gauge field $B_{\mu\nu}$, an $SU(2)$ -Majorana-Weyl spinor λ^i of negative chirality, and a real scalar ϕ . The two sets of matter fields correspond to two inequivalent versions of the $40+40$ $N=2$, $d=6$ Weyl multiplet, which from now on are denoted as I and II, respectively:

$$\text{Formulation I: } e^a_\mu, \psi^i_\mu, V^{ij}_\mu, T_{abc}, \chi^i, D, \quad (1)$$

$$\text{Formulation II: } e^a_\mu, \psi^i_\mu, V^{ij}_\mu, B_{\mu\nu}, \lambda^i, \phi. \quad (2)$$

An unusual feature of formulation II, not shared by I is that the scalar ϕ and the spinor λ^i transform inhomogeneously under dilatations and conformal supersymmetry transformations. This enables one to impose the gauge $\phi=1$, $\lambda=0$ after which, one ends up with the field content given by $\{e^a_\mu, \psi^i_\mu, V^{ij}_\mu, B_{\mu\nu}\}$.

It is not difficult to derive the superspace constraints which yield the component results given in (1) and (2). In fact, the constraints leading to formulation I have already been given in ref. [16]. It is convenient to formulate these constraints in an $SU(2)$ superspace, i.e. a superspace whose tangent space group not only contains Lorentz rotations, but also $SU(2)$ gauge transformations. The basic superfield which contains all the component fields given in (1), is an antiselfdual tensor superfield T_{abc} whose $\theta=0$ component can be identified with the component field T_{abc} . This is very similar to the situation in four dimensions, where the $N=2$ conformal supergravity theory is described by a chiral antiselfdual superfield T_{ab} [17]. The basic torsion constraints describing formulation I are given by [16]¹²

$$T^a_{\alpha i, \beta j} = \gamma^a_{\alpha\beta} \epsilon_{ij}, \quad T^a_{bc} = T^{a-}_{bc}, \quad (3)$$

$$T^a_{\beta i, c} = T^{\alpha i}_{\beta j, \gamma k} = 0,$$

where $\alpha=1, \dots, 8$ is a spinor index, $i=1, 2$ an $SU(2)$ index, $a=1, \dots, 6$ a Lorentz vector index, and T_{abc} is the basic superfield of formulation I.

Substitution of the constraints (3) into the super Bianchi identities leads to a further restriction on the torsion component $T^{\alpha i}_{b, \gamma j}$. The only nonzero component of $T^{\alpha i}_{b, \gamma j}$ is given by an $SU(2)$ vector superfield D^i_a that transforms inhomogeneously under super scale transformations. In order to see how the components given in (1) fit into the torsion components one should go to a Wess-Zumino gauge for the super scale transformations [17]. This is all very similar to what happens for the $N=2$ conformal supergravity theory in four dimensions.

The superspace constraints leading to formulation II have not been worked out before. They resemble the constraints describing the conformal supergravity theory in ten dimensions [19] and also the constraints of the on-shell Poincaré theory in six dimensions [20]. The fundamental superfield containing all the components of formulation II in the gauge $\phi=1$, $\lambda=0$ is a three-index antisymmetric tensor superfield H_{abc} that satisfies a Bianchi identity $dH=0$, and whose $\theta=0$ component can be identified with the field strength of the gauge field $B_{\mu\nu}$. The basic set of torsion constraints is given by¹³ [21]

$$T^a_{\alpha i, \beta j} = \gamma^a_{\alpha\beta} \epsilon_{ij}, \quad T^a_{bc} = H^a_{bc},$$

$$T^a_{\beta i, c} = T^{\alpha i}_{\beta j, \gamma k} = 0. \quad (4)$$

A remarkable feature is that after substitution of the constraints (4) into the super Bianchi identities, one still has the freedom to impose the additional constraint

$$T^{\alpha i}_{b, \gamma j} = 0. \quad (5)$$

Note that the set of constraints (4), (5) is no longer invariant under superscale transformations. The torsion component $T^{\alpha i}_{bc}$ is now equal to the gravitino curvature ψ_{ab} . We have checked that the constraints (4), (5) exactly lead to the component result given in ref. [12] in the gauge $\phi=1$, $\lambda=0$.

An interesting feature of formulation II is that it allows for the existence of a closed super three-form $\mathcal{H}_{ABC}$ with nonzero components given by

¹² For our superspace notation and conventions see ref. [18]; however, we have labelled the indices as in ref. [12].

¹³ The same set of constraints has been derived independently by J. Gates Jr., H. Nishino, and E. Sezgin [21].

$$\mathcal{H}_{abc} = H_{abc}, \quad \mathcal{H}_{a,\alpha i, \beta j} = \epsilon_{ij}(\gamma_a)_{\alpha\beta}. \quad (6)$$

This concludes our description of the off-shell supergeometry.

3. The R^2 -action. We first construct the supersymmetric (Riemann tensor)² action in components. As explained in ref. [10] such an action exists only for formulation II of the Weyl multiplet. Therefore, from now on we will no longer consider formulation I. The component supersymmetry transformations that follow from the superspace constraints (4), (5) agree with the ones given in refs. [10,12]. Using the same notation and conventions as in refs. [10,12], the basic transformation rules in the gauge $\phi=1$, $\lambda=0$ are given by

$$\begin{aligned} \delta \omega_{\mu}{}^{ab}{}_{-} &= -\frac{1}{2} \bar{\epsilon} \gamma_{\mu} \psi^{ab}, \\ \delta \psi^{abi} &= \frac{1}{4} \gamma^{cd} \epsilon^i \hat{R}_{cd}{}^{ab}(\omega_{-}) + \frac{1}{2} \hat{V}^{abij} \epsilon_j, \\ \delta \hat{V}^{abij} &= \epsilon^{(i} \gamma^{\mu} \hat{\mathcal{D}}_{\mu}(\omega, \omega_{-}) \psi^{abj)} \\ &\quad - \frac{1}{24} \epsilon^{(i} \gamma^{def} \psi^{abj)} H_{def}, \end{aligned} \quad (7)$$

where

$$\omega_{\mu ab \pm} = \omega_{\mu ab}(e, \psi) \pm \frac{1}{2} H_{\mu ab},$$

$$\psi_{\mu\nu} = \mathcal{D}_{\mu}(\omega_{+}) \psi_{\nu} - (\mu \leftrightarrow \nu),$$

$$\begin{aligned} \hat{\mathcal{D}}_{\mu}(\omega, \omega_{-}) \psi^{abi} &= \partial_{\mu} \psi^{abi} + \frac{1}{4} \omega_{\mu}{}^{cd} \gamma_{cd} \psi^{abi} + 2 \omega_{\mu}{}^{[ac} \psi^{cb]i} - \frac{1}{2} V_{\mu j}^i \psi^{abj} \\ &\quad - \frac{1}{4} \gamma^{cd} \psi_{\mu}^i \hat{R}_{cd}{}^{ab}(\omega_{-}) - \frac{1}{2} \hat{V}^{abij} \psi_{\mu j}, \end{aligned} \quad (8)$$

and $\hat{R}_{\mu\nu}{}^{ab}(\omega_{-})$, $\hat{V}_{\mu\nu}^{ij}$, $H_{\mu\nu\rho}$ are the supercovariant curvatures of $\omega_{\mu}{}^{ab}{}_{-}$, V_{μ}^{ij} , and $B_{\mu\nu}$ respectively [10,12].

In order to simplify the construction of the R^2 -action it is now crucial to make the following observation. Consider an 8+8 off-shell Yang-Mills multiplet which contains a vector A_{μ}^I (I is the Yang-Mills index), a spinor χ^I of positive chirality and an SU(2) triplet auxiliary field Y^{ij} . The transformation rules of these Yang-Mills fields coupled to formulation II in the gauge $\phi=1$, $\lambda=0$ are given by [12]

$$\delta A_{\mu}^I = -\bar{\epsilon} \gamma_{\mu} \chi^I, \quad (9)$$

$$\delta \chi^I = \frac{1}{8} \gamma^{cd} \epsilon^i \hat{F}_{cd}^I - \frac{1}{2} Y^{lij} \epsilon_j,$$

$$\delta Y^{lij} = -\bar{\epsilon}^{(i} \gamma^{\mu} \hat{\mathcal{D}}_{\mu}(\omega, A) \chi^{lj)} + \frac{1}{24} \bar{\epsilon}^{(i} \gamma^{def} \chi^{lj)} H_{def}, \quad (9, \text{cont'd})$$

where

$$\begin{aligned} \hat{\mathcal{D}}_{\mu}(\omega, A) \chi^{li} &= \mathcal{D}_{\mu}(A) \chi^{li} + \frac{1}{4} \omega_{\mu}{}^{cd} \gamma_{cd} \chi^{li} - \frac{1}{2} V_{\mu j}^i \chi^{lj} \\ &\quad - \frac{1}{8} \gamma^{cd} \psi_{\mu}^i \hat{F}_{cd}^I + \frac{1}{2} Y^{lij} \psi_{\mu j}, \end{aligned} \quad (10)$$

and $\hat{F}_{\mu\nu}^I$ is the supercovariant field strength of A_{μ}^I [12]. We now note that there exists a formal relationship between the transformation rules (7) and (9). More precisely, the transformation rules (7) become identical to (9) upon the following substitution:

$$\begin{aligned} \omega_{\mu}{}^{ab}{}_{-} &\mapsto A_{\mu}^I \quad \text{or} \quad \hat{R}_{\mu\nu}^{ab}(\omega_{-}) \mapsto \hat{F}_{\mu\nu}^I, \\ \psi^{ab} &\mapsto 2\chi^I, \\ \hat{V}^{abij} &\mapsto 2Y^{lij}. \end{aligned} \quad (11)$$

In other words the conformal supergravity fields $\{\omega_{\mu}{}^{ab}{}_{-}, \psi^{ab}, \hat{V}^{abij}\}$ have the same transformation rules as an SO(6) vector multiplet $\{A_{\mu}^I, \chi^I, Y^{ij}\}$. Of course the fields of the conformal multiplet satisfy additional Bianchi identities such that they describe less components than are contained in a SO(6) multiplet.

Using the above observation it is now difficult to construct a supersymmetric R^2 -action. We first note that a supersymmetric Yang-Mills F^2 action has already been constructed [12]. In the gauge $\phi=1$, $\lambda=0$ the lagrangian is given by

$$\begin{aligned} e^{-1} \mathcal{L}(F^2) &= -\frac{1}{4} F_{\mu\nu}^I F_{\mu\nu}^I - 2 \bar{\chi}^I \gamma^{\lambda} \mathcal{D}_{\lambda}(\omega, A) \chi^I + Y^{ij} Y_{ij}^I \\ &\quad + \frac{1}{4} (F_{\mu\nu}^I + \hat{F}_{\mu\nu}^I) \bar{\psi}_{\lambda} \gamma^{\mu\nu} \gamma^{\lambda} \chi^I + \frac{1}{2} \bar{\chi}^I \gamma^{\lambda} H_{\lambda}^I \\ &\quad + \frac{1}{16} i e^{-1} \epsilon^{\mu\nu\rho\sigma\lambda\tau} F_{\mu\nu}^I F_{\rho\sigma}^I B_{\lambda\tau}. \end{aligned} \quad (12)$$

The idea is now to use (12) as a more general action formula: Any multiplet which transforms like a Yang-Mills multiplet has an invariant action given by (12)! In other words, we can now immediately use (12) to construct a supersymmetric R^2 -action. Upon making the formal substitution (11), one obtains the following result:

$$\begin{aligned}
e^{-1} \mathcal{L}(R^2) = & R_{\mu\nu}^{ab}(\omega_-) R_{\mu\nu}^{ab}(\omega_-) \\
& + 2\bar{\psi}^{ab} \gamma^\lambda \mathcal{D}_\lambda(\omega, \omega_-) \psi^{ab} - \hat{V}_{ab}^{ij} \hat{V}_{ij}^{ab} \\
& - \frac{1}{2} [R_{\mu\nu}^{ab}(\omega_-) + \hat{R}_{\mu\nu}^{ab}(\omega_-)] \bar{\psi}_\lambda \gamma^{\mu\nu} \gamma^\lambda \psi^{ab} \\
& - \frac{1}{12} \bar{\psi}^{ab} \gamma \cdot H \psi^{ab} \\
& - \frac{1}{4} i e^{-1} \epsilon^{\mu\nu\rho\sigma\lambda\tau} R_{\mu\nu}^{ab}(\omega_-) R_{\rho\sigma}^{ab}(\omega_-) B_{\lambda\tau}. \quad (13)
\end{aligned}$$

We have checked that the up to quartic terms in (13) exactly reproduce the result of ref. [10].

It is now not difficult to reformulate the component F^2 and R^2 actions as superspace actions. A general method for constructing superspace actions in six dimensions is not well understood. The only justification for their explicit expressions is that they reproduce the component results. We have checked that the following superspace action reproduces the component F^2 lagrangian given in (12):

$$S(F^2) = \int d^6 x \mathcal{D}_\alpha^{(i} \mathcal{D}_\beta^{j)} M_{\gamma\delta}^{\alpha\beta} (\chi_i^\gamma \chi_j^{\delta I}), \quad (14)$$

where χ_α^I is the Yang-Mills spinor superfield whose $\theta=0$ component is identified with the Yang-Mills spinor χ^I . Our normalization is such that the supersymmetry variation of χ^I (see (9)) is given by the relation $\delta\chi^I = (1/\sqrt{2}) \epsilon^{\alpha i} \mathcal{D}_\alpha \chi^I|_{\theta=0}$. The relevant components of the measure function $M_{\gamma\delta}^{\alpha\beta} = -M_{\gamma\delta}^{\beta\alpha} = -M_{\delta\gamma}^{\alpha\beta}$ are given by

$$\begin{aligned}
M_{\gamma\delta}^{\alpha\beta}|_{\theta=0} = & \frac{1}{12} e [\delta_{[\gamma}^{\alpha} \delta_{\delta]}^{\beta]} - \frac{1}{2} B_{\mu\nu} (\gamma^{\mu\nu})^{\alpha}{}_{[\gamma} \delta_{\delta]}^{\beta]}, \\
\mathcal{D}_\alpha^i M_{\gamma\delta}^{\alpha\beta}|_{\theta=0} = & (e/24\sqrt{2}) (\gamma^\mu)_{\gamma\delta} (\psi_\mu^{\beta i} + (\gamma_\mu{}^\nu \psi_\nu^i)^\beta \\
& + B_{\mu\nu} (\gamma^{\nu\rho} \psi_\rho^i)^\beta + \frac{1}{2} B_{\nu\rho} (\gamma^{\nu\rho} \psi_\mu^i)^\beta], \\
\mathcal{D}_\alpha^{(i} \mathcal{D}_\beta^{j)} M_{\gamma\delta}^{\alpha\beta}|_{\theta=0} = & -\frac{1}{16} e (\gamma^\rho)_{\gamma\delta} \{\bar{\psi}_\mu^{(i} \gamma^{\mu\nu}{}_\rho \psi_\nu^{j)} \\
& + \frac{1}{8} B^{\mu\nu} \bar{\psi}^{\lambda(i} \gamma_{[\mu\nu\lambda} \psi^{\tau j)} g_{\tau\rho}\}. \quad (15)
\end{aligned}$$

The action (14) generalizes the flat superspace action of ref. [22] to curved superspace. Note that (14) and (15) are just the analogue of the chiral density formula in four dimensions [23].

Using (14) we can now immediately derive the superspace action that reproduces the component R^2 lagrangian given in (13). We just replace in (14) the Yang-Mills spinor superfield χ_α^I by the gravitino superfield ψ_α^{iab} . We thus obtain

$$S(R^2) = \int d^6 x \mathcal{D}_\alpha^{(i} \mathcal{D}_\beta^{j)} M_{\gamma\delta}^{\alpha\beta} (\psi_i^{\gamma ab} \psi_j^{\delta ab}). \quad (16)$$

4. Comments. In this note we have given the fully supersymmetrization (including quartic fermions) of (Riemann tensor)² in six dimensions both in components and in superspace. The construction was enormously simplified by the observation that there exists a formal relationship between supergravity and Yang-Mills in six dimensions. We have also given the superspace constraints describing the two inequivalent $d=6$ Weyl multiplets. Concerning these constraints, it is interesting to see that the constraints of formulation II are consistent with the coupling of off-shell conformal supergravity to the Green-Schwarz superstring [24].

Although supersymmetric Yang-Mills F^4 actions in six dimensions have been recently constructed (up to order α^2) [25], it is not obvious how to use these results to construct supersymmetric R^4 actions. The results of ref. [24] suggest that the transformation rules of the off-shell Weyl multiplet will receive order α^2 corrections.

A related question to ask is whether there also exists a formal relationship between supergravity and Yang-Mills in ten dimensions. Such a relationship could then be used to simplify the construction of $d=10$ supersymmetric R^2 actions. Unfortunately, in ten dimensions things are more complicated. First of all, there does not exist an off-shell $d=10$ Weyl multiplet, and one cannot compare the off-shell $d=10$ Yang-Mills multiplet. Of course one could try to compare the transformation rules of the on-shell supergravity fields $\{\omega_\mu{}^{ab}, \psi^{ab}\}$ with those of on-shell Yang-Mills $\{A_\mu^I, \chi^I\}$. It turns out that after making suitable field redefinitions, the transformation rules are very similar, but they are not quite the same. The origin of this difference is that the supersymmetry transformation of the on-shell supergravity field $B_{\mu\nu}$ contains a term proportional to δA_μ^I :

$$\delta B_{\mu\nu} = A_\mu^I \delta A_\nu^I - (\mu \leftrightarrow \nu). \quad (17)$$

This term is related to the presence of a Yang-Mills Chern-Simons term in the theory [13]. One could now still try to construct out of a Yang-Mills action, a $d=10$ R^2 action which is supersymmetric under all

variations except the one given in (17). One should then try to restore the supersymmetry invariance (at least up to order α^2) by adding additional terms to the action. We hope to comment on this program in a future publication.

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References

- [1] S. Deser, Energy in gravity theories with quadratic curvatures, Brandeis University preprint BRXTH-195 (1985); D.G. Boulware and S. Deser, String generated gravity models, University of Washington preprint 40048-27 P5 (1985).
- [2] J. Scherk and J.H. Schwarz, Nucl. Phys. B 81 (1974) 118.
- [3] P. Candelas, G.T. Horowitz, A. Strominger and E. Witten, Nucl. Phys. B 258 (1985) 46.
- [4] L.J. Romans and N.P. Warner, Nucl. Phys. B 273 (1986) 320;
S.K. Han, J.K. Kim, I.G. Koh and Y. Tanii, Phys. Rev. D 34 (1986) 553;
S.J. Gates Jr. and H. Nishino, Manifestly supersymmetric $O(\alpha')$ superstring corrections in new $D=10$, $N=1$ supergravity Yang-Mills theory, University of Maryland preprint (1985).
- [5] S.K. Han, J.K. Kim, I.G. Koh and Y. Tanii, Phys. Lett. B 117 (1986) 167;
S.J. Gates Jr. and H. Nishino, Manifestly supersymmetric extensions of (curvature)²-terms in six dimensional $N=2$ supergravity, University of Maryland preprint 86-137 (1986).
- [6] B. Zwiebach, Phys. Lett. B 156 (1985) 315.
- [7] S. Cecotti, S. Ferrara, L. Girardello and M. Porrati, Phys. Lett. B 164 (1985) 46.
- [8] M. Kaku, P.K. Townsend and P. van Nieuwenhuizen, Phys. Rev. D 17 (1978) 3179;
M. Kaku, P. Townsend, P. van Nieuwenhuizen and S. Ferrara, Nucl. Phys. B 129 (1977) 125;
P.K. Townsend and P. van Nieuwenhuizen, Phys. Rev. D 19 (1979) 3166.
- [9] E. Bergshoeff, M. de Roo and B. de Wit, Nucl. Phys. B 182 (1981) 173.
- [10] E. Bergshoeff, A. Salam and E. Sezgin, Phys. Lett. B 173 (1986) 73; Nucl. Phys. B 279 (1987) 659.
- [11] H. Nishino and E. Sezgin, Phys. Lett. B 144 (1984) 187; Nucl. Phys. B 278 (1986) 353.
- [12] E. Bergshoeff, E. Sezgin and A. Van Proeyen, Nucl. Phys. B 264 (1986) 653.
- [13] A.H. Chamseddine, Nucl. Phys. B 185 (1981) 403;
E. Bergshoeff, M. de Roo, B. de Wit and P. van Nieuwenhuizen, Nucl. Phys. B 195 (1982) 97;
G.F. Chapline and N.S. Manton, Phys. Lett. B 120 (1983) 105.
- [14] K.S. Narain, Phys. Lett. B 169 (1986) 41.
- [15] E. Bergshoeff, Proc. 1st Torino Meeting on Superunification and extra dimensions, eds. R. D'Auria and P. Fré (1985).
- [16] P. Breitenlohner and A. Kabelschacht, Nucl. Phys. B 148 (1979) 96;
S.J. Gates Jr., Phys. Lett. B 84 (1979) 205; Nucl. Phys. B 162 (1979) 79.
- [17] See e.g. P.S. Howe, Nucl. Phys. B 199 (1982) 309.
- [18] S.W. MacDowell and M.T. Rakowski, Nucl. Phys. B 274 (1986) 589.
- [19] B.E.W. Nilsson and A. Tollstén, Phys. Lett. B 169 (1986) 369;
P.S. Howe and A. Umerski, Phys. Lett. B 177 (1986) 163.
- [20] M. Awada, G. Sierra and P.K. Townsend, Class. Quant. Grav. 2 (1985) L85.
- [21] E. Sezgin, private communication.
- [22] L.N. de Mator Pimentao, Diplom Arbeit, Institut für Theoretische Physik, Universität Karlsruhe (1980);
B.E.W. Nilsson, Nucl. Phys. B 174 (1980) 335;
J. Wess, Proc. 5th John Hopkins Workshop on Current problems in particle physics (Baltimore, 1981);
P.S. Howe, G. Sierra and P.K. Townsend, Nucl. Phys. B 221 (1983) 337.
- [23] See e.g. S.J. Gates Jr., M.T. Grisaru, M. Roček and W. Siegel, Superspace (Benjamin-Cummings, Menlo Park, CA, 1983).
- [24] E. Witten, Nucl. Phys. B 266 (1986) 245;
E. Bergshoeff, E. Sezgin and P.K. Townsend, Phys. Lett. B 169 (1986) 191.
- [25] E. Bergshoeff, M. Rakowski and E. Sezgin, to be published in Phys. Lett. B.
- [26] E. Bergshoeff, M. de Roo and B. de Wit, Nucl. Phys. B 217 (1983) 489.