Lawrence Berkeley Laboratory
UNIVERSITY OF CALIFORNIA

Physics Division

Submitted for publication

COMPOSITE ANOMALIES IN SUPERGRAVITY

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February 1985

Prepared for the U.S. Department of Energy under Contract DE-AC03-76SF00098

February 1985
LBL-19082
UCB-PTH-85/4

Composite Anomalies in Supergravity[1]

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Abstract

It is argued that the vectors of N-extended supergravity theories should be thought of as self-dual tensor fields, and that they contribute to the anomalies of the $(S)U(N)$ currents. Some implications of these anomalies are discussed. The total anomaly is found to vanish for $N > 4$, so the composite gauging scenario is consistent in these cases.

[1] This work was supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contracts DE-AC03-76SF00098

[2] Miller Fellow in Theoretical Physics

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The scalars of an extended supergravity theory parameterize a coset space G/H, where G is a noncompact Lie group and H its maximally compact subgroup. The theory has a physical global G symmetry acting on all the fields, but can be written with a global G symmetry acting only on the bosons and a local H symmetry acting on the fermions [1]. In this formalism the scalars are treated as a vielbein $\mathcal{V}$ transforming under both G and H. After noting that the compact (H) part of the "left invariant one form" $\mathcal{V}^{-1}\mathrm{d}\mathcal{V}$ transforms as a gauge field under H transformations, Cremmer and Julia conjectured that a dynamical mechanism could make these composite vector fields physical, resulting in an effective low energy theory with a gauged H group [1]. This scenario has important consequences but is very difficult to judge quantitatively, as it involves bound state problems.

In a recent paper [2] it was suggested that insight into this conjecture could be found by studying the question of anomalies in the H symmetry. This work was based on the study of anomalies in nonlinear sigma models [3]. Such anomalies would imply that the formulation of the theory with auxiliary scalar fields and a local H invariance is not equivalent to a "physical gauge" formalism, and that the H symmetry could not become dynamically gauged. The authors argued that only fermions contribute to this anomaly and, adding the contribution of the spinors and Rarita-Schwinger fields, concluded that the symmetry was indeed anomalous. This calculation is somewhat delicate since the scalar vielbein transforms under both G and H. Thus while the vector field strengths transform under G, their product with the (inverse) vielbein transforms under H. In this letter this problem is avoided by choosing a physical gauge. The results suggest that the calculation of ref. [2] is incorrect.

A natural physical parameterization of the scalar vielbein is $\mathcal{V} = e^{\phi}$, with ϕ in the noncompact part of the algebra [4]. In order to remain in this gauge, G transformations must be accompanied by particular H transformations, and the only remaining symmetry is a global G invariance. Unlike the $G \otimes H$ symmetry, this transformation acts on the states of the theory, giving a physical symmetry. The compact transformations act linearly on the fields generating a global H symmetry. The noncompact transformations are nonlinear, with the scalars transforming as Goldstone bosons to lowest order. The other fields transform linearly under the induced local H transformations. It must be emphasized that the parameters of these H transformations

helicity	2	3/2	1	1/2	0	$-1/2$	-1	$-3/2$	-2
SU(8)	1	8	28	56	70	$\overline{56}$	$\overline{28}$	$\overline{8}$	1
field	e^a_μ	$\psi^I_{L\mu}$	?	χ^{IJK}_L	ϕ^{IJKL}	χ^c_{IJK}	?	$\psi^c_{\mu I}$	e^r_μ

Table 1: The Particles and Fields of $N = 8$ Supergravity

are not arbitrary, but are particular combinations of the group elements and the scalar fields. Thus, while the theory has the structure of a gauge theory, it has no *elementary* gauge fields. The composite gauge fields are seen by defining the operator

$$\mathcal{V}^{-1}\mathbf{d}\mathcal{V} \equiv \mathbf{d} + \mathbf{Q} + \mathbf{P} \equiv \mathbf{D} + \mathbf{P} \quad . \tag{1}$$

Here $\mathbf{P} \sim \mathbf{d}\phi$ is in the coset and transforms covariantly. $\mathbf{Q}$ is in the subgroup and acts as a gauge field for the induced composite gauge transformations, so that $\mathbf{D}$ is the H-covariant derivative. The lagrangian has the structure of a locally gauge invariant theory with respect to $\mathbf{P}$ and $\mathbf{Q}$, and it is again reasonable to invoke dynamical gauging. It should be noted, however, that the field strength of the gauge symmetry is restricted by the Maurer-Cartan equation

$$\mathbf{F}[\mathbf{Q}] \equiv (\mathbf{d} + \mathbf{Q}) \wedge (\mathbf{d} + \mathbf{Q}) = -\mathbf{P} \wedge \mathbf{P} \quad . \tag{2}$$

This immediately follows upon squaring the operator of eq. (1).

To study the anomalies of the H transformation, it is necessary to know how the fields of the theory transform. At this point it is useful to be specific and to consider $N = 8$ supergravity. Here the symmetry group is $\mathrm{E}_{7,7}$, which has SU(8) as its maximally compact subgroup [1]. The SU(8) transformation is that of the physical supersymmetry algebra [5]

$$\{Q^I, \overline{Q}_J\} = \delta^I_J \quad .$$

The particle and field content of the theory are summarized in table 1. As indicated there, there is a problem with writing the helicity one particles. The other particles are represented by fields transforming covariantly under SU(8) (with the 70 scalars satisfying the "reality" condition $\phi^{IJKL} = \epsilon^{IJKLMNOP}\phi^*_{MNOP}$). The helicity one states, however, cannot be

represented by 28 vector fields as there is no *real* 28 dimensional representation of SU(8). The solution of this dilemma is that in the SU(8) transformation $F'^{IJ}_{\mu\nu} = \Lambda^I{}_K \Lambda^J{}_L F^{KL}_{\mu\nu}$, $iF_{\mu\nu}$ should be interpreted as

$$iF_{\mu\nu} \rightarrow \tilde{F}_{\mu\nu} \equiv 1/2\, \epsilon_{\mu\nu\rho\sigma} F^{\rho\sigma} \quad .$$

This implies that the SU(8) transformations rotate the Bianchi identity $\mathbf{d}F = 0$ into the equation of motion $\mathbf{d}\tilde{F} = 0$, and that SU(8) covariance only holds on shell. Some insight into the peculiar form of the SU(8) transformations can be found by examining the form of the theory with eight Majorana (real) gravitini $\Psi^I_\mu \equiv \psi^I_\mu + \psi^c_{\mu I}$. This is mathematically equivalent to using the Weyl gravitini, but now $i\,\Psi^I_\mu$ must be interpreted as $i\gamma^5\Psi^I_\mu$. For the helicity one particles the manifestly covariant object is

$$\mathcal{F}^{IJ}_{\mu\nu} \equiv F^{IJ}_{\mu\nu} - i\,\tilde{F}^{IJ}_{\mu\nu} \quad , \tag{3}$$

which is a 28 of SU(8). At the linearized level, and using the equations of motion, $\mathcal{F}^{IJ}$ satisfies

$$\mathcal{F}^{IJ} = i\,\tilde{\mathcal{F}}^{IJ} \tag{4}$$

and

$$\mathbf{d}\mathcal{F}^{IJ} = 0 \quad . \tag{5}$$

It is thus a "self-dual" tensor. As required by CPT invariance, its complex conjugate is an anti-self-dual tensor in a $\overline{28}$ of SU(8). Since it is a chiral field, it can contribute to anomalies. This is also plausible [6] since it is not possible to write a lagrangian that implies eqs (4) and (5), without breaking $\mathcal{F}^{IJ}_{\mu\nu}$ up as in eq. (3), thereby losing manifest SU(8) covariance. This is analogous to the case of a real self-dual tensor field in $4n+2$ dimensions, for which no manifestly *Lorentz* covariant lagrangian exists [7].

The full nonlinear equation of motion of this field can be written relatively simply using this formalism[4]. It is

$$\hat{\mathcal{F}}^{IJ} = i\,\tilde{\hat{\mathcal{F}}}^{IJ} \quad , \tag{6}$$

where $\hat{\mathcal{F}}^{IJ}$ is the "supercovariantized" generalization of $\mathcal{F}^{IJ}$ (appropriate fermion bilinears are added). The Bianchi identity, eq. (5), becomes

$$\mathcal{V}\mathbf{d}\mathcal{V}^{-1}\mathcal{F} = 0 \quad . \tag{7}$$

[4]This is derived from ref [1], where $\mathcal{F}^{IJ}_{\mu\nu}$ is denoted $\mathcal{F}'_{\mu\nu}$ and equals $\mathcal{V}'\mathcal{F}^F_{\mu\nu}$.

Expanding and using explicit indices, this gives

$$(\mathbf{D}\mathcal{F})^{IJ} = \mathbf{P}_{IJKL} \wedge \mathcal{F}^{KL} \quad . \tag{8}$$

This equation shows the typical gauge invariant "look" of the theory. One can again show that $\mathbf{F}$ must satisfy the Maurer-Cartan equation (2) for the consistency of the theory by applying $\mathbf{D}$ to eq. (8). The same result also arises in the gravitino sector.

We can now examine the question of anomalies. In general there are two types of anomalies: a pure "gluon" anomaly derived from the six form $\mathrm{tr}\,\mathbf{F}\wedge\mathbf{F}\wedge\mathbf{F}$ [8], and a mixed anomaly derived from $\mathrm{tr}\,\mathbf{F}\wedge\mathrm{tr}\,\mathbf{R}\wedge\mathbf{R}$, where $\mathbf{R}$ is the Riemann two form. The latter anomaly only exists if there is a U(1) factor in H. There is a question of interpretation in these formulae, since in the fundamental theory $\mathbf{F}$ is a composite object restricted by the Maurer-Cartan equation. However, if one is concerned with dynamical gauging, the resulting field strength should be thought of as that of a physical gluon, and the naive interpretation should be correct.

In principle the anomalies could be calculated directly from Feynman diagrams, but this would be somewhat tedious. It is also unclear how to evaluate the self-dual tensor contribution. The gauge field $\mathbf{Q}$ couples to an object like $F_{\mu\nu}A^{\mu}$, which is not covariant, so the trick used for the gravitational anomaly calculation in ref. [6] cannot be used. Fortunately it is now known that anomalies can be easily evaluated by the Atiyah-Singer index theorem [9]. The gravitational indexes of the three fields, used for the axial anomaly, are

$$\begin{aligned}
\mathcal{I}(1/2) &= \hat{\mathrm{A}} \equiv \prod_i \frac{x_i/2}{\sinh x_i/2} \rightarrow (1 - p_1/24) \\
\mathcal{I}(1) &= \hat{\mathrm{A}} \times \frac{1}{2}\prod_i 2\cosh x_i \rightarrow 2\,(1 + p_1/12) \\
\mathcal{I}(3/2) &= \hat{\mathrm{A}} \times \left(2\sum_i \cosh x_i - 1\right) \rightarrow 3\,(1 + 7p_1/24) \quad .
\end{aligned} \tag{9}$$

The expansions above are written for four dimensions, with

$$p_1 \equiv \sum_i x_i^2 = -\frac{1}{8\pi^2}\,\mathrm{tr}\,\mathbf{R}\wedge\mathbf{R} \quad .$$

For anomaly calculations, fermionic fields give an additional minus sign. The results for the spin 1/2 and spin 3/2 fields are given in refs [10]. The gauge

	1	X^I	X^{IJ}	X^{IJK}	X^{IJKL}
C_0	1	N	$\frac{N(N-1)}{2!}$	$\frac{N(N-1)(N-2)}{3!}$	$\frac{N(N-1)(N-2)(N-3)}{4!}$
C_2	0	1	$\frac{N-2}{1!}$	$\frac{(N-2)(N-3)}{2!}$	$\frac{(N-2)(N-3)(N-4)}{3!}$
C_3	0	1	$\frac{N-4}{1!}$	$\frac{(N-6)(N-3)}{2!}$	$\frac{(N-8)(N-3)(N-4)}{3!}$

Table 2: Group Theory Factors for SU(N)

anomaly is given by eq. (9) multiplied by the Chern class $(\mathrm{tr}\,e^{i\mathbf{F}/2\pi})$. In four dimensions this means that the anomaly contribution of a helicity h field is summarized by the six-form

$$\frac{-i}{3!(2\pi)^3}\,(-)^{2h}\,2h\left(\mathrm{tr}\,\mathbf{F}\wedge\mathbf{F}\wedge\mathbf{F} - \frac{(2-(2h)^2)}{8}\,\mathrm{tr}\,\mathbf{F}\wedge\mathrm{tr}\,\mathbf{R}\wedge\mathbf{R}\right) \quad , \tag{10}$$

The extension to arbitrary helicity fields is found by considering them to be symmetric products of Dirac spinors.

The other ingredient needed for the calculation is the value of the group theory traces. In N-extended supergravity the symmetry group is SU(8) for $N = 8$ and SU(N)⊗U(1) for $N \leq 6$. If there is a U(1) factor, eq. (10) gives rise to four possible anomalies. These involve the group theory quantities

$$\begin{aligned}
\mathrm{str}\,\lambda &= q\,C_0 \\
\mathrm{str}\,\lambda\,\lambda\,\lambda &= q^3\,C_0 \\
\mathrm{str}\,\lambda\,\lambda^a\lambda^b &= q\,\delta^{ab}\,C_2 \\
\mathrm{str}\,\lambda^a\lambda^b\lambda^c &= d^{abc}\,C_3 \quad ,
\end{aligned} \tag{11}$$

where str denotes the symmetrized trace and q is the charge of the representation. The dimensions C_0 and the Casimirs C_2 and C_3 are given in table 2. They can be calculated using the Chern class, but are more easily written by finding their zeros, and using

$$C_n(\overline{R}) = (-)^n\,C(R) \quad .$$

	SU(6)	W_{grav}	$q\,A$	W_{YM}	D	$q\,C$	$q^3\,A$
ψ_1^I	6_1	−21	6	3	1	1	6
F_2^{IJ}	15_2	4	30	−2	2	8	120
F_{-6}	1_{-6}	″	−6	″	0	0	−216
χ_3^{IJK}	20_3	1	60	1	0	18	540
χ_{-5}^I	6_{-5}	″	−30	″	1	−5	−750
Σw			0		0	0	0

Table 3: The Anomaly Contributions for $N = 6$ supergravity. The W's are the weights of the different helicities.

For $N = 8$ there is only one possible anomaly. The Casimir C_3 has the ratios 1 : 4 : 5 for the 8, 28 and 56 of SU(8). Substituting this into eq. (10), one sees that the anomaly vanishes :

$$(-3) \times 5 + (2) \times 4 + (-1) \times 1 = 0 \quad !$$

There is thus no obstruction to the dynamical gauging of SU(8) in $N = 8$ supergravity. It is also interesting to examine the work of Ellis, Gaillard and Zumino [11] in this light. They attempted to study the results of a composite gauging by considering the linear $N = 8$ supermultiplet containing the composite gauge field. Some of this multiplet was then interpreted as the physical composite particles. One constraint used was that the spin 1/2 particles be in an anomaly free representation, which led to the breaking of SU(8) to SU(5). It is amusing that, using eq. (10), the entire multiplet is already anomaly free. However, as shown above, it is inconsistent to couple unconstrained gauge fields to self-dual tensors (and gravitini), so the restriction to spinors may be justified.

The $N = 6$ case is summarized in table 3. There is a remarkable cancellation and all four anomalies vanish. A similar cancellation occurs for $N = 5$.

In $N = 4$ supergravity both the SU(4) and the U(1) symmetries are

anomalous. In that case, the self-dual tensor is a 6 of SU(4), which is a real representation. The entire SU(4) anomaly thus comes from the 4 spinor and the $\bar{4}$ spin 3/2 field. As the anomaly is proportional to the helicity, there is a relative factor of 3 between these contributions, and they do not cancel. The coefficient of the pure U(1) anomaly is

$$(-3) \times 4 \times 1^3 + (2) \times 6 \times 2^3 + (-1) \times 4 \times 3^3 = -24 \quad ,$$

and that of the mixed U(1) gravitational anomaly is

$$(-21) \times 4 \times 1 + (4) \times 6 \times 2 + (1) \times 4 \times 3 = -24 \quad ,$$

so they, too, do not vanish. As the scalars of $N = 4$ supergravity are in the coset SU(1,1)/U(1), there is no composite SU(4) gauge field and that symmetry would not be expected to become dynamical. The U(1) symmetry, on the other hand, would appear to be on the same footing as the SU(8) of $N = 8$ supergravity or the U(1)'s of $N = 5$ and $N = 6$ supergravity, and it is rather surprising that it is anomalous. For $N < 4$, all the symmetries are anomalous (except for SU(2) which has no complex representations), and there are no scalars to form composite gauge fields.

Aside from the issue of dynamical gauging, it is unclear what, if any, is the significance of these anomalies. Certainly if the anomaly vanishes, as in $N > 4$ supergravity, there can be no problems. If there are no scalars, as in the SU(4) of $N = 4$ supergravity, the Maurer-Cartan equation is simply $\mathbf{F} = 0$ so, presumably, the anomalies should be irrelevant. For the U(1) subgroup of $N = 4$ supergravity $\mathbf{F} = -\mathbf{P} \wedge \mathbf{P}^*$. This means that the pure U(1) anomaly becomes

$$\mathrm{tr}\,\mathbf{F} \wedge \mathbf{F} \wedge \mathbf{F} \to -(\mathbf{P} \wedge \mathbf{P}^*)^3 \quad , \tag{12}$$

which vanishes. However, the mixed anomaly becomes

$$-\mathbf{P} \wedge \mathbf{P}^* \wedge \mathrm{tr}\,\mathbf{R} \wedge \mathbf{R} \quad , \tag{13}$$

which can give rise to a U(1) anomaly in the fundamental theory.

The main issue has now been settled, but for completeness the question of composite anomaly cancellations in other dimensions can also be studied. I shall consider the maximal supergravity in six dimensions in some detail, and simply give the results of the other cases. In six dimensions particles are classified by the little group $SO(4) \approx SU(2) \otimes SU(2)$. The "$N = 4 + 4$" supergravity has two Sp(4) symmetries. The physical supersymmetries form a left-handed spinor in a four of one Sp(4) and a right-handed spinor in a four of the other Sp(4). As befits a maximally supersymmetric theory, there are a total of sixteen real components. The supersymmetry algebra is

$$\{Q^I, Q^J\} = \delta^{IJ} \quad ,$$

where the SO(16) indices of the Clifford algebra are broken into a $(2,1,4,1)$ and a $(1,2,1,4)$ of $SU(2) \otimes SU(2) \otimes Sp(4) \otimes Sp(4)$. The states of the theory are thus given by the spinors of SO(16) broken into representations of this algebra. The bosons are given by

$$(3,3,1,1) \oplus (1,1,5,5) \oplus (2,2,4,4) \oplus (3,1,1,5) \oplus (1,3,5,1)$$

and the fermions by

$$(3,2,1,4) \oplus (2,3,4,1) \oplus (1,2,5,4) \oplus (2,1,4,5) \quad .$$

The chiral fields of the theory are thus a self-dual tensor $F^+_{1,5}$, a left-handed gravitino $\psi^\mu_{1,4}$, a left-handed spinor $\chi_{4,5}$ and the opposite chirality fields with the two Sp(4)'s interchanged.

In six dimensions anomalies come from square graphs, and there are six possible nonvanishing combinations of external fields. Using the notation $y_n \equiv \sum_i (y_i)^n$, the Chern classes of the 4 and 5 of Sp(4) can be written to the desired order as

$$\begin{aligned} \mathrm{ch}(4) &\simeq 4 + \frac{y_2}{2} + \frac{y_2^2}{32} - \frac{y_4}{48} \\ \mathrm{ch}(5) &\simeq 5 + y_2 + \frac{y_4}{12} \quad . \end{aligned}$$

Using these formulae (with z instead of y for the second Sp(4)) together with eq. (10), one obtains the family index of the fields :

$$\begin{aligned} I(\chi_{4,5}) &= \left(1 - \frac{x_2}{24} + \frac{x_4}{2880} + \frac{x_2^2}{1152}\right)\left(4 + \frac{y_2}{2} + \frac{y_2^2}{32} - \frac{y_4}{48}\right)\left(5 + z_2 + \frac{z_4}{12}\right) \\ I(F^+_{1,5}) &= \left(4 + \frac{x_2}{3} - \frac{7x_4}{360} + \frac{x_2^2}{72}\right)\left(5 + z_2 + \frac{z_4}{12}\right) \\ I(\psi_{\mu\,1,4}) &= \left(5 + \frac{19x_2}{24} + \frac{49x_4}{576} - \frac{43x_2^2}{1152}\right)\left(4 + \frac{z_2}{2} + \frac{z_2^2}{32} - \frac{z_4}{48}\right) \qquad (14) \end{aligned}$$

One must next add these expressions, with an extra minus sign for the fermions, and antisymmetrize in y and z to take account of the right-handed fields. When this is done, the entire expression vanishes up to and including eighth order terms! The anomaly, which is derived from the eight-form piece of the expression thus vanishes.

This method can be used in all the remaining cases. In six dimensions the only other pure supergravity which does not have a fatal gravitational anomaly is the $N = 2 + 2$ theory, which has an $Sp(2) \otimes Sp(2)$ symmetry. Unlike the $N = 4 + 4$ theory, the Sp(2)'s are anomalous, but there are no scalars to make composite gauge fields. In eight dimensions the $N = 1$ supergravity has a U(1) invariance, and the $N = 2$ a $U(1) \otimes SU(2)$ invariance. The latter case is similar to four dimensions in that the three-index tensor from eleven dimensions rotates under the U(1), and should be thought of as a complex self-dual tensor. The SU(2) is trivially anomaly free, as its representations are real. In both theories the U(1) symmetry is anomalous. However only the $N = 2$ theory has the relevant fundamental scalars.

Finally we come to ten dimensions. There the $N = 1$ and the nonchiral $N = 1 + 1$ supergravities have no symmetries, and the chiral $N = 2$ supergravity has an $SO(2) \approx U(1)$ invariance. This symmetry is anomalous and no dynamical vector should form. This case is somewhat unusual since when the Maurer-Cartan equation is used all the anomalies vanish. The only possible anomaly remaining is analogous to the mixed anomaly of eq. (13). However, in ten dimensions this becomes

$$-\mathbf{P} \wedge \mathbf{P}^* \wedge \mathrm{tr}\,\mathbf{R}^5 \quad ,$$

which vanishes by the antisymmetry of $\mathbf{R}$. There are thus no anomalies in the fundamental theory.

The main results of this paper can be summarized as follows. In supergravity theories one should work with the physical symmetry, and not the fake local symmetry. The self-dual combination of the vector field strengths and their duals transforms covariantly under this symmetry and contribute to the anomalies of the compact part of the group. Such anomalies may affect the consistency of the theories and make different formalisms inequivalent. Furthermore, they ruin the possibility of dynamical gauging in the theory. In four dimensions only the U(1) of $N = 4$ supergravity has these anomalies. All the symmetries of $N > 4$ supergravity are anomaly free and could become gauge symmetries in the effective action.

I would like to thank Orlando Alvarez, Augusto Sagnotti, Jean Thierry-Mieg and Barton Zwiebach for many useful discussions.

This work was supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contracts DE-AC03-76SF00098.

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