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SUPERGRAVITY, A GAUGE THEORY FOR PARTICLE UNIFICATION

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A B S T R A C T

Attempts to use supersymmetric theories to describe physics below the Planck scale are reviewed. Different options for the embedding of a Yang-Mills internal symmetry in supersymmetric schemes are discussed.

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There are several options for the unification of particle interactions in the framework of supersymmetry (SUSY). The final goal is to reproduce an effective renormalizable gauge theory based on the group $SU(3)_C \times SU(2)_{\text{weak}} \times U(1)_Y$ which is eventually broken down to $SU(3)_C \times U(1)_{\text{em}}$.

The internal symmetries are restricted by no-go theorems¹⁾ to be compact Lie groups, which commute with the space-time Poincaré symmetry. However, the supersymmetry spinorial generators Q_α^i ($i = 1 \dots N$ in N -extended supersymmetry) need not commute with the internal symmetries because they belong to a graded Lie algebra and obey anticommutation relations. If we denote by G the internal gauge symmetry with the property that

$$G \supset SU(3) \times SU(2) \times U(1)$$

then two possible alternatives can be envisaged:

$$[G, \text{SUSY}] = 0 \quad (1)$$

$$[G, \text{SUSY}] \neq 0 \quad (2)$$

The second alternative is only possible in the framework of extended supergravity (2). In fact, if G does not commute with the spinorial generators Q_α^i , supersymmetry transformations will involve space-time dependent anticommuting spinors $\varepsilon(x) = T_g(x)\varepsilon$ where T_g is the representation of the group G on the spinorial charges. But when the supersymmetry transformations become space-time dependent from the basic anticommutation relations

$$\{Q_\alpha^i, Q_\beta^j\} = (\gamma^\mu C)_{\alpha\beta} P_\mu \delta^{ij} \quad (3)$$

it follows that also translations must be local (general coordinate transformations). This argument indeed shows that possibility (2) can only be realized in extended supergravity theories. Needless to say, this second option is the most attractive from a theoretical point of view because it realizes a true marriage of gravitation with the other fundamental forces. In this situation, the internal symmetry group which does not commute with the spinorial charges is severely restricted by a simple

analysis of particle supermultiplets in N-extended supergravity.

From the structure of massless supermultiplets³⁾, it follows that the state of maximal helicity $\lambda_{\max}$ in any multiplet satisfies the bound

$$\lambda_{\max} \geq \frac{N}{4} \left(\frac{N+1}{4} \text{ for odd } N \right) \quad (4)$$

where N is the number of spinorial charges. Therefore interactions with elementary fields with spin not greater than two, can only exist for $N \leq 8$. For $N > 8$, spins higher than two must be present in any supermultiplet and no consistent interaction for such high spin fields is known. We will consider later, in more detail, some attempts based on extended supergravities for the ultimate unification.

Let us temporarily focus on the possibility (1). In this case, the gauge group G commutes with the supersymmetry generators. As a consequence, particles of a given supermultiplet will lie in the same representation of G. If we want to describe non-gravitational interactions, the elementary particles will be described by spin 1, $\frac{1}{2}$ and 0 fields. Because of (4), this limits the choice of the extended supersymmetry algebra to $N = 1, 2$ and 4 ($N = 3$ reproduces the same multiplet with $\lambda_{\max} = 1$ as $N = 4$). For $N = 1, 2$ we have two kinds of supermultiplets, matter multiplets with spin 0 and $\frac{1}{2}$ and gauge multiplets with spin 1 and $\frac{1}{2}$ for $N = 1$ and spin 1, $\frac{1}{2}$ and 0 for $N = 2$. For $N = 4$ we have a unique supermultiplet with spin ranging from 0 to 1. Since spin-one particles are described by gauge vector fields which belong to the adjoint representation of the group G, it follows that $N = 4$ supersymmetry can only describe a unique self-interaction of the $N = 4$ gauge supermultiplet in which all particles sit in the adjoint representation of the gauge group. The only arbitrariness in this $N = 4$ theory is the choice of the gauge group G. An interesting option would be to identify G with the exceptional group E_6 because for this group the fundamental and adjoint representations coincide⁴⁾ and therefore the spin $\frac{1}{2}$ partners of the gauge vector bosons could be naturally identified with quarks and leptons. However, there are many problems in trying out this identification; in particular, the fact that our low-energy world seems to be

remarkably well described by left-handed fermions, some of them having no right-handed counterpart. If we really believe that chiral fermions are fundamental then the only supersymmetry which permits left-handed fermions in a complex representation of the group G is $N = 1$ supersymmetry. This is because the $N = 1$ matter supermultiplet contains only a $\lambda = -\frac{1}{2}$ helicity state. In $N = 2$ supersymmetry the matter multiplet contains $\lambda = \pm\frac{1}{2}$ helicity states related by supersymmetry transformations. Since the group G commutes with the Q_α^i , this means that they belong to the same representation R of G . This means that the $N = 2$ matter multiplet describes Dirac particles in some representation R of the group G .

The use of supersymmetric gauge theories with a global $N = 1$ supersymmetry⁵⁾ to describe particle physics in the low-energy ($\ll M_{\text{Planck}}$) regime has been recently explored in some detail⁶⁻⁸⁾.

The breaking of $N = 1$ supersymmetry at an energy scale in the TeV region is perhaps the unique solution to the so-called hierarchy problem of Grand Unified Theories⁶⁻⁸⁾. This is related to the absence of scalar quadratic mass renormalization in $N = 1$ supersymmetric gauge models based on a semi-simple gauge group G ⁹⁾. Moreover, if supergravity is broken in the 1 TeV region, the so-called desert of conventional GUTs would be eventually substituted by a plethora of new particles and new physics may be the welcome surprise for the next generation of particle accelerators. A most dramatic consequence of this supersymmetric scenario is that these new particles in the 1 TeV region would probably push the proton decay rate far beyond the present experimental capability¹⁰⁾ and this may be one of the crucial tests for the existence of a low-energy supersymmetry.

We would like now to come back to the alternative scenario given by possibility (2), namely that the gauge symmetry of low-energy physics, as well as the group of Grand Unification do not commute with the supersymmetry generators.

In this case, a possible scenario is that extended supersymmetry is broken at the Planck scale and the first main constraint imposed on the theory would be the internal symmetry group, its breaking to $SU(3) \times SU(2) \times U(1)$ and the choice of fermion and scalar representations.

GUTS can be embedded in extended supergravity in essentially two ways. N-extended supergravity has a manifest $SO(N)$ global symmetry. This symmetry can be gauged¹¹⁾ and the associated gauge fields are just the $\lambda = \pm 1$ helicity partners of the $\lambda = \pm 2$ graviton. In fact supersymmetry assigns to these $\lambda = \pm 1$ states just the right multiplicity $N(N-1)/2$ in order to transform according to adjoint representation of $SO(N)$. Since supergravities seem to exist up to $N = 8$, the biggest gauge group we can think of in this way is $SO(8)$. This group does not contain $SU(3) \times SU(2) \times U(1)$. In addition, the spin $\frac{1}{2}$ spectrum of $SO(8)$ supergravity is not rich enough (56 states) to accommodate all observed quarks and leptons with the right $SU(3) \times U(1)_{em}$ quantum numbers¹²⁾. A last minor difficulty within this unification scheme is that GUTs based on orthogonal groups with particles in tensor representations would give completely vector-like theories and it would be difficult to explain the chirality of our low-energy world without some additional hypotheses. The second and more promising possibility to embed GUTs in supergravity is to use the hidden local symmetries discovered by Cremmer and Julia¹³⁾. These hidden symmetries are closely related to the fact that in extended supergravities with $N \geq 4$ a global non-compact invariance H is present¹⁴⁾. The elementary scalar fields which belong to the graviton supermultiplet parametrize a coset space H/G where $G = SU(N)$, $[U(N) \text{ for } N \neq 8]$ is isomorphic to the hidden local symmetry of the theory.

The gauge symmetry G is hidden because there are no elementary gauge fields associated to G in the supergravity Lagrangian. These gauge fields are composite operators and do not correspond to any degree of freedom. However, in analogy to CP^N models in two-dimensions, the gauge fields could become dynamical through non-perturbative effects of supergravity and eventually describe an effective GUT based on a subgroup of G ¹³⁾.

Ellis, Gaillard, Maiani and Zumino (EGMZ)¹⁵⁾ have further suggested that if composite gauge fields of N-extended supergravity bound to give massless particles, since in a supergravity theory these states must occur in supermultiplets, it is natural to assume that also the spin $\frac{1}{2}$ and spin 0 particles of the effective GUT are bound states of the preonic fields of the basic supergravity Lagrangian.

In Ref. (15), the massless supermultiplet containing the $\lambda = \pm 1$ helicity states in the adjoint representation of $SU(N)$ was identified to be

$$\left\{ \frac{3}{2}, [1]_N = N \right\} = \left(\frac{3}{2} \right)^A (1)^A_B \left(\frac{1}{2} \right)^A_{BC} \dots \left(\frac{3-N}{2} \right)^A \quad (5)$$

+ PCT conjugate states

Here we denote by $[K]_N$ the K -fold antisymmetric representation of $SU(N)$ and $[\bar{K}]_N = [N-K]_N$. The PCT conjugate states are those with reversed helicities

$$\left(\frac{N-3}{2} \right)_A \left(\frac{N-4}{2} \right)_{A,B} \left(\frac{N-5}{2} \right)_{A,B,C} \dots \left(-\frac{3}{2} \right)_A$$

EGMZ have assumed that all particles of the effective GUT are contained in the supermultiplet given by (5). Ellis, Gaillard and Zumino¹⁶⁾ (EGZ) have formulated some criteria which

- i) select $N = 8$ supergravity
- ii) select $SU(5)$ as the effective GUT
- iii) obtain three families of left-handed spin $\frac{1}{2}$ fermions.

Let us remind the hypotheses of Ref. (16).

- a) All states of the effective GUT come from the same supermultiplet (5). In particular the $\lambda = 0, \frac{1}{2}$ and 1 states are in the following $SU(8)$ representations:

$$\lambda=1 : 63 + 378 + (1+70) ; \lambda=0 : 420 (+28) ; \lambda=-\frac{1}{2} : 504 + 2\overline{16} + (\overline{56} + 8) \quad (6)$$

where the representations in brackets refer to single column Young-Tableaux (traces).

- b) $|\lambda| > 1$ states, $|\lambda| = 1$ states not in adjoint representations and traces are unbound.
- c) The effective GUT is defined by the maximal set of states $\lambda = -\frac{1}{2} \subset 504 + 2\overline{16}$ with an anomaly-free gauge group contained in $SU(8)$.

From (a), (b), and (c) we get the following results:

$SU(6)$ and $SU(7)$ give completely vector-like GUTs.

SU(5) gives a GUT with the maximal number of $\lambda = -\frac{1}{2}$ states and leads precisely to three $10 \pm \bar{5}$ families, plus a set of states which are real under SU(5).

In order to understand these results, let us write the general SU(N) assignment for a state of helicity λ which belongs to a PCT conjugate supermultiplet

$$\{\lambda_{\max}, R\} = \{\frac{N}{2} - \lambda_{\max}, \bar{R}\} \quad (7)$$

where R is the irreducible SU(N) representation under which $\lambda_{\max}$ transforms.

In the multiplet given by (7) the state of helicity λ transforms according to:

$$[2\lambda_{\max} - 2\lambda] \times R + [2\lambda_{\max} + 2\lambda] \times \bar{R} \quad (8)$$

and the second term is absent if the multiplet is PCT self-conjugate ($\lambda_{\max} = N/4$, $R = \bar{R}$).

It is evident that the only supermultiplets which contain states with $|\lambda| = \frac{1}{2}$ are those with $2 \leq \lambda_{\max} \leq 9/2$ for $N = 8$. More precisely the $\lambda = -\frac{1}{2}$ content is

$$\begin{aligned} \{9/2, R\} &= \bar{R} \\ \{4, R\} &= \bar{8} \times \bar{R} \\ \{7/2, R\} &= \bar{28} \times \bar{R} + R \\ \{3, R\} &= \bar{56} \times \bar{R} + 8 \times R \\ \{5/2, R\} &= 70 \times \bar{R} + 28 \times R \\ \{1/2, R\} &= 56 \times (R + \bar{R}) \end{aligned} \quad (9)$$

The multiplet considered by EGMZ corresponds to $\{5/2, \bar{8}\}$ and contains the left-handed $-\frac{1}{2}$ fermions in the

$$70 \times 8 + 28 \times \bar{8} = 504 + \bar{216} + (56 + 8) \quad (10)$$

representation of SU(8).

The EGZ solution for $SU(5) \subset SU(8)$ which fulfils criteria (a), (b), and (c) is obtained by splitting the $504 + \overline{216}$ representations in a real part

$$45 + \overline{45} + 4(24) + 9(1) + 9(10 + \overline{10}) + 6(5 + \overline{5}) \quad (11)$$

and a complex part

$$2(45) + 3(40) + 15 + 6(5) + 3(10) \quad (12)$$

and by substituting $3(10)$ for $3(5)$ in (11). Then the anomaly free set of states is:

$$3(10 + \overline{5}) + SU(5)\text{-real} \quad (13)$$

where $SU(5)$ -real is:

$$45 + \overline{45} + 4(24) + 9(1) + 9(10 + \overline{10}) + 3(5 + \overline{5}) \quad (14)$$

We stress that in order to obtain this result it is crucial to neglect the trace part $\overline{56} + 8$ in (10). In fact this trace under $SU(5)$ decomposes as:

$$\overline{56} + 8 = (10 + \overline{10} + 5 + \overline{5} + 4(1))(\text{real}) + (2(\overline{10}) + 2(\overline{5}))(\text{complex}) \quad (15)$$

and the same substitution as before would give only one $(10 + \overline{5})$ family. There is no sensible reason why the spin $\frac{1}{2}$ states belonging to traces should be unbound. They are just there. However, it has been suggested¹⁷⁾ that there is a way to cure this point. Namely, to add the spin $\frac{1}{2}$ left-handed fermions which belong to the graviton multiplet and which transform as the 56-dimensional representation of $SU(8)$. Then these fermions can be paired with the trace part of the multiplet in (5) to give rise again to an anomaly-free $SU(5)$ spectrum with three $10 + \overline{5}$ families. Preliminary calculations¹⁸⁾ on the dynamics of $N = 8$ supergravity are in support of the idea that the states belonging to the graviton multiplet should participate to the effective $SU(8)$ gauge theory and that traces should not be neglected.

EGZ gave another argument to select $SU(5)$ as the effective GUT below the Planck scale. They considered the $\lambda = 0$ states as high fields and noticed that the scalar representation 420 contains an $SU(5)$ singlet but not $SU(6)$ or $SU(7)$ singlets.

However, there are two problems associated with this analysis¹⁷⁾: there is no SU(8) gauge theory to start with because SU(8) is spoiled by anomalies. The stability group of the 420-dimensional representation of SU(8) is not SU(5) but SU(5) × SU(2) × U(1). The trace (28) would reduce it to SU(5) × U(1). The scalars are not sufficient to break U(1) which remains an unbroken gauge symmetry of the theory, unless some other mechanism is advocated. It is easy to see that there is no complex set of states which is anomaly-free with respect to the full SU(5) × U(1) unbroken subgroup of SU(8).

In order to show this, it is sufficient to give the SU(5) × SU(2) × U(1)_Q decomposition of the 504+216 representations of SU(8)¹⁷⁾:

$$\begin{aligned}
 216 = & (1,2)_{Q=0} + 2(1,1)_{Q=0} + (5,1)_{Q=1} + (\bar{5},1)_{Q=-1} + (1,3)_{Q=0} + (\bar{5},3)_{Q=-1} \\
 & + (\bar{10},2)_{Q=-2} + (\bar{5},2)_{Q=4} + (24,2)_{Q=0} + (1,2)_{Q=0} + (5,2)_{Q=6} + (\bar{5},2)_{Q=-6} \\
 & + (45,1)_{Q=-1} + (24,1)_{Q=5} + (\bar{10},1)_{Q=-7} + (\bar{5},1)_{Q=-1} \\
 \bar{504} = & (\bar{45},1)_{Q=1} + (24,1)_{Q=-5} + (10,1)_{Q=7} + (5,1)_{Q=1} + (\bar{10},2)_{Q=-2} \\
 & + (10,2)_{Q=2} + (\bar{5},2)_{Q=4} + (5,2)_{Q=-4} + (10,3+1)_{Q=-3} + (\bar{10},3+1)_{Q=3} \\
 & + (\bar{10},2)_{Q=8} + (\bar{40},2)_{Q=2} + (\bar{45},2)_{Q=-4} + (10,2)_{Q=2} + (\bar{40},1)_{Q=-3} \\
 & + (\bar{10},1)_{Q=3} + (\bar{15},1)_{Q=3} + (\bar{5},1)_{Q=9}
 \end{aligned} \tag{16}$$

From (16) one can read out the U(1)_Q charge of the 10s and 5s representations to be

$$\begin{aligned}
 216 = & (\bar{10},2)_{Q=-2} + (\bar{10},1)_{Q=-7} + (5,1)_{Q=1} + (5,2)_{Q=6} \\
 & + (\bar{5},1)_{Q=-1} + (\bar{5},3)_{Q=-1} + (\bar{5},2)_{Q=4} + (\bar{5},2)_{Q=-6} + (\bar{5},1)_{Q=-1} \\
 \bar{504} = & (10,1)_{Q=7} + (10,2)_{Q=2} + (10,3+1)_{Q=-3} + (10,2)_{Q=2} \\
 & + (\bar{10},2)_{Q=-2} + (\bar{10},3+1)_{Q=3} + (\bar{10},2)_{Q=8} + (\bar{10},1)_{Q=3} \\
 & + (5,1)_{Q=1} + (5,2)_{Q=-4} + (\bar{5},2)_{Q=4} + (\bar{5},1)_{Q=9}
 \end{aligned} \tag{17}$$

Because of the previously mentioned difficulties, we tried to formulate an alternative scheme¹⁷⁾ by relaxing the assumption that all low-energy states should belong to the same supermultiplet which contains the $\lambda = 1$ states as given by (5). Our assumptions for extracting an effective GUT from N-extended supergravity are:

- i) Non-perturbative dynamics of N-extended supergravity should be responsible for an effective renormalizable SU(N) GUT ($N = 5, 6, 7, 8$).
- ii) $\lambda = -\frac{1}{2}$ fermions should belong to a complex anomaly-free SU(N) representation R obtained from a combination of $\{\lambda_{\max}, [k]_N\}$ massless supermultiplets.
- iii) When decomposed with respect to SU(5), R should contain $n(10+\bar{5}) + \text{real}$ ($n =$ number of family generations).
- iv) $\lambda \geq 1$ states, except the $\lambda = 1$ states in the adjoint representation, should not participate in the effective GUT because of the Veltman decoupling conjecture.

Hypothesis (ii) is connected to the fact that we consider the quantum number assignment of two-particle bound states.

Since for each N we must only consider $\lambda_{\max} \leq N+1/2$, only a finite number of supermultiplets are available in order to satisfy the anomaly condition. Incidentally, the absence of SU(N) anomaly, for $N \leq 8$, is not only sufficient but also necessary, in order to have SU(5) anomaly free representations for the $\lambda = -\frac{1}{2}$ particle states.

If we don't pay attention to the quantum number assignments of the preonic fields, we find that there are precisely three sets of different combinations of $N = 8$ supermultiplets which give rise to the anomaly-free SU(8) family

$$56 + 2\bar{8} + \bar{8} \supset 10 + \bar{5} + \text{real} \quad (18)$$

This is uniquely obtained by a combination of supermultiplets, each occurring at most once, as follows:

$$\begin{aligned} & \{9/2, \overline{56} + 28 + 8\} \\ & \{2 + 4 + 7/2, 1\} \\ & \{7/2 + 5/2, 56 + \overline{28} + 8\} + \{2 + 4 + 7/2, 70\} \end{aligned} \quad (19)$$

(15 supermultiplets)

In $N = 8$ supergravity some of these supermultiplets seem to be excluded because fermions have an odd number of $SU(8)$ indices, so any composite supermultiplet must have an even or odd number of $SU(8)$ indices depending whether $\lambda_{\max}$ is integer or half-integer.

According to this selection rule the $56 + \overline{28} + 8$ family is excluded and the only alternative with single-column Young-tableaux is:

$$56 + 5(\overline{8}) > 2(10 + \overline{5}) + 2eal \quad (20)$$

Of course this $SU(8)$ anomaly-free combination requires supermultiplets with higher multiplicities, since from the multiplets $\{\lambda_{\max}, [K]_N\}$ we can, at most, get three $\overline{8}$ representations if repetitions are excluded.

Interestingly enough, the following general theorem can be proved:

If only one and two-columns Young-tableaux for the $\lambda = -\frac{1}{2}$ families are considered, there are only three independent combinations of $SU(N)$ anomaly-free representations without exotics. They always give rise to an even number of $(10+5)$ $SU(5)$ families. It is conjectured that these combinations could arise from massless bound states. The $SU(N)$ representation content of bound states can be analysed in terms of massive representations¹⁹⁾. These bound states have the quantum numbers of multilinear particle states of the preonic fields of the supergravity Lagrangian:

$$\begin{aligned} & \{2, 1\} && \text{preons} \\ & \{2, 1\} \times \{2, 1\} && (J_{\max} = \text{singlet}) \text{ bilinear} \\ & \vdots && \\ & \{2, 1\} \times \{2, 1\} \dots \times \{2, 1\} && \text{n-multilinear} \end{aligned} \quad (21)$$

($J_{\max}$, n-2 columns Young-tableaux)

It appears clear that in order to have a control on the massless bound states, a dynamical set-up must be realized. Recently, a preliminary analysis of massless bound states in terms of Regge poles of $N = 8$ supergravity has been given¹⁸⁾. A possible scenario is that bound states do not form finite multiplets but rather they belong to an infinite-dimensional irreducible unitary representation of E_7 ¹⁶⁾. Recently a general procedure²⁰⁾ to build a certain series of irreducible unitary infinite-dimensional representations (discrete series) of the non-compact groups which occur in extended-supergravities has been invented.

We believe that exciting developments and more insight the complexity of the dynamics of extended supergravity theories will be achieved in the near future.

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