Universal scheme of minimal reduction of usual and dual N=1,D=10 supergravity to the Minkowsky space

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Abstract

The reduction from N=1, D=10 to N=4, D=4 supergravity with the Yang-Mills matter is considered. For this purpose a set of constraints is imposed in order to exclude six additional abelian matter multiplets and conserve the supersymmetry. We consider both the cases of usual and dual N=1, D=10 supergravity using the superspace approach. Also the effective potential of the deriving theory is written.

1 Introduction

The action of the N=1, D=10 supergravity is supposed to be the effective low-energy limit of the superstring action where all the massive degrees of freedom are integrated out. All the terms of this action may be characterized by the number

\[ n = N_\partial + \frac{1}{2} N_f \]

where \( N_\partial \) and \( N_f \) are numbers of derivatives and fermions. The minimal supergravity has only the terms with \( n = 2 \). Usual and dual supergravity are equivalent each other at the minimal level. But requirement of anomalies cancellation implies that the Chern-Simons term must be presented in the field-strength \( H \) of the usual supergravity. It leads to appearance of terms with higher \( n \) in the lagrangian. Here the dual supergravity becomes more preferable because there is a good reason to believe that only the terms with \( n = 4 \) must be added to the dual supergravity lagrangian while the usual supergravity lagrangian turns out to be the infinite series in \( n \). We are not able to take into account terms with \( n > 2 \) now because they have not yet written explicitely but this work is on the way.

We suppose that the ten-dimensional space-time \( M^{10} \) is the product \( M^4 \times K \) of the Minkowsky space \( M^4 \) and some internal manifold \( K \) with unknown structure. It means that every vector \( V^M \) in \( M^{10} \) decomposes into vector \( V^\mu + 6 \) scalars \( V^m \) in \( M^4 \) (and vice versa in \( K \)) and every Majorana-Weyl spinor in \( M^{10} \) decomposes into 4 Majorana spinors in \( M^4 \). So

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one can write down how the N=1,D=10 supergravity degrees of freedom disintegrate under
the reduction $D = 10 \rightarrow D = 4$ (in brackets the fields carried them are denoted):

| N = 1 | D = 10 | $\rightarrow$ | D = 4 |
|-------|--------|---------------|------|
| **Fermions:** | | | |
| gravitino | $56(\psi_M)$ | $= 4 \times 2(\psi_\mu) + 24 \times 2(\psi_m)$ |
| dilatino | $8(\chi)$ | $= 4 \times 2(\chi)$ |
| **Bosons:** | | | |
| graviton | $35(E_M^A)$ | $= \frac{2(e_\mu^a) + 6 \times 2(B^m_{\mu}) + 21(e_m^a)}{1(\phi)}$ |
| dilaton | $1(\phi)$ | $= 1(\phi)$ |
| 2-potential (usual) | $28(B_{MN})$ | $= \frac{1(B_{\mu\nu}) + 6 \times 2(B_{m\mu}) + 15(B_{mn})}{1(B_{\mu\nu})}$ |
| 6-potential (dual) | $28(M_{M_1...M_6})$ | $= \frac{1(M_{m_1...m_6}) + 6 \times 2(M_{\mu m_1...m_5}) + 15(M_{\mu m_1...m_4})}{1(M_{m_1...m_6})}$ |

(for notations see below). It is easy to see that the underlined fields constitute the multiplet of
the N=4, D=4 supergravity. All other fields are put together into 6 abelian matter
multiplets. Our aim here is to eliminate the abelian multiplets in order to obtain the N=4,
D=4 supergravity as a part of the N=1, D=10 supergravity where 6 coordinates $y^m$ are
compactified on $K$. We don’t know whether these multiplets are essential in the low-energy
limit; but, definitely, the features of the theory become more simple to analysis without them
and they always can be taken into account as a perturbation to the N=4, D=4 supergravity.

In usual N=1, D=10 supergravity without additional Yang-Mills (YM) matter the separa-
tion of the abelian multiplets from pure N=4,D=4 supergravity has been realized in [6]
and some attempts to analise the dual case were made in [7]. In the usual case with YM-
matter the abelian multiplets have been eliminated in [8] by means of some constraints. In
this paper we find the constraints applying to both the cases of usual and dual supergravity
coupled with the YM-matter. It is demonstrated also that these constraints are unique ones.

Unfortunately the minimal N=4, D=4 supergravity has some problems which make it
difficult to obtain a realistic model. One of them is the cosmological term which appears in
the lagrangian where the internal $SU(2) \times SU(2) \sim O(4)$ symmetry is gauged [9]. From the
ten-dimensional point of view this gauging corresponds to the $y$-dependent compactification
by Scherk-Schwarz [10]. Nonminimal terms can cancel the cosmological term in the action
and it is one of the reasons why they could be important. We hope that the scheme of
reduction described here will be the most convenient one in the case of nonminimal N=1,
D=10 supergravity too. It is a matter for future speculations.

In section 2 we fix the notations; in section 3 the constraints are derived; in section 4 the
effective potential of the N=4,D=4 supergravity is written.
2 Notations

The following index notations are used here:

| dimension | flat          | world        |
|-----------|---------------|--------------|
| $D = 10$  | $A, B, C, \ldots$ | $M, N, P, \ldots$ |
| $D = 4$   | $\alpha, \beta, \gamma, \ldots$ | $\mu, \nu, \lambda, \ldots$ |
| $D = 6$   | $a, b, c, \ldots$ | $m, n, p, \ldots$ |

A number of formulae are taken from the superspace approach where the 16–component representation for the Majorana-Weil spinors is the most convenient one. We are not interesting in decomposition them into four 4–component spinors under the reduction $D = 10 \rightarrow D = 4$ here (this procedure has been described in many papers). So we use the $16 \times 16$ $\Gamma$–matrices with upper and lower indices. The spinorial indices will be omitted usually.

The fields of pure N=1,D=10 supergravity are given in the introduction; the fields of the YM-multiplet are: $A_M$ – the gauge potential and $\lambda$ – the gaugino field. They take values in the Lie algebra of the gauge group G:

$$A_M = i A^i_M t^i, \quad \lambda = i \lambda^i t^i,$$

where $t^i$ – the hermitian G-group generators.

The superspace description of the N=1, D=10 supergravity is the most convenient one, especially in the nonminimal case. The superspace (10 ordinary coordinates + 16 spinoral coordinates) has a nonzero torsion $\hat{T}^\hat{A}\hat{B}\hat{C}$, where $\hat{A}$, $\hat{B}$, $\hat{C}$ take vector or spinoral values. Due to some set of constraints, defining the field parametrization, and Bianchi identities the components of the superspace torsion and curvature are expressed through the fields of the supergravity multiplet. The field parametrization used here has been introduced in [11] and slightly modified in [5, 12]. This is not the parametrization with canonical kinetic terms in the lagrangian [2, 3] but it is sufficiently convenient one from the superspace point of view. Nevertheless, the connection with all other parametrizations can be restored unambiguously if the supersymmetry transformations are given. In our case they have the form [12]:

$$\epsilon_M^A = \psi_M \Gamma^A \epsilon$$

$$\delta \psi_M = \epsilon;_M - \frac{1}{144} \left( 3 \hat{T} \Gamma_M + \Gamma_M \hat{T} \right) \epsilon - \frac{1}{4} \Gamma^{PQ} \epsilon S_{MPQ}$$

$$\delta \phi = - \chi \epsilon$$

$$\delta \chi = - \frac{1}{2} \phi;_A \Gamma^A \epsilon + \frac{1}{36} \phi \hat{T} \epsilon + \frac{1}{2} \Gamma^A \epsilon \left( \psi_A \chi \right) + \frac{1}{2} \text{Sp}[\Gamma_A \lambda (\lambda \Gamma^A \epsilon)]$$

$$\delta M_{1, \ldots, 6} = - 3 \psi_{[M_1 \Gamma_{M_2 \ldots M_6}] \epsilon} \quad \text{(dual)}$$

$$\delta B_{MN} = \phi \psi_{[M} \Gamma_{N]} \epsilon - \frac{1}{2} \lambda \Gamma_{MN} \epsilon + \frac{1}{\sqrt{2}} \text{Sp}(A_{[M} \lambda \Gamma_{N]} \epsilon) \quad \text{(usual)}$$

$$\delta \lambda = - \frac{1}{2 \sqrt{2}} \hat{F} \epsilon$$

$$\delta A_M = \frac{1}{\sqrt{2}} \lambda \Gamma_M \epsilon \quad \text{(1)}$$
where the semicolon denotes the ordinary covariant derivative (without torsion);

\[ S_{ABC} = \frac{1}{2} (2 \psi_A \Gamma_{[B} \psi_{C]} + \psi_B \Gamma_A \psi_C) ; \]  
\(^{(2)}\)

\[ \hat{F} = \mathcal{F}_{AB} \Gamma^{AB} , \mathcal{F}_{AB} \text{ is the superspace matter field-strength:} \]

\[ \mathcal{F}_{AB} = F_{AB} + \sqrt{2} \psi_{[A} \Gamma_{B]} \lambda , \]  
\(^{(3)}\)

\( F_{AB} \) is the ordinary matter field-strength:

\[ F_{AB} = 2 A_{[B,A]} - 2 g A_{[AAB]} , \]  
\(^{(4)}\)
g is the charge; \( \hat{T} = T_{ABC} \Gamma^{ABC} \), where \( T_{ABC} \) is the superspace torsion with three vector indeces. It is expressed in a different way in the cases of usual and dual supergravity:

\[ \phi T_{ABC} = - 2 H_{ABC} + 3 \phi \psi_{[A} \Gamma_{B]} \psi_{C]} - 3 \psi_{[A} \Gamma_{BC]} \chi + \frac{1}{2} \text{Sp}(\lambda \Gamma_{ABC} \lambda) \]  
\(^{(5)}\)
in usual case and

\[ T_{ABC} = 2 M_{ABC} - \frac{1}{2} \psi^D \Gamma_{DABCE} \psi^E \]  
\(^{(6)}\)
in dual. Here

\[ H_{ABC} = 3 B_{[AB,C]} - 6 \text{Sp}(A_{[AAB,C]} + \frac{2}{3} g A_{[AABAC]}) \]  
\(^{(7)}\)
is the usual supergravity field-strength and

\[ M_{ABC} = \frac{1}{6!} \varepsilon_{ABC}^{D_1...D_7} M_{D_1...D_6;D_7} \]  
\(^{(8)}\)
the field-strength of the dual supergravity. The transition from flat indices to world ones is fulfilled by means of the 10-bein \( E_{M}^{A} \).

By means of the O(1.9) rotation over the flat index one can vanish, as usual, the \( E_{m}^{\alpha} \) -component of the 10-bein:

\[ E_{M}^{A} = \left( \begin{array}{cc} E_{\mu}^{\alpha} & E_{\mu}^{a} \\ E_{m}^{\alpha} & E_{m}^{a} \end{array} \right) = \left( \begin{array}{cc} e_{\mu}^{\alpha} & B_{\mu}^{a} e_{n}^{a} \\ 0 & e_{m}^{a} \end{array} \right) \]  
\(^{(9)}\)

The Scherk-Schwarz compactification procedure [10] used here. It means that any tensor with 4-indices and flat 6-indices is independent of the coordinates \( y^m \) of the internal manifold but the tensors with world 6-indices depend on \( y^m \) in the following way:

\[ V_{m...n...}(x,y) = V_{p...q...}^{(0)} (x) \bar{U}^p \ (y) \ldots \bar{U}^{n}_{q} \ (y) \ldots , \]  
\(^{(10)}\)
where \( \bar{U}^{m}_{n} \) is inverse to \( U^{m}_{n} \). (We shall see, however, that some exceptions to this rule are needed.) Due to [10] the \( y \)-dependence appears in the physical formulae only in the form:

\[ C_{lm}^{k} = 2 \bar{U}^{p}_{l} \bar{U}^{q}_{m} \partial_{[q} U^{p}_{k]} \]  
\(^{(11)}\)
Consequently it is necessary to require that all $C^k_{lm}$ must be constants. Then they are the structural constants of some group with generators $L_m = U^p_m \partial_p$

$$[L_m, L_n] = C^p_{mn} L_p$$

(12)

and hence obey the Jacobi identity:

$$C^p_{q[k} C^q_{lm]} = 0$$

(13)

In different expressions they enter usually in the following combination with $U$-matrices:

$$C^k_{lm} \equiv U^k_n C^n_{pq} U^p_l U^q_m$$

(14)

## 3 Constraints

We start from the search of a constraint in the fermionic sector because it is simpler than bosonic one. Moreover, the fermionic constraint has the same form both in usual and dual supergravity while bosonic constraints have not. As we have seen in introduction there are 24 Majorana spinors $\psi_m$ in $M^4$ which don’t enter in the multiplet of $N=4, D=4$ supergravity. So they must be eliminated by means of a condition like that:

$$\psi_m + a \Gamma_m \chi + b \text{Sp}(A_m \lambda) = 0$$

(15)

If we don’t want to break the supersymmetry algebra than we must to require the vanishing of the supersymmetry variation of the relation (15). One can expand this variation in powers of $\Gamma$-matrices:

$$[X + X(2) \Gamma(2) + X(4) \Gamma(4)] \varepsilon$$

The vanishing of this expression for arbitrary $\varepsilon$ implies

$$X = X(2) = X(4) = 0.$$ 

In general case these conditions have only trivial solution. But there are unique values of $a$ and $b$ such that $X$ and $X(4)$ are equal to zero identically and the only restriction $X(2) = 0$ has a nontrivial solution. So let us require that $a$ and $b$ take exactly these values. Then all the factors in (15) are fixed unambiguously. It explains also why terms of any other type are not written in (15).

The explicit form of $a$ and $b$ depends on the field parametrization. In our notations (1) the constraint (15) takes the form:

$$2 \phi \psi_m - \Gamma_m \chi + 2 \sqrt{2} \text{Sp}(A_m \lambda) = 0$$

(16)

The supersymmetry variation of (16) leads to:

$$\phi \bar{\omega}_{mAB} + \frac{1}{2} \psi_m \Gamma_{AB} \chi - E_m[A(\phi, B) - \psi_B | \chi] +$$

$$+ 2 \text{Sp}(A_m \mathcal{F}_{AB} - \frac{1}{8} \lambda \Gamma_{mAB} \lambda) = 0,$$

(17)
where \( \tilde{\omega}_{mAB} = e_m \tilde{\omega}_{cAB} \) is the superspace spin-connection:

\[
\tilde{\omega}_{ABC} = \omega_{ABC} + \frac{1}{2} T_{ABC} + S_{ABC},
\]

\( \omega_{ABC} \) is the ordinary spin-connection depending only on the 10-bein, \( S_{ABC} \) is given in (2).

One can show that the supersymmetry variation of (17) does not lead to any other restrictions at the mass-shell level. The following formula helps to do it:

\[
\delta \tilde{\omega}_{mAB} = \epsilon^\alpha R_{\alpha mAB},
\]

where \( \alpha \) – spinorial index, \( R \) is the supercurvature from [5].

Consequently (16) and (17) are all the constraints we must impose. In fact the relation (17) contains the five independent conditions:

\[
\partial_\mu [\phi g_{mn} + 2 \text{Sp}(A_mA_n)] = 0 \quad (19)
\]

\[
[\phi g_{q(m} + 2 \text{Sp}(A_qA_{m})] \overline{C}^\nu_{np} = 0 \quad (20)
\]

\[
K_{mn\nu} = -2 \partial_\mu [\phi B^n_{\nu}g_{mn} + 2 \text{Sp}(A_{\nu}A_m)] \quad (21)
\]

\[
K_{mn\nu} = -[\phi g_{qm}B^q_{\nu} + 2 \text{Sp}(A_qA_{\nu})] \overline{C}^\rho_{mn} \quad (22)
\]

\[
K_{mnp} = -[\phi g_{qm} + 2 \text{Sp}(A_qA_m)] \overline{C}^\nu_{np} \quad (23)
\]

Where \( K_{MNP} = E_M^A E_N^B E_P^C K_{ABC} \) denotes the following tensor:

\[
K_{ABC} = \phi T_{ABC} - 3 \psi[\Lambda B \psi C] + 3 \psi[\Lambda A \psi BC] \chi -
\]

\[
- \frac{1}{2} \text{Sp}(\lambda \Gamma_{MNP} \lambda) - 12 \text{Sp}(A_{[M}A_{N];P} + \frac{2}{3} g A_{[M}A_NA_P]) \quad (24)
\]

Until we don’t substitute the explicit expression for \( T_{ABC} \) in (24), the constraints (19) – (24) have the same form both in the usual and dual supergravity.

The condition (19) connects the 6-metric with other fields:

\[
\phi g_{mn} = \eta_{mn} - 2 \text{Sp}(A_mA_n),
\]

where \( \eta_{mn} = \eta^{(0)}_{pq} U^p_m U^q_n, \) \( \eta^{(0)}_{mn} \) is invariant tensor of the group (12). Consequently it must be the Killing tensor:

\[
\eta^{(0)}_{mn} = - C^p_{mn} C^q_{np} \quad (26)
\]

So the condition (24) is fulfilled automatically because the structural constants are completely antisymmetric over all indices due to (20) and (14):

\[
\overline{C}_{mnp} \equiv \eta_{mq} \overline{C}^l_{np} = \overline{C}_{[mnp]} \quad (27)
\]

But the conditions (21) – (23) have a different meaning for the reduction of usual and dual N=1, D=10 supergravity.

In usual supergravity one must use the expressions (5), (7) for the torsion \( T_{ABC} \). Here the tensor \( K_{MNP} \) gets a simple meaning

\[
K_{MNP} = -6 B_{[MN;P]} \quad (28)
\]
and the equation (21) may be integrated:

\[ 2B_{m\mu} = -\phi g_{mn}B^n_{\mu} - 2\text{Sp}(A_mA_{\mu}) \]  

Equations (22), (23) are transformed to:

\[ \partial_\mu B_{mn} = 0 \]  

\[ 6\partial_{[m}B_{np]} = \overline{C}_{mnp} \]  

If the \( B_{mn} \) component of the potential obeys the condition (31) it must depend on the \( y \)-coordinates in the way different from (11).

The results (16), (29) – (31) have the same form as in [8] (the different field parametrization used there) but we take into account all the terms in the formulae, not only the lowest order in fermionic fields.

The conditions (16), (29) – (31) don’t contain derivatives (30) and (31) may be easily integrated and consequently can be imposed at the lagrangian level. So in order to obtain the \( N=4, D=4 \) lagrangian one can substitute them into \( N=1, D=10 \) lagrangian and then express the \( B_{\mu\nu} \) through the pseudoscalar field \( B \) by means of the dual transformation. At the level of bosonic fields it has been done in [8].

In dual supergravity \( T_{ABC} \) is given in (6),(8) and the tensor \( K_{ABC} \) takes the form:

\[ K_{ABC} = 2\phi M_{ABC} - \frac{1}{2}\phi \psi_D \Gamma^{[D} \Gamma_{ABC} \Gamma^{E]} \psi_E - 3 \psi_{[A} \Gamma_{BC]} \chi - \frac{1}{2} \text{Sp}(\lambda \Gamma_{ABC} \lambda) - 12 \text{Sp}(A_{[A} A_{B;C]} + \frac{2}{3} g A_{[A} A_B A_C]) \]  

Relations (24) – (23) define all of the components of the field-strength \( M_{MNP} \) at the the reduction with except of \( M_{m1...m6} \); the potential \( M_{m1...m6} \), which enter in \( M_{\mu\nu\lambda} \), becomes directly the pseudoscalar field \( B \) of the \( N=4,D=4 \) supergravity:

\[ M_{m1...m6} = \epsilon_{m1...m6} B \]

where \( \epsilon_{m1...m6} = 1 \) if \( \{m_1...m_6\} = \{12...6\} \).

We see, that all the components of the potential \( M_{M1...M6} \) (with except of \( M_{m1...m6} \)) are expressed through other fields in a nonlocal manner. It is not a problem at the level of equations of motion because they contain only the field-strenght. But constraints (24) – (23), containing derivatives, cannot be imposed at the lagrangian level: if we try to obtain the \( N=4, D=4 \) lagrangian substituting them into \( N=1, D=10 \) one we would get a wrong result.

### 4 Potential

Finally we write the potential of the \( N=4, D=4 \) theory. It has been obtained in [8] for other field parametrization and therefore we omit many intermediate formulae.

The pseudoscalar field \( B \) doesn’t form a part of the potential because the theory is invariant relative to the transformation

\[ B \rightarrow B + C \]
where $C$ is a constant. Consequently in order to obtain the potential it is necessary to keep the terms only with $\phi$ and $A_m$ – fields in the usual supergravity lagrangian$^2$

$$L^{N=1,D=10} = \frac{1}{4} \phi R + \frac{1}{12} \phi^{-1} H^2 + \frac{1}{4} \text{Sp}(F^2)$$

(it is the famous lagrangian [2] rewritten in fields used here).

To derive the lagrangian with correctly normalized kinetic terms we replace our fields by primed ones

$$e_\mu^{\alpha'} = (\phi E)^{1/2} e_\mu^{\alpha}$$

$$e^{-2\phi'} = E$$

$$A_m' = A_m$$

$$\int e' L' = \int e L$$

(34)

where $e = \text{det} e_\mu^{\alpha}$, $E = \text{det} e_m^{\alpha}$, $L$ is the lagrangian,

$$g_{mn}' = \phi g_{mn} = \eta_{mn} - 2 \text{Sp}(A_mA_n)$$

and omit all the primes later on.

The scalar field lagrangian has the form:

$$L = L_T - U$$

$L_T$ is the kinetic part:

$$L_T = \frac{1}{4} R + \frac{1}{2} \phi,_{\mu} \phi^{;\mu} + \frac{1}{2} g^{mn}\text{Sp}(A_m;_{\mu}A_n^{;\mu}) +$$

$$+ g^{mp}g^{nq}\text{Sp}(A_mA_n^{;\mu})\text{Sp}(A_pA_q^{;\mu})$$

(36)

$U$ – the potential:

$$U = \frac{1}{16} e^{2\phi} \left\{ C_m^{np} g^{nq} (2 C_p^{mq} - g_{mr} C^r_{qs} g^{sp}) -$$

$$- \frac{1}{3} \left[ C_{mnp} - 2 \text{Sp}(3 A_qA[mC^n_{np} + 4 g A_{[m}A_nA_p]) \right]^2 -$$

$$- 4 \text{Sp} \left[ (A_pC^p_{mn} + 2 g A_{[m}A_nA_p])^2 \right] \right\}$$

(37)

where

$$g_{mn} = \eta_{mn} - 2 \text{Sp}(A_mA_n)\, , \quad C_{mnp} = \eta_{mq} C^n_{np}$$

$g^{mn}$ is inverse to $g_{mn}$; the contraction of the indices $m, n, \ldots$ is fulfilled by means of the $g^{mn}$ – tensor.

But as it was mentioned in [8], the potential (37) cannot lead to a realistic model.

Indeed, in the case $C_{mnp} = 0$ this potential is unbounded from below because $-g^{mn}$ is not positively definite and singular.

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$^2$As mentioned in previous section the dual supergravity is not convenient for this purpose because constraints (21) – (23) contain derivatives in this case.
In the case $C_{mnp} \neq 0$ there is a field configuration ($A_m = 0$) where $U$ takes a negative value:

$$U = \frac{1}{24} e^{2\phi} (C_{mnp})^2$$

It falls down infinitely together with the rise of the vacuum expectation value $<\phi>$. Moreover, if $<\phi> = 0$ by some reasons and the potential has a minimum, it must lie below zero. Hence, after spontaneously symmetry breaking the theory gets an enormous cosmological constant $\sim M_P^2$.

It is obviously that the abelian matter fields, eliminated here, cannot solve this problem.

## 5 Conclusion

We have described how to choose the N=4, D=4 supergravity degrees of freedom from the N=1, D=10 supergravity coupled with the YM-matter. The main problem of this theory is the non-positively definite potential. It is possible to solve this taking into account the Chern-Simons term in the field-strength $H$. In this case supersymmetry transformations have nonminimal corrections and the starting condition breaks the supersymmetry. Consequently this condition must be modified by adding appropriate nonminimal terms. But then the constraint becomes the equation of third order in $T_{ABC}$ and we don’t know whether it is solvable in a nonperturbative way or not.

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