$N = 8$ SUPERGRAVITY WITH LOCAL $SO(8) \times SU(8)$ INVARIANCE

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We present an extension of $N = 8$ supergravity in which the natural symmetry group SO(8) is gauged. Local SO(8) invariance is shown to be consistent with the dynamically realized SU(8) symmetry. We mention possible implications of this result for superunification.

It is widely believed that among all supergravity theories (see, e.g., ref. [1], where further references may be found) the $N = 8$ theory [2–4] stands out as the unique and most promising candidate for a unification of all fundamental interactions. Nevertheless, our knowledge of this model remains rather limited at present; although the basic lagrangian has been constructed, its possible extensions are not known and it is unclear how many free parameters it may contain. The question of its ultra-violet behaviour cannot be tackled as yet since we lack sufficiently well-developed techniques to deal with a model of such complexity. In view of these deficiencies, any study of the dynamics and the phenomenology of the model is preliminary, at best, even though there have been some encouraging results [5].

In this article we present an extension of the $N = 8$ theory which, in addition to the local SU(8) symmetry discovered in ref. [4], is invariant under a local SO(8) group of the conventional Yang-Mills type. Thus, in addition to the gravitational coupling constant κ (which we put equal to one in what follows), the $N = 8$ theory in the form given here admits a gauge-coupling constant g which, by the results of ref. [6], has a vanishing one-loop β function. As the gauging also generates a scalar field potential which is absent in the ungauged version, there arises the possibility of non-vanishing vacuum expectation values of the scalar fields or their composites which, in turn, might trigger spontaneous breaking of supersymmetry, SU(8) or both. Whether such a mechanism is compatible with previously proposed ones to obtain spontaneously broken $N = 8$ supergravity without gauging [7] remains an interesting problem for further study.

As is well known, the N independent supersymmetries of N extended supergravity are naturally combined with a rigid $SO(N)$ symmetry group which, by means of the $\binom{N}{2}$ vector fields of the graviton multiplet, may be promoted to a local $SO(N)$. This program has been carried out for $N \leqslant 4$ [8] and, more recently, for $N = 5$ [9]. The latter paper also contains lowest order results for $N = 8$.

In their pioneering paper [4], Cremmer and Julia demonstrated that extended supergravity theories have more symmetries than the aforementioned rigid $SO(N)$, which are hidden in the conventional approach. First, there is an (off-shell) local group H which is an invariance of the lagrangian and which is realized without a kinetic term for the H gauge connections. Secondly, there is an (on-shell) non-compact symmetry group G which is an invariance of the equations of motion. H is then isomorphic to the maximal compact subgroup of G, and the scalar fields of the theory are described as the coset space G/H. For $N = 8$, $G = E_7$ and

$H = SU(8)$, and it is this case which we will consider in this paper. All lower N supergravities may be obtained by consistent truncation.

The central result of our work is that the local invariance group SU(8) can be preserved in the process of gauging SO(8). To clarify the significance of this fact we recall how the gauging was achieved for $N \leqslant 5$. There one started from the theory in its manifestly $SO(N)$ invariant formulation, which is obtained by imposing a local H gauge. The $SO(N)$ group that was gauged was thus embedded in the rigid $SU(N) \times SU(N)$ subgroup of $G \times H$. Here we proceed differently by leaving the local SU(8) intact and by gauging the SO(8) subgroup of E_7 instead. The conventional form with local SO(8) and without $E_7 \times SU(8)$ may be recovered by fixing the local SU(8) gauge afterwards. At present, it is not known whether this theory follows from a non-trivial reduction from 11 dimensional supergravity [10].

There are essentially two reasons for keeping the local SU(8) throughout the calculation, the first being technical: the $E_7/SU(8)$ coset structure turns out to be just as indispensable for the construction and the consistency proof of gauged $N = 8$ supergravity as in the ungauged case, even though E_7 is no longer a symmetry of the theory since it is broken in a rather specific way. The second reason is even more important. It has been known for quite some time [11] that an SO(8) Yang-Mills group is too small to comprise the observed particle states, and on the basis of this observation it has been argued that a gauging is undesirable. Our results show that such an objection is not valid since both SU(8) and SO(8) may be relevant for the particle spectrum.

It is important to realize that the existence of the $N = 8$ theory with local $SO(8) \times SU(8)$ widens the range of possible unification scenarios. One which we find particularly appealing is the following. Unbroken Yang-Mills theories are known to have a confining phase; therefore, one may assume that the SO(8) group provides the force which binds the preons together ‡. In that case, all "observable" states would have to be SO(8) singlets. Since the only fields of the gravitational multiplet which carry SO(8) indices are the spin 0 and spin 1 fields (see table 1), we are led to the conclusion that, of the graviton multiplet, only the graviton, the gravitinos and the spin 1/2 fields are "observable", which would explain why the graviton is unconfined as was assumed in ref. [5]. Furthermore, since only part of the basic multiplet is "observable", the supersymmetry must be broken by the very same mechanism that confines some of the preons. Needless to say, this "preconfinement" mechanism would be non-perturbative and hence invisible in perturbation theory. In any case, it should be clear that, in presence of the gauging, the E_7 group may cease to play a significant role as far as unification is concerned.

‡ One possible objection to this picture is the conjectured vanishing of the β function to all orders (J. Ellis, private communication).

In our conventions we will closely follow refs. [3] and [9]. The field multiplet of the $N = 8$ theory contains one graviton e_μ^a which is a singlet with respect to $SO(8) \times SU(8)$, 8 gravitinos ψ_μ^i and 56 spin-1/2 fields χ^{ijk} which are assigned to representations of chiral SU(8) and singlets under SO(8), and 28 vector fields A_μ^{IJ} which transform as a 28 under SO(8). The scalar fields are represented by a 56-bein ("Sechsundfünfzigbein")

$$\mathcal{V} = \begin{bmatrix} u_{ij}{}^{IJ} & v_{ijKL} \\ v^{klIJ} & u^{kl}{}_{KL} \end{bmatrix}, \qquad (1)$$

which is an element of E_7 in the fundamental representation. $\mathcal{V}$ transforms under local SU(8) from the left, and under E_7 from the right; its inverse can be written in terms of the same submatrices u and v

$$\mathcal{V}^{-1} = \begin{bmatrix} u^{ij}{}_{IJ} & -v_{klIJ} \\ -v^{ijKL} & u_{kl}{}^{KL} \end{bmatrix}. \qquad (2)$$

All assignments have been collected in table 1. Note that these are off-shell assignments in contradistinction to the ungauged case with E_7(rigid) × SU(8)(local) where only the field strengths and their duals, but not the vector fields themselves, could be fitted into on-shell representations of E_7.

Table 1

	e_μ^a	ψ_μ^i	A_μ^{IJ}	χ^{ijk}	$u_{ij}{}^{IJ}$	v^{ijIJ}
SO(8)	1	1	28	1	28	28
SU(8)	1	8	1	56	$\overline{28}$	28

We write $N = 8$ lagrangian as follows:

$$\mathcal{L} = -\tfrac{1}{2} eR(\mathrm{e}, \omega) - \tfrac{1}{2} \epsilon^{\mu\nu\rho\sigma} \overline{\psi}^i_\mu \gamma_\nu \overleftrightarrow{\mathrm{D}}_\rho \psi_{\sigma i}$$
$$- \tfrac{1}{8} e [F^+_{\mu\nu IJ} F^+_{\mu\nu KL} (2S^{IJ,KL} - \delta^{IK}\delta^{JL}) + \text{h.c.}]$$
$$- \tfrac{1}{2} e [F^+_{\mu\nu IJ} O^{+\,KL}_{\mu\nu} S^{IJ,KL} + \text{h.c.}]$$
$$- \tfrac{1}{2} e \overline{\chi}^{ijk} \overleftrightarrow{\not{\mathrm{D}}} \chi_{ijk} - \tfrac{1}{96} e |\mathcal{A}^{ijkl}_\mu|^2$$
$$- \tfrac{1}{12} e [\overline{\chi}_{ijk} \gamma^\nu \gamma^\mu \psi_{\nu l} \mathcal{A}^{ijkl}_\mu + \text{h.c.}]$$
$$+ \text{SU(8)-invariant four-fermion terms} \qquad (3)$$

where $F^+_{\mu\nu IJ}(F^{-\,IJ}_{\mu\nu})$ denotes the self-dual (antiself-dual) abelian field strengths of the 28 vectors A^{IJ}_μ

$$F_{\mu\nu}{}^{IJ} = F^+_{\mu\nu IJ} + F^-_{\mu\nu}{}^{IJ} = 2\partial_{[\mu} A^{IJ}_{\nu]} \,. \qquad (4)$$

$S^{IJ,KL}$ is defined in terms of the submatrices of the 56-bein by the condition

$$(u_{ij}{}^{IJ} + v_{ijIJ}) S^{IJ,KL} = u_{ij}{}^{KL} \qquad (5)$$

and $O^{+}_{\mu\nu}{}^{IJ}$ is defined by

$$u^{ij}{}_{IJ} O^{+}_{\mu\nu}{}^{IJ} - (\sqrt{2}/144) \eta \epsilon^{ijklmnpq} \overline{\chi}_{klm} \sigma_{\mu\nu} \chi_{npq}$$
$$- \tfrac{1}{2} \overline{\psi}_{\lambda k} \sigma_{\mu\nu} \gamma^\lambda \chi^{ijk} + (\sqrt{2}/2) \overline{\psi}^i_\rho \gamma^{[\rho} \sigma_{\mu\nu} \gamma^{\sigma]} \psi^j_\sigma \,. \qquad (6)$$

The derivatives in (3) are covariant with respect to local Lorentz and local SU(8) transformations; hence, besides the standard spin connection ω^{ab}_μ, we have SU(8) gauge fields $\mathcal{B}_\mu{}^i{}_j$ which satisfy

$$(\mathcal{B}_\mu{}^i{}_j)^* = -\mathcal{B}_\mu{}^j{}_i; \quad \mathcal{B}_\mu{}^i{}_i = 0 \,; \qquad (7)$$

These gauge fields occur in D_μ according to

$$\mathrm{D}_\mu \phi^i = \partial_\mu \phi^i + \tfrac{1}{2} \mathcal{B}_\mu{}^i{}_j \phi^j \,. \qquad (8)$$

where ϕ^i is an SU(8) vector in the fundamental representation.

The SU(8) gauge fields $\mathcal{B}_\mu{}^i{}_j$ do not correspond to dynamic degrees of freedom (at least classically); they can be expressed in terms of the physical fields of $N = 8$ supergravity. This dependence can be viewed as the result of an algebraic equation of motion (first order form) or of a conventional constraint (second order form). For our purposes it is most convenient to choose the second option; this brings our results in direct correspondence with those of ref. [3]. Also, the quantity $\mathcal{A}^{ijkl}_\mu$, which characterizes the scalar kinetic terms in (3), is dependent. The dependence of $\mathcal{A}_\mu$ and $\mathcal{B}_\mu$ on the 56-bein is determined by the requirement of E_7 invariance; the only quantity of that kind which contains one derivative is given by $\mathrm{D}_\mu \mathcal{V} \cdot \mathcal{V}^{-1}$. This 56 × 56 matrix transforms covariantly under local SU(8), and takes its values in the Lie algebra of E_7. The dependence of $\mathcal{A}_\mu$ and $\mathcal{B}_\mu$ is now defined by

$$\mathrm{D}_\mu \mathcal{V} \cdot \mathcal{V}^{-1} = \begin{bmatrix} 0 & -\tfrac{1}{4}\sqrt{2}\, \mathcal{A}_{\mu ijkl} \\ -\tfrac{1}{4}\sqrt{2}\, \mathcal{A}^{mnpq}_\mu & 0 \end{bmatrix} . \qquad (9)$$

The diagonal blocks of (9) characterized the SU(8) subalgebra and are used to express $\mathcal{B}_\mu$ in terms of the submatrices of $\mathcal{V}$ and their derivatives. The part of the algebra orthogonal to SU(8) defines $\mathcal{A}_\mu$ in a similar fashion. Note that $\mathcal{A}_\mu$ does not explicitly depend on $\mathcal{B}_\mu$ in this way. Its classification as a component of the E_7 Lie algebra implies

$$\mathcal{A}_\mu{}^{ijkl} = \tfrac{1}{24} \eta \epsilon^{ijklmnpq} \mathcal{A}_{\mu mnpq} \,. \qquad (10)$$

This is a typical example of the kind of argument that is of crucial importance throughout this paper. It is not possible to show the validity of (10) directly from the explicit dependence of $\mathcal{A}_\mu$ on the 56-bein; instead, we have to rely on group-theoretic arguments based on E_7.

It is possible to view the matrix

$$\begin{bmatrix} \mathcal{B} & \mathcal{A}^* \\ \mathcal{A} & \mathcal{B}^* \end{bmatrix}$$

as the connection of a local E_7 group. In that context, (9) specifies that the connection is pure gauge and has vanishing E_7 field strengths. Indeed, application of a second SU(8) covariant derivative D_ν on (9) and antisymmetrization in μ and ν leads to

$$([\mathrm{D}_\mu, \mathrm{D}_\nu]\mathcal{V})\mathcal{V}^{-1}$$
$$= -\frac{1}{8} \begin{bmatrix} 0 & \mathcal{A}^*_\mu \\ \mathcal{A}_\mu \mathcal{V} & 0 \end{bmatrix} \begin{bmatrix} 0 & \mathcal{A}^*_\nu \\ \mathcal{A}_\nu & 0 \end{bmatrix}$$
$$- \frac{\sqrt{2}}{4} \begin{bmatrix} 0 & \mathrm{D}_\mu \mathcal{A}^*_\mu \\ \mathrm{D}_\mu \mathcal{A}_\nu & 0 \end{bmatrix} - (\mu \leftrightarrow \nu) \,. \qquad (11)$$

On the other hand, the commutator of two covariant derivatives is equal to the field strength; as SU(8) acts on $\mathcal{V}$ from the left, the left-hand side of (11) takes the simple form

$$([\mathrm{D}_\mu, \mathrm{D}_\nu]\mathcal{V}) \cdot \mathcal{V}^{-1} = \begin{bmatrix} \delta^{[i}_{[m} \mathcal{F}(\mathcal{B})_{\mu\nu n]}{}^{j]} & 0 \\ 0 & \delta^{[k}_{[p} \mathcal{F}(\mathcal{B})_{\mu\nu}{}^{l]}{}_{q]} \end{bmatrix} \qquad (12)$$

where $\mathcal{F}(\mathcal{B})_{\mu\nu}{}^i{}_j$ denotes the SU(8) field strengths. Equations (11) and (12) play a crucial role in establishing the invariance of the supergravity action; this requires a number of partial integrations which lead to SU(8) field strengths by means of the Ricci identity (12) or to derivatives on $\mathcal{A}_\mu$. By using (11) and (12) we find that the antisymmetric derivative of $\mathcal{A}_\mu$ vanishes, and that the SU(8) field strength can be expressed in terms which cancel against other variations. Indeed, (11) and (12) were previously found by requiring supersymmetry invariance of the action [3].

The gauging of SO(8) is now effected by further extending the covariant derivative with respect to local SO(8) embedded in the E_7 group. For instance, we define

$$\mathrm{D}_\mu u_{ij}{}^{IJ} \equiv \partial_\mu u_{ij}{}^{IJ} + \mathcal{B}_\mu{}^k{}_{[i} u_{j]k}{}^{IJ} - 2g A_\mu^{K[I} u_{ij}{}^{J]K} , \qquad (13)$$

where A_μ^{IJ} are the 28 vectors of $N = 8$ supergravity. At the same time, we replace the field strengths $F_{\mu\nu}^{IJ}$ by their fully SO(8) covariant counterparts

$$F_{\mu\nu}{}^{IJ} \equiv 2\partial_{[\mu} A_{\nu]}^{IJ} - 2g A_{[\mu}^{IK} A_{\nu]}^{KJ} . \qquad (14)$$

The presence of the order g terms in the lagrangian and transformation rules violates the supersymmetry invariance of the original action. To re-establish the invariance, one has to introduce new terms in the lagrangian and transformations. These can be parametrized in terms of three tensorial functions A_{1-3} which depend on the scalars contained in the 56-bein. The parametrization takes the following form [9]

$$\begin{aligned} \delta_g \overline{\psi}^i_\mu &= -\sqrt{2} g \overline{\epsilon}_j \gamma_\mu A_1^{ji} \\ \delta_g \overline{\chi}^{ijk} &= -2g \overline{\epsilon}^l A_{2l}{}^{ijk} , \\ \mathcal{L}_g &= \sqrt{2} g e A_{1ij} \overline{\psi}^i_\mu \sigma^{\mu\nu} \psi^j_\nu + \text{h.c.} \\ &\quad + \tfrac{1}{6} g e A^i_{2jkl} \overline{\psi}^\mu_i \gamma_\mu \chi^{jkl} + \text{h.c.} \\ &\quad + g e A_3^{ijk,lmn} \overline{\chi}_{ijk} \chi_{lmn} + \text{h.c.}, \\ \mathcal{L}_{g^2} &= g^2 e (\tfrac{3}{4} |A_1^{ij}|^2 - \tfrac{1}{24} |A^i_{2jkl}|^2) . \end{aligned} \qquad (15)$$

Note that the SU(8) tensors A_{1-3} must satisfy certain symmetry properties as a consequence of the way in which they appear in (15).

The new variations of the lagrangian induced by the SO(8) covariantization come from three sources. One set corresponds to a direct covariantization of the variations of the original lagrangians; because of the supersymmetry invariance of the original action, these covariantizations vanish as well. The second class of terms originates from the standard variation of A_μ^{IJ} in the covariant derivatives. The third class is generated through the Ricci identity which now, in addition to the SU(8) field strengths, lead to new g dependent terms proportional to the SO(8) field strengths. However, only vectors and scalars are not inert under SO(8) (see table 1) and, in fact, it is easy to show that all new variations containing the SO(8) field strength are governed by eqs. (11) and (12), but now with derivatives that are also covariant with respect to local SO(8).

The variations proportional to the SO(8) field strength are sufficient to determine the functions A_{1-3}. One first considers the SO(8) modification in (12) which takes the form

$$([\mathrm{D}_\mu, \mathrm{D}_\nu]\mathcal{V}) \cdot \mathcal{V}^{-1} = \text{eq. (12)} + 2g \mathcal{V} F_{\mu\nu} \mathcal{V}^{-1}, \qquad (16)$$

where

$$F_{\mu\nu}(A) \equiv \begin{bmatrix} \delta^{[I}_{[M} F(A)_{\mu\nu N]}{}^{J]} & 0 \\ 0 & \delta^{[K}_{[P} F(A)_{\mu\nu}{}^{L]}{}_{Q]} \end{bmatrix} . \qquad (17)$$

It is of crucial importance in the explicit calculations that the right-hand side of (16) takes values in the Lie algebra of E_7, even after the introduction of the gauge coupling.

The extra term (16) occurs in $g\overline{\epsilon}\psi$ and $g\overline{\epsilon}\chi$ variations of the lagrangian whose cancellations require the following values of A_{1-3}

$$\begin{aligned} A_1^{ij} &= -(4/21)\, T_m{}^{ijm} , \\ A_{2m}{}^{ijk} &= -(4/3)\, T_m{}^{[ijk]} , \\ A_3^{ijk,lmn} &= (\sqrt{2}/108)\, \eta \epsilon^{ijkpqr[lm} T^{n]}{}_{pqr} , \end{aligned} \qquad (18)$$

where the SU(8) tensor T is defined through

$$T_l{}^{kij} \equiv (u^{ij}{}_{IJ} + v^{ijIJ}) \times (u_{lm}{}^{JK} u^{km}{}_{KI} - v_{lmJK} v^{kmKI}) , \qquad (19)$$

which is manifestly antisymmetric in i and j. Note that T is not invariant with respect to the full E_7 group, but it is invariant under local SO(8) and covariant under SU(8).

In order for the solution (18) to be consistent, the tensor T has to obey a number of non-trivial identities; these can be shown to follow from the E_7 structure of the 56-bein. For instance, one can prove that T can be decomposed as follows:

$$T_l{}^{kij} = T_l{}^{[kij]} + \tfrac{2}{7}\delta_l^{[i}T_m{}^{j]mk} , \qquad (20)$$

where $T_m{}^{jmk}$ is symmetric in j, k and $T_l{}^{[kij]}$ is completely antisymmetric in kij and traceless, i.e.,

$$T_k{}^{[kij]} = 0. \qquad (21)$$

Having found the solution, one must now establish the invariance of the lagrangian in order g and g^2. This requires further identities on T, such as

$$\begin{aligned} D_\mu T_l{}^{kij} = &-\tfrac{1}{4}\sqrt{2} \\ &\times \{-2\mathcal{A}_\mu^{kmn[i}T^{j]}{}_{lmn} - \mathcal{A}_\mu^{ijmn}T^k{}_{lmn} \\ &+ \tfrac{2}{3}\delta_l^{[i}T^{j]}{}_{mnp}\mathcal{A}_\mu^{kmnp} + \tfrac{1}{3}\delta_l^k\mathcal{A}_\mu^{mnp[i}T^{j]}{}_{mnp}\} , \end{aligned}$$

$$\begin{aligned} &A_{2p}{}^{ijk}A^p_{2\,lmn} - 2\delta^{ijk}_{lmn}|A_2|^2 \\ &+ 18\delta^{[ij}_{[lm}A^r_{2\,n]pq}A_{2r}{}^{k]pq} - 6\delta^{[ij}_{[lm}A^{k]}_{2\,pqr}A_{2n]}{}^{pqr} \\ &- 9\delta^{[i}_{[l}A^j_{2\,mpq}A_{2\,n]}{}^{k]pq} - 9A_{2[l}{}^{p[ij}A^{k]}_{2\,mn]p} = 0 , \end{aligned} \qquad (22)$$

where we remind the reader that, by (18), A_2 is just the antisymmetric part of T.

It is instructive to impose an SU(8) gauge choice which corresponds to an explicit parametrization of the $E_7/SU(8)$ coset space. In the symmetric gauge [4], the 56-bein becomes

$$\mathcal{V} = \exp\begin{bmatrix} 0 & -\tfrac{1}{4}\sqrt{2}\,\phi_{ijkl} \\ -\tfrac{1}{4}\sqrt{2}\,\phi^{mnpq} & 0 \end{bmatrix} . \qquad (23)$$

Inserting this parametrization into T, we obtain the functions A_{1-3} as infinite series of the self-dual scalar fields ϕ^{ijkl}. We here give the expansions for A_1 and A_2 up to cubic order

$$\begin{aligned} A_1^{ij} = &(1 - \tfrac{1}{96}|\phi|^2)^{-1/2}\delta^{ij} \\ &+ \tfrac{1}{96}\sqrt{2}\,\phi^{ikmn}\phi_{mnpq}\phi^{pqkj} + O(\phi^4) , \end{aligned}$$

$$\begin{aligned} A_{2l}{}^{ijk} = &-\tfrac{1}{2}\sqrt{2}(1 - \tfrac{1}{144}|\phi|^2)\phi^{ijkl} - \tfrac{3}{8}\phi_{mnl[i}\phi^{jk]mn} \\ &+ \tfrac{1}{16}\sqrt{2}\,\phi_{lpqr}\phi^{pqs[i}\phi^{jk]rs} + O(\phi^4) , \\ |\phi|^2 \equiv &\ \phi^{ijkl}\phi_{ijkl} , \end{aligned} \qquad (24)$$

A_3 is related to A_2 by (18). In the reduction to $N = 5$, these results can be compared with those of ref. [9]. Both SO(8) and SU(8) are affected by the gauge choice (23); however, the non-trivial SO(8) subgroup acting on both SO(8) and SU(8) indices is preserved. Determining $\mathcal{B}^i_{\mu j}$ from (9), we find

$$\mathcal{B}_\mu{}^i{}_j = -2gA^{ij}_\mu + \text{nonlinear SO(8)-covariant terms.} \qquad (25)$$

Observe that in the gauge (23) there is no longer a distinction between SO(8) and SU(8) indices. Equation (25) shows that A^{IJ}_μ will now couple minimally to all fields that carry SO(8) and/or SU(8) indices. This corresponds to the standard formulation of supergravity with local SO(N) that has been obtained for $N \leqslant 5$.

A more detailed exposition of these results will be given in a forthcoming publication [12].

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