Lawrence Berkeley Laboratory
UNIVERSITY OF CALIFORNIA

Physics, Computer Science & Mathematics Division

Presented at the International Symposium in Memory of Weiner Heisenberg, "Unified Theories of Elementary Particles: Critical Assessment and Prospects", Munich, W. Germany, July 16-21, 1981

AN APPROACH TO THE UNIFICATION OF ELEMENTARY PARTICLE INTERACTIONS

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September 1981

Prepared for the U.S. Department of Energy under Contract W-7405-ENG-48

September 1981 LBL-13371

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I. INTRODUCTION

Grand unified theories[1,2,3] (GUTs) have met with a few rather impressive successes. First, they are experimentally viable in that the assumption that the three coupling constants of the strong and electroweak gauge theories become equal at a common renormalization scale is consistant[4,5,3] with the measured value of the weak angle, and under more restrictive assumptions,[2] with the b-quark to τ-lepton mass ratio.[6,5,7] Secondly they predict the instability of matter, giving an estimated nucleon life-time that should be measurable in the near future, and, for the first time, the possibility of understanding the cosmological predominance of matter over anti-matter. In addition, these theories ensure the quantization of electric charge.

While possessing many attractive features, grand unified theories are, in the view of most theorists, clearly incomplete. The first obvious defect is the proliferation of arbitrary parameters. Among these are: the Yukawa couplings which determine the fermion mass spectrum and flavor mixing in weak decays; the scalar self-couplings which determine the pattern of symmetry breaking which in turn determines the masses of vector mesons associated with broken gauge symmetries. Both the Yukawa and scalar couplings may play a role in the CP violating parameters of the resultant broken-symmetric theory. A unified gauge theory as such possesses no criterion for fixing its fermion and scalar content, nor for the initial choice of gauge group. The second notorious difficulty is the infamous "gauge hierarchy problem"; this is really a specific case of the above-mentioned arbitrariness, but it is particularly acute in that it entails the understanding of a ratio of mass scales which differ by many orders of magnitude, or -equivalently- by many powers of the coupling constant. Finally the force of gravity is not included in our present picture of unified interactions. Our theories extrapolate in energy to about 13 orders of magnitude beyond presently observed energies while ignoring quantum gravitational effects which should become significant at an energy scale only four orders of magnitude larger than that presently accepted as the "unification energy".

It is clear that we need physical criteria which lie outside the scope of gauge theories in order to further restrict our model building. During the last decade the underlying criterion has been renormalizability. The reason why the scalar and spin 1/2 content and coupling constants remain arbitrary is that any interactions among

* This work was supported by the Director, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract No. W-7405-ENG-48.

these fields of dimension ≤ 4 gives a renormalizable theory. On the other hand we know how to construct a renormalizable theory of spin 1 fields only if it encorporates an exact or spontaneously broken gauge symmetry. This fixes the spectrum of vector fields — their multiplicity is given by the adjoint representation of the gauge group — and determines their self couplings in terms of a single arbitrary constant. Once the spin 1/2 and scalar content has been specified, their couplings to vector fields are also determined in terms of the same constant. For higher spins there is no known renormalizable theory, but supergravity[8] offers the promise of a highly convergent theory for spin ≤ 2. For spin > 2 one does not even know how to write down a field theory.

If we wish to take a lesson from recent history, we may note the following. Physical observation, namely, the existence of spin 1 charged weak currents, together with the criterion of renormalizability led to the construction[9] of the experimentally successful electroweak gauge theory. An analogy might be that physical observation namely, the existence of gravitational interactions, requiring spin 2 in a quantum formulation, together with the criterion of a sufficiently convergent theory to allow for the calculation of S-matrix elements in terms of a limited number of input parameters (only the gravitational constant κ?) would lead uniquely to supergravity as the theory of fundamental interactions. Extended supergravity, which embeds internal symmetries as well as supersymmetries,[10] may determine uniquely the elementary particle content of the theory.

The above conjecture leads naturally to the question as to whether the gauge theories which apparently describe well observed phenomena can be embedded into an extended supergravity theory.[11] Extended supersymmetry is characterized by a number N which specifies both the number of independent supersymmetry transformations and the number of degrees of freedom associated with internal symmetries. For a massless supermultiplet, the helicities λ run over the range [12]

$$J - \frac{N}{2} \leq \lambda \leq J \ . \qquad (1)$$

If we restrict ourselves to known field theories we must impose

$$|\lambda| \leq 2 \quad \text{or} \quad N \leq 8 \ . \qquad (2)$$

This constraint restricts the possible supergravity theories to those of Table 1. A theory which is symmetric under N supersymmetries possesses[13] a rigid (usually called global) U(N) symmetry except for N = 8, in which case the basic (unique for $|\lambda| \leq 2$) supermultiplet is self-conjugate, and does not admit a U(1) symmetry (N = 7 and N = 8 theories have the same spectra and are believed to be the same theory). However we are interested in gauge symmetries. In conventional theories these require vector fields in the adjoint representation. Within this framework supergravity theories allow at most the gauging of an orthogonal group SO(N), $N \leq 8$, which is not large enough to

TABLE I. Possible supergravity theories

N	Rigid symmetry	Number of vector fields	Possible gauge symmetry
≤ 5	U(N)	$\frac{1}{2}N(N-1)$	SO(N)
6	U(6)	16	SO(6) × U(1)
7,8	SU(8)	28	SO(8)

accomodate phenomenologically acceptable grand unified gauge theories for which the minimal gauge group is SU(5). In fact SO(8) does not even contain the "observed" broken gauge group $SU(3)_c \times SU(2)_L \times U(1)$ which is believed to describe the strong and electroweak interactions. On the other hand, it does contain[11] the group $SU(3)_c \times U(1)_{em}$ of the presumably exact strong and electromagnetic gauge symmetries (along with an additional U(1) factor) and a possible point of view would be that only the exact gauge symmetries are described by fundamental gauge couplings while the others are somehow dynamically generated. This of course runs counter to the whole idea of grand unification and renders the above mentioned successes fortuitous. In addition one finds[11] that the observed quarks and leptons cannot all be accomodated in the spectrum of elementary particles, which is unique for N = 8 supergravity.

One is thus led to the idea of a composite structure for at least some fields of the conventional theories. Once this idea is accepted we shall see that not only is the allowed particle spectrum enlarged, but the possible gauge group is also enlarged in such a way that the conventional grand unification approach remains intact, except that the "observed" gauge interactions emerge as the effective couplings of composite states, and the grand unification group is no longer arbitrary.

II. HIDDEN SYMMETRIES OF SUPERGRAVITY

The idea of embedding observed gauge symmetries within an extended supergravity theory was given renewed impetus by the analysis of Cremmer and Julia[14] who showed that extended supergravity theories for $4 \leq N \leq 8$ possess an invariance under a group which is the direct product of a rigid non-compact group $\mathcal{G}$ with a local compact group $\mathcal{K}$, where $\mathcal{K}$ is U(N) for N = 4, 5, 6 and SU(8) for N = 7, 8. The scalar fields of the theory are valued on the coset space $\mathcal{G}/\mathcal{K}$. For N = 8 supergravity the invariance group is

$$\mathcal{G} \times \mathcal{K} = E_{7(+7)} \times SU(8) \quad . \tag{2}$$

In a manifestly gauge invariant formulation, there are 133 scalars, corresponding to the adjoint representation of E_7; these decompose into 63 (adjoint of SU(8)) which can be removed by a gauge transformation and 70 (coset space $\mathcal{G}/\mathcal{K}$) which are the physical scalars of the "special" or "physical" gauge. Expressed in terms of physical fields alone, with only 70 scalars, the theory has no manifest SU(8) gauge invariance but only a rigid SU(8) invariance, corresponding to the compact part of $E_{7(+7)}$, while the non-compact part of E_7 is realized non-linearly with parameters which depend on the scalar fields. The full $E_7 \times SU(8)$ invariance becomes manifest in a general gauge where the 63 unphysical scalars are present; gauge invariance is implemented in the usual way through the covariant derivative

$$D_\mu = \partial_\mu - Q_\mu \tag{3}$$

where Q_μ represents a set of vector fields transforming as an SU(8) adjoint, for which, however, there is no kinetic energy term. The equations of motion then determine the Q_μ as composite (bilinear and higher order) operators in terms of the elementary fields of the N = 8 supermultiplet. At first sight SU(8) gauge invariance appears to be an artifact of the general gauge formulation which disappears when the theory is expressed in terms of physical particles. However, Cremmer and Julia conjectured that the two-point function for Q_μ developes a pole at the origin, resulting in massless spin-one states which might be identified with the gauge vectors of the "observed" theory. They supported their conjecture by analogy with CP^{n-1} models in two dimensions where a similar effect has been found[15] to occur in the 1/n expansion.

An immediate consequence of the Cremmer-Julia conjecture is that the grand unified theory (excluding generation unification) is necessarily the "minimal" one, i.e. the Georgi-Glashow[2] SU(5). The larger GUTs which are phenomenologically acceptable such as SO(10) and (compact) E_6 cannot be embedded in SU(8).

A second consequence is that, since the Q_μ necessarily belong to a supermultiplet of fields, one is naturally led to the EGMZ conjecture[16, 17] that the superpartners of the ϱ_μ also develope poles at the origin, a conjecture which is similarly supported by analogy with supersymmetric CP^n models.[18] In this way we would obtain a set of composite, massless states which should include all the fermions, vectors and scalars of the "observed" gauge theory. In this framework the Lagrangian of the "observed" theory, by which we mean the fully "grand" unified theory (i.e. SU(5)), is viewed as an effective Lagrangian describing the interactions of composite states which can be treated as local interactions of elementary fields at energies well below the Planck mass m_P, or equivalently over distances large compared with the Planck length $\kappa = m_P^{-1}$.

A problem immediately arises. Any N = 8 supermultiplet contains spins at least as large as 2, while a renormalizable field theory can have maximum spin 1. In order for an effective renormalizable theory to emerge at the grand unification mass $m_{GU} \sim 10^{-4} m_P$, supersymmetry must be broken at a scale $\gg m_{GU}$, perhaps already at m_P. For this we appeal to some unknown dynamical mechanism, and to "Veltman's theorem":[19]

> Given a set of bound states of radius r and mass $m \ll r^{-1}$, their couplings at energy scales $E \ll r^{-1}$ can be described by an effective local perturbation theory only if they are renormalizable.

This is not really a theorem,[20] but rather a statement of self consistancy. A perturbative treatment of non-renormalizable interactions requires the introduction of a cut-off Λ for which the only scale available is $\Lambda \sim r^{-1} \sim \kappa^{-1}$ in our case. It could be this highly divergent behavior which triggers the breaking of supersymmetry in such a way that a sensible low mass effective theory remains. This seems to be the conslusion that one is inevitably led to if N = 8 supergravity is the underlying theory, since experiment tells us that quarks and leptons, at least some of which are necessarily bound states with $m \ll \kappa^{-1}$, have interactions which are successfully described by a local, renormalizable perturbation theory at energy scales $E \ll m_P$.

Accepting the above set of conjectures we must address ourselves to two major questions.

1. What good is supersymmetry if it is so badly broken? For one thing it may determine[16] the particle spectrum just as the hadron spectrum is determined in terms of flavor SU(F) which is badly broken by quark mass differences. Secondly it may relate [21] the so-far unconstrained parameters of the GUT, just as badly broken SU(5) relates otherwise unconstrained parameters, such as $\sin^2\theta_w$ and m_b/m_τ, of the "low energy" $SU(3)_c \times SU(2)_L \times U(1)$ gauge theory. It is also possible that part of the supersymmetry remains unbroken down to a much lower energy scale, a possibility which could be relevant to the gauge hierarchy problem.[22]

2. What happened to all the other states of the bound supermultiplet? Possibilities include:[23] they acquire masses $O(m_P)$; their couplings are suppressed by powers of κ; they never bound in the first place.

In the remainder of this talk I will review the progress, or lack of it, on confronting these issues.

III. THE SPECTRUM OF COMPOSITE STATES

The EGMZ hypothesis[16] is that the full particle content of the grand unified theory should be found among the states of the irreducible massless supermultiplet which includes the zero-mass-shell projection of the composite SU(N) vector operators $\mathcal{Q}_\mu$ of $4 \leq N \leq 8$ supergravity. Various arguments[16,24,25] suggest that the appropriate multiplet has the structure

$$(-3/2)^A,\ (-1)^A_B,\ (-1/2)^A_{[BC]},\ (0)^A_{[BCD]},\ (+1/2)^A_{[BCDE]}\ \ldots(5/2)^A \tag{4}$$

+ T.C.P. conjugate states,

where A, B,... = 1, ..., N are SU(N) indices, brackets denote total antisymmetrization, and helicity is indicated in parenthesis.[12] The first remark to be made is that the spectrum (4) contains the usual three generations of fermions only for N = 8. The helicity and SU(8) content for this case is displayed in Table 2 where for $-1 \leq \lambda \leq 2$ the smaller representation is obtained by saturating the upper index with one lower index in the tensors (4), and the larger representation is the traceless part. The states enclosed by a dashed line correspond to the possible

TABLE II. Spectrum of composite states for N = 8 (TCP conjugates understood)

helicity	- 3/2	- 1	- 1/2	0	1/2	1	3/2	2	5/2
SU(8)	$\overline{8}$	63	216	420	$\overline{504}$	$\overline{378}$	$\overline{168}$	$\overline{36}$	$\overline{8}$
content		1	8	28	56	70	$\overline{56}$	$\overline{28}$	

content of a renormalizable gauge theory. However it must be further restricted so that the surviving fermion states form an anomaly-free set with respect to the surviving gauge group. This immediately rules out SU(8) as the effective renormalizable gauge theory. If we further impose the observational constraint that left and right handed fermions transform according to inequivalent (chiral) representations of the GUT, but equivalent (vector-like) representations under $SU(3)_c \times U(1)_{em}$ we find[21] that SU(6) and SU(7) are also eliminated as candidate effective theories. Since SU(5) is phenomenologically the minimal GUT group we are led uniquely to this choice. We note in addition[21] that if the 420 of scalars acquires a vacuum expectation value of order m_P, the natural scale of the theory, the largest surviving simple gauge group is again SU(5). Finally, we observe that if (why? in analogy with the retention of only the 63 of spin 1?) we retain only the $216 + \overline{504}$ of left-handed ($\lambda = -1/2$) fermions, we find that the maximal subset which is anomaly free under SU(5) is

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

The $3(\overline{5} + 10)$ can then be identified with the "observed" three generations of light

fermions, while the remaining states can acquire SU(5) invariant Dirac or Majorana masses which are expected to be of order m_{GU} or larger. In addition to the fermions (5) and their right-handed charge conjugate states the particle content of the effective GUT will include the 24 gauge vectors of SU(5) and some (or all) of the scalars of Table II. We note that the scalar content is sufficiently rich to satisfy any phenomenological needs of conventional SU(5).

IV. SCENARIOS FOR THE MISSING STATES

In this section I shall elaborate somewhat on the possibilities mentioned above for disposing of those states in Table II which cannot be embedded in the effective renormalizable GUT.

<u>They acquire masses $O(m_p)$.</u> In the present picture we have the following pattern of gauge symmetry breaking

$$SU(8) \xrightarrow[10^{19}GeV]{} SU(5) \xrightarrow[10^{15}GeV]{} SU(3)_c \times SU(2)_L \times U(1) \xrightarrow[10^{2}GeV]{} SU(3)_c \times U(1)_{em} \qquad (6)$$

where the first step at $m_p = 10^{19}$ GeV is attributed to dynamical effects and the subsequent breaking of SU(5) proceeds according to the conventional Higgs mechanism, with the Higgs scalars being suitable composites among those of Table II. Those masses which are generated at the first step, and which can be $O(m_p)$, must necessarily be SU(5) invariant. It is easy to see[16] that the high spin states of Table II cannot all obtain masses $O(m_p)$ in this way. It has further been found[16] that starting with any irreducible supermultiplet one cannot give masses to high spin states which are invariant even under the unbroken gauge group $SU(3)_c \times U(1)_{em}$. This could be achieved by disregarding conventional SU(5) and reinterpreting the quantum numbers of the basic multiplet so as to make the theory completely vector-like under $SU(3)_c \times U(1)_{em}$, but this cannot be done[16] in a way which reproduces the observed charge spectrum of fermions. There are [25] at least some sets of irreducible supermultiplets which allow SU(8) invariant masses, namely those which are obtained in the zero mass limit from an irreducible massive supermultiplet. However this set allows group invariant masses for <u>all</u> states and no light mass sector would be expected to survive. Attempts to find finite sets of supermultiplets other than the above type have given negative results.[26]

The remaining possibility[24] is that the surviving low mass spectrum is derived from an infinite set of supermultiplets. This idea has some support both from dynamical[27] and group theoretical[23,24] considerations. Recall that N = 8 supergravity is invariant under a non-compact E_7 which is realized non-linearly on the elementary fields (preons):

$$\psi \to F(\phi)\psi \qquad (7)$$

where ψ is a preon state of arbitrary spin and ϕ represents the 70-plet of scalar fields in the special gauge. If we wish to obtain an effective bound state theory where no preon fields appear, E_7 must be realized linearly, in which case the unitary representations are infinite. We may again appeal to an analogy with CP^n models with a non-compact group, where studies[28] suggest that the bound state spectrum forms a linear representation of the full group.

Consider a preon field ψ which transforms according to some representation R

of SU(8). Then under an infinitesimal E_7 transformation of parameter ϵ

$$\delta\psi \sim \epsilon\phi\psi \qquad (8)$$

which transforms according to the reducible SU(8) representation

$$R' = R \times 70 \quad . \qquad (9)$$

Now consider a bound state B which is the zero-mass-shell projection of an operator

$$\mathcal{O} = \psi_1\psi_2 + \text{higher order terms} \qquad (10)$$

$\mathcal{O}$ and therefore B will transform according to some representation

$$R_B \in R_1 \times R_2 \qquad (11)$$

of SU(8). From (9) we find

$$\delta\mathcal{O} \sim \epsilon\psi_1\psi_2\phi + \ldots \qquad (12)$$

so that the state B' obtained from an infinitesimal transformation on B

$$\delta B \sim \epsilon B' \qquad (13)$$

should transform according to a reducible representation

$$R_{B'} \in 70 \times R_B \quad . \qquad (14)$$

A further transformation will generate a representation

$$R_{B''} \in (70 \times 70)_{sym} \times R_B \quad , \qquad (15)$$

and so on.

From a general group theoretical analysis it is easy to see that for a given representation R of SU(8), the set of states

$$R_n = \{70\}^n \times R, \; n = 0, \ldots, \infty, \qquad (16)$$

where $\{70\}^n$ represents a totally symmetrized product of 70's, is a representation of E_7. It is a reducible representation which is infinitely degenerate in each SU(8) representation which it contains. However if the binding dynamics respect both E_7 and supersymmetry, we must look for a set of states which represent the full algebra of symmetry generators. A more detailed analysis[23] shows that the set (16), where in our case the SU(8) representations R for each helicity are those of Table II, does represent the algebra and suggests that no reduction is possible. In this case we obtain[23] the following result: it is possible to give masses to any arbitrary set of the states (16), while keeping any arbitrary set massless in a way which is invariant under a subgroup of SU(8) which can be as large as SU(6). This result also holds if the set (16) is reduced only by limiting the multiplicity of each distinct SU(8) representation appearing in (16) to some finite (non-zero) value.

What is to be concluded from the above analysis? Simply that the scenario (6), with group invariant masses generated at each step, is a self-consistent possibility. However in this framework we have no criterion for determining which set of states survives at low energy, nor even for assuring that the low mass states all belong to our original supermultiplet of Eq. (4) and Table II, except that this is perhaps the most aesthetic hypothesis and the one which is most likely to constrain the effective GUT.

<u>They have couplings $O(\kappa^n)$.</u> S-matrix theorems[29] for threshold behavior imply that a state of helicity λ must couple with a dimensionful parameters

$$G_\lambda \propto m^{|2\lambda|-3} \tag{17}$$

where in our case we expect $m = \kappa^{-1}$. This implies that states of helicity $|\lambda| \geq 5/2$ have couplings $|G| < \kappa^2$ which are weaker than gravity and unobservable. States of helicity $|\lambda| = 2$ could have couplings of gravitational strength but their effect would be incoherent and unobservable for states which are not $SU(3)_c \times U(1)$ invariants; the others can have masses $\gtrsim O(m_W)$. Vector fields which are in a non-self-conjugate SU(8) representation have well defined transformation properties only when expressed in terms of the field strength and its dual, so that their SU(8) invariant couplings must again involve a dimensional constant $G \lesssim \kappa$. These arguments do not directly imply the decoupling of $|\lambda| = 3/2$ states nor of the $|\lambda| = 1/2$ states which are excluded from the anomaly free sector.

An alternative approach[30] is to disregard (as decoupled?) the high spin sector $|\lambda| \geq 1$ (except for the adjoint of vectors) and demand that the full $|\lambda| = 1/2$ spectrum form an anomaly free set under SU(8). Taking account of the selection rule[23] which requires that fermionic bound states of the supergravity preons must have an odd number of SU(8) indices, it is found that an even number of SU(5) generations of light $(\bar{5} + 10)$ fermions will occur, and that the initial bound state spectrum must include a repetition of some supermultiplets (to be interpreted as radial excitations?)

One might also ask whether there are any cosmological constraints limiting the number of decoupled massless helicity states. This turns out[23] not to be the case as long as the number of states which decoupled at the Planck temperature is sufficiently larger than, or smaller than, the number of states remaining in equilibrium, or if the latter number is itself sufficiently large, which is the case for the set (5) and also for smaller sets which one might be led to consider[23] if one wished to retain an unbroken simple supersymmetry for the effective gauge theory.

<u>They never bind.</u> There are arguments[31] based on Lorentz covariance that massless states with helicity $|\lambda| \geq 1$ cannot carry a quantum number [e.g. SU(8), SU(5)]

associated with a local vector current. Fortunately this argument is not applicable[32] in our case, as it would also exclude gauge vectors as bound states. The conserved SU(8) current[32] of N = 8 supergravity is not invariant under abelian gauge transformations

$$A_\mu \to A_\mu + \partial_\mu \lambda$$

on the preon vector fields A_μ. This means that the current can only be specified once a gauge is fixed, and that its Lorentz transformation properties are not well defined.

On the other hand, as discussed above, "Veltman's theorem " implies that supersymmetry must break dynamically at the Planck scale in order that an effective low energy theory may emerge. Perhaps this symmetry breaking is associated with the binding mechanism itself, preventing the formation of those states which would spoil renormalizability.

IV. CONCLUSIONS AND PROSPECTS

Ideally, the assumption that the grand unified theory emerges as an effective theory describing bound states of $N = 8$ supergravity preons should determine the GUT particle spectrum and constrain their couplings. Analysis of the spectrum has led to some possibly encouraging indications. At the least, the particle content in scalars, vectors and fermions needed to reproduce SU(5) phenomenology can be found among the states of the EGMZ multiplet.[16]

As for constraints on couplings, little progress has been made at present. One possibility is that some supersymmetry survives below the Planck mass. Renormalizability requires $|\lambda| \leq 1$, which from (1) implies $N \leq 4$ for the number of surviving supersymmetries. However, the empirical observation that fermions form chiral representations of the gauge group restricts this number to 1. The further study of supersymmetric GUTs, as well as of the non-renormalization properties of supersymmetric theories may prove enlightening, particularly with regard to the "gauge hierarchy" problem.

Electroweak phenomenology requires the Higgs scalar of the Glashow-Weinberg-Salam model[9] to have a mass smaller than about a TeV. On the other hand, when the electroweak model is embedded in a GUT, the scalar mass generally acquires a contribution of order 10^{15} GeV through its coupling to the 24-plet of scalars which generate the break-down of SU(5) to the $SU(3)_c \times SU(2)_L \times U(1)$ of GWS. In order to fit phenomenology, this contribution must be made to vanish by artificially adjusting coupling constant ratios or by cancelling it against equally large explicit mass terms. The required cancellation can indeed be arranged at any order in perturbation theory but it is highly unstable against radiative corrections which for n loops give contributions of order $(\alpha/4\pi)^n \times 10^{15}$ GeV, where $\alpha \cong 1/40$ is the GUT coupling constant.

A natural solution to the problem would occur if there were some symmetry which forced the mass of the GWS Higgs scalar to vanish. In gauge theories gauge invariance forces vectors to be massless and the chiral nature of the gauge couplings requires massless fermions. Since supersymmetry requires fermion-boson degeneracy, a supersymmetric chiral gauge theory would require massless scalars as well. Within this scenario supersymmetry must remain valid at energy scales down to 1 TeV or less. However this more or less obvious way of implementing the requirement $m_H \lesssim$ TeV turns out to have other phenomenological difficulties.

There is another property of supersymmetric theories which is not entirely understood: the so-called non-renormalization theorem[33] which ensures that if cancellations of the type needed to keep the Higgs mass small are imposed on the Lagrangian at the tree level, they will be stable against radiative corrections. SU(5) models[34] of this type yield satisfactory phenomenology if the symmetry breaking arises from terms of dimension two or less with parameter $m^2 \lesssim (1\ \text{TeV})^2$.

Theories of this type have observable consequences for proton decay, and predict

a large spectrum of new states of mass $\lesssim 1$ TeV. One might ask if anything can be salvaged in the event that the supersymmetry breaking scale is much higher (e.g. m_{GU} or m_P). A better understanding of the physics behind the "non-renormalization" properties found in supersymmetric theories could perhaps shed some light on this question.

ACKNOWLEDGMENT

The ideas collected here are based on work in collaboration with John Ellis and Bruno Zumino. I am also happy to acknowledge numerous discussions with Murray Gell-Mann on these issues.

This work was supported by the Director, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract No. W-7405-ENG-48.

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