PHENOMENOLOGICAL SUPERGRAVITY *)

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ABSTRACT

The rich phenomenological consequences of simple supergravity theories are discussed in detail. Simple ($N = 1$) supergravity theories may provide the framework for consistently describing physics from energies close to (but below) the Planck scale ($\sim 10^{19}$ GeV) down to low energies. It is shown that gravitational effects, as contained in supergravity theories, may play, for the first time, a rather fundamental role in particle physics at low as well as at superhigh energies. For example, gravitational induced $SU(5)$ and $SU(2) \times U(1)$ breaking occurs naturally at vastly different energy scales, thus evading the cumbersome gauge hierarchy problem. In this extended, updated, written version of my lecture, I concentrate on the physical consequences of simple ($N = 1$) supergravity theories. After providing a few reasons which make SUPERGRAVITY theories relevant and maybe necessary to physics (Section 1), the physical structure of $N = 1$ supergravity is discussed at some length (Section 2), followed by a detailed discussion of the phenomenological consequences of supergravity theories, both at the grand unified (GUT) and electroweak energy scales (Section 3), and some final remarks (Section 4). Physics, not formalism, is our focus.

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1. WHY SUPERGRAVITY?

Grand unified theories (GUTs), despite all their successes and aesthetical appeal, seem to be problematic¹⁾. Their main defect is the gauge hierarchy problem. It seems very difficult to create, and then keep intact to all orders in perturbation theory, the energy gap of thirteen decades that exists between the electroweak scale $M_W (\sim 100 \text{ GeV})$ and the GUT scale $M_X (\geq 10^{15} \text{ GeV})$:

$$\frac{M_W}{M_X} \leq 10^{-13} \quad (1)$$

Supersymmetry (SUSY), or Fermion-Boson symmetry²⁾, solves at least the technical aspect of this problem automatically³⁾. Given (1) at the tree level, SUSY theories are able to keep it intact to all orders in perturbation theory.

In supersymmetric theories, fermions and bosons are sharing the same multiplet (supermultiplet), and thus new relations between fermion and boson masses, as well as between gauge and scalar self-couplings, are enforced. This is analogous to GUTs where, by putting together quarks and leptons, we get relations between their masses, as well as relations between different gauge coupling constants¹⁾. In the case of SUSY gauge theories, because of their high symmetry, a lot of "miraculous" cancellations take place and render the theory much more ultra-violet convergent than usual gauge theories. A quantitative expression of this behaviour is shown by the so-called non-renormalization theorems⁴⁾. There is wave function renormalization for each chiral and gauge supermultiplet and gauge coupling renormalization

BUT no separate renormalization of masses, scalar and Yukawa couplings is needed. The non-renormalization theorems⁴⁾ go beyond the fact that parameters related by SUSY are expected to be renormalized by the same constant. They say that some parameters are not renormalized at all. THE SET IT AND FORGET IT principle. It is exactly this property of SUSY theories that solves the technical aspect of the gauge hierarchy problem, as mentioned before³⁾. Also in SUSY theories, the technical aspect of the strong CP problem is solved automatically⁵⁾. The parameter θ_{QCD} which characterizes the magnitude of the effect ($\theta_{QCD} < 10^{-9}$ experimentally), if set equal to zero, will stay zero, if SUSY is good, because the fermion mass matrix is not renormalized⁵⁾. When SUSY is spontaneously broken, θ_{QCD} acquires a small, finite, acceptable value⁵⁾. Furthermore, if SUSY is exact, no vacuum renormalization is needed⁶⁾, a property that eventually may shed light on the absence of the cosmological constant.

Clearly, SUSY cannot be an exact symmetry in Nature, since then all members of a supermultiplet will have the same mass, which is experimentally excluded. For example, there is no scalar electron (selectron) with mass of 0.5 MeV! Actually, present experimental limits⁷⁾ put the masses of any charged SUSY particle above 0(20 GeV). SUSY needs to be broken. We may consider spontaneous, or explicit but soft breaking, SB or ESB respectively. Arbitrary ESB is ad hoc and problematic (see below). SB of global SUSY is associated with the existence of a Goldstone fermion, the goldstino, ψ_g . If g_{bf} denotes the coupling of the goldstino ψ_g to a supermultiplet that contains a boson b and a fermion f , then the boson-fermion mass splitting is given by a Goldberger-Treiman type relation:

$$m_b^2 - m_f^2 = g_{bf} M_S^2 \quad (2)$$

where M_S denotes the SUSY breaking scale. A priori, g_{bf} and M_S are arbitrary parameters, but they have to satisfy certain physical constraints: (i) absence of anything "new" below 20 GeV⁷⁾ imposes a lower bound, at least for the charged SUSY partners of the "observed" particles:

$$g_{bf} M_S^2 \geq 0 \left((20 \text{ GeV})^2 \right) \quad (3)$$

(ii) absence of radiative corrections to the mass of the electroweak Higgs (E-Higgs) larger than $O(M_W)$, i.e., solution of the gauge hierarchy problem (technical aspect), imposes an upper bound:

$$g_{bf} M_S^2 \leq 0 \left((M_W)^2 \right) \quad (4)$$

Clearly, the most naïve way to satisfy (3) and (4) is to identify g_{bf} with some gauge coupling and M_S with M_W :

$$\begin{aligned} g_{bf} &\sim g \\ M_S &\sim O(M_W) \end{aligned} \quad (5)$$

This approach has been tried for many years now by Fayet, with rather limited success⁸⁾. The Fayet-type models⁸⁾ suffer from incurable diseases, such as broken colour and/or electromagnetism, and/or Adler-Bell-Jackiw anomalies, and/or dangerous flavour-changing neutral currents, etc.⁹⁾. We need something better. The other way is to diminish g_{bf} with an approximate increase of M_S , such that (3) and (4) are still satisfied. This approach, instigated by the work of Barbieri, Ferrara and myself¹⁰⁾, and followed by several other groups¹¹⁾, seems to be more successful. The idea is very simple¹⁰⁾. One has to find a natural way, so that the effective goldstino coupling to "ordinary" matter (quarks, leptons, E-Higgs, $W^\pm$, Z , γ , g) and their SUSY partners is extremely weak. This tiny effective coupling can be arranged at the tree level¹⁰⁾, or if absent at the tree level, may be induced, after some gymnastics, at some number of loops by radiative corrections¹¹⁾. In any case, in a large class of this type of models, one gets the following relations, as discussed in Ref. 10):

$$\begin{aligned} g_{bf} &\sim \frac{M_W}{M_{Pl}} \\ M_S^2 &\sim M_W M_{Pl} \sim (10^{10} \text{ GeV})^2 \end{aligned} \quad (6)$$

which automatically satisfy the constraints (3) and (4). In other words, despite the fact that the whole theory is characterized by a SUSY breaking scale $M_S \sim 10^{10} \text{ GeV}$, the "low energy" world ($\equiv$ "ordinary" matter plus SUSY partners) is feeling an effective SUSY breaking

$$(M_S)_{eff} \sim O(M_W) \ll M_S \quad (7)$$

Thanks to the "magic" properties⁴⁾ of SUSY theories, this decoupling has been shown to persist to all orders in perturbation theory^{10), 12)}.

Up to now, we have considered global SUSY gauge theories and have shown that the construction of "realistic"([?]), phenomenologically accepted models^{10), 11)}, has enforced upon us Eq. (6). This fact is extremely important because the supergravitational (SUGRA)¹³⁾ coupling between the goldstino and "ordinary" matter ($m \sim M_W$)

$$(g_{bf})_{\text{SUGAR}} \sim \sqrt{G_N} m \simeq \frac{M_W}{M_{Pl}} \quad (8)$$

is exactly of the same order of magnitude as the corresponding non-gravitational coupling, given by Eq. (6). In other words, (super) gravitational effects cannot be neglected anymore^{14),15),16)}.

We may go even further and assume that non-gravitational g_{bf} 's are altogether absent and only those induced by supergravity remain. This is a spectacular idea, because it is not only very exciting that gravity, for the first time, plays a major role in low energy physics, but also because the resulting models are the simplest of all the realistic models of broken SUSY.

If we follow this line of thought, then (8) and (4) set an upper limit¹⁵⁾ on M_S :

$$M_S^2 \leq O(M_W M_{Pl}) \quad (9)$$

while (8) and (3) set a lower limit on M_S :

$$M_S^2 \geq O\left(\left(\frac{90 \text{ GeV}}{M_W}\right)^2 M_W M_{Pl}\right) \quad (10)$$

More or less in SUGAR-type models, M_S is determined uniquely from physics constraints to be

$$M_S^2 \sim O(M_W M_{Pl}) \quad (11)$$

The coincidence between (6) and (8), (11) is rather amazing. Intuitively, one can easily understand the need to move from global to local SUSY or supergravity¹³⁾. Consider the one-loop correction to the E-Higgs mass due to the graviton

$$\delta m_H^2 \sim \frac{\Lambda^4}{M_{Pl}^2} \quad (12)$$

This is catastrophic because the natural scale for the cut-off Λ is M_{Pl} , which, through (12), implies electroweak breaking at the Planck scale! We need either extremely accurate fine tuning or unavoidable cancellation of the graviton (boson) loop by a fermion loop, i.e., supersymmetric gravity, thus SUPERGRAVITY¹³⁾. The needed spin 3/2 fermion sits on the same supermultiplet with the graviton (spin 2) and it is called gravitino. As the graviton may be considered the "gauge boson" of the Poincaré algebra, the gravitino is the

"gauge fermion" of local SUSY. After SB of local SUSY, the gravitino "eats" the goldstino, thus becoming massive¹⁷⁾:

$$M_{3/2} \sim \sqrt{G_N} M_S^2 \sim \frac{M_S^2}{M_{Pl}} \quad (13)$$

which implies, by using (11) 10):

$$M_{3/2} \sim O(M_W) \quad (14)$$

One may argue that, because of the anticommutation relation, which in standard notation²⁾ reads

$$\{Q_\alpha, \bar{Q}_\beta\} = -2(\gamma_\mu)_{\alpha\beta} P_\mu$$

global SUSY plus gravity implies automatically local SUSY or supergravity, so what's new? Well, that is correct, but it was thought for a long time that, as in the case of ordinary gravity, the effects of supergravity, at least for low energy physics, would be negligible. We have proved here that this is not necessarily the case. Only if the Fayet-type models⁸⁾ were correct, would (8) indeed be much smaller than (5) and SUGRA effects would be negligible. Nevertheless, Fayet-type models⁸⁾ do not work⁹⁾, thus supergravity enters the "low energy" physics world in full strength and glory.

2. PHYSICAL STRUCTURE OF SIMPLE (N = 1) SUPERGRAVITY

We are then led to consider local SUSY gauge theories¹⁸⁾. The effective theory below the Planck scale must be¹⁹⁾ N = 1 supergravity. The restriction to N = 1 follows from the apparent left-right asymmetry of the "known" gauge interactions. Since we are dealing with local SUSY, the breaking of SUSY must be spontaneous, not explicit, if Lorentz invariance or unitarity are not to be violated. It is remarkable that the effective theory below M_{Pl} has been uniquely determined¹⁹⁾ to be a spontaneously broken N = 1 local SUSY gauge theory¹⁸⁾.

We start with a reminder of the structure of N = 1 supergravity actions¹⁸⁾ containing gauge and matter fields (if not explicitly stated, we use natural units $k^2 \equiv 8\pi G_N = (8\pi/M_{Pl}^2) \equiv 1/M^2 = 1$)

$$A = \int d^4x d^4\theta E \left\{ \Phi(\varphi, \bar{\varphi} e^{2V}) + \text{Re}[R^{-1}g(\varphi)] + \text{Re}[R^{-1}f_{\alpha\beta}(\varphi) W_\alpha^\alpha \varepsilon^{ab} W_b^\beta] \right\} \quad (15)$$

where E is the superspace determinant, Φ is an arbitrary real function of the chiral superfields ϕ and their complex conjugates $\bar{\phi}$, V is the gauge vector supermultiplet, R is the chiral scalar curvature superfield, g is the chiral superpotential, $f_{\alpha\beta}$ is another chiral function of the chiral superfields ϕ , and W_α^a is a gauge-covariant chiral superfield containing the gauge field strength. In addition to all the obvious general co-ordinate transformations, local supersymmetry and gauge invariance, the action (15) is also invariant¹⁸⁾ under the transformations

$$\begin{aligned} J &\equiv 3\ln(-\frac{1}{3}\Phi) \rightarrow J + K(\phi) + K^*(\bar{\phi}) \\ g(\phi) &\rightarrow e^{K(\phi)} g(\phi) \end{aligned} \quad (16)$$

These transformations can be related to a description of the chiral superfields ϕ as co-ordinates on a Kähler manifold with Kähler potential Φ , and the transformations (16) are known as Kähler gauge transformations²⁰⁾. One particular manifestation of this Kähler gauge symmetry is in the effective scalar potential

$$V = -\exp(-G) \left(3 + G_i' G_j''{}^{i-1} G'^j \right) + (\text{gauge terms}) \quad (17)$$

where

$$G \equiv J - \ln \left(\frac{1}{4} |g|^2 \right) \quad (18)$$

which is clearly invariant under the transformations (16). In general, the action depends on a real function

$$\hat{\Phi} \equiv \Phi / |g(\phi)|^{2/3} \quad (19)$$

and on the chiral function $f_{\alpha\beta}(\phi)$. The most familiar forms of these functions are $J = -\phi\bar{\phi}/2$, giving canonical kinetic energy terms for the chiral superfields, $g(\phi)$ a cubic polynomial giving renormalizable matter interactions of dimension ≤ 4 , and $f_{\alpha\beta} = \delta_{\alpha\beta}$. We expect that more complicated functions will contain terms $O(\phi/M_{Pl})^n$ relative to these canonical leading terms.

Ellis, Tamvakis and myself have suggested¹⁹⁾ interpreting Eq. (15) as an effective action suitable for describing particle interactions at energies $\ll M_{Pl}$ just as chiral $SU(N) \times SU(N)$ Lagrangians were suitable for describing hadronic interactions at energies $\ll 1$ GeV. In much the same way as we know that physics gets complicated at $E = 1$ GeV, with many new hadronic degrees of freedom having masses of this order, we also expect many new "elementary particles" to exist with masses $O(M_{Pl})$. It may well be that all the known light "elementary particles", as well as these

heavy ones, are actually composite, and that at energies $\gg M_{Pl}$ a simple preonic picture will emerge, analogously to the economical description of high-energy hadronic interactions in terms of quarks and gluons. It may even be that these preonic constituents are themselves ingredients in an extended supergravity theory²¹⁾. But let us ignore these speculations for the moment and return to our pedestrian phenomenological interpretation of the action (15).

The well-known rules of phenomenological Lagrangians²²⁾ are that one should write down all possible interactions consistent with the conjectured symmetries [e.g., chiral $SU(2) \times SU(2)$], and only place absolute belief in predictions which are independent of the general form of the Lagrangian (e.g., $\pi\pi$ scattering lengths). These are the reliable results which could also be obtained using current algebra arguments. It does not make sense to calculate strong interaction radiative corrections (read: supergravity loop corrections) to these unimpeachable predictions: these are ambiguous until we know what happens at the 1 GeV scale (read: M_{Pl}), and our ignorance can be subsumed in the general form of the phenomenological Lagrangian, in which any and all possible terms are present a priori (read: non-trivial J , non-polynomial g and $f_{\alpha\beta}$). On the other hand, non-strong interaction radiative corrections can often be computed meaningfully (e.g., the $\pi^+-\pi^0$ mass difference, large numbers of pseudo-Goldstone boson masses in extended technicolour theories). Similarly, it makes sense to compute matter interaction (gauge, Yukawa, Higgs) corrections to the tree-level predictions of the effective action (15).

Since the supergravity action is non-renormalizable, and since both the ϕ and $f_{\alpha\beta}$ terms in the action (15) have a $\int d^4\theta$ form, we expect general variants of them to be generated by loop corrections. Presumably, radiative corrections maintain the essential geometry of the Kähler manifold²⁰⁾. Therefore, we expect loop corrections to fall into the class of Kähler gauge transformations (16). The only analogous transformation allowed in a conventional renormalizable theory is $K = \text{constant}$, corresponding to a wave function renormalization. In our case, more general gauge functions $K(\phi)$ might appear.

In $N = 1$ SB local SUSY gauge models, called for abbreviation supergravity or SUGRA models, one usually distinguishes two sectors: (i) the "observable" sector containing quarks, leptons, Higgs and gauge bosons of electroweak and GUT types, as well as their SUSY partners; (ii) the "hidden" sector containing at least the goldstino and its SUSY associate. The "observable" and "hidden" sectors both couple to supergravity, but not directly to each other. In other words, ignoring supergravity, physics in the "observable" sector would be completely supersymmetric. A realization of this programme occurs, for example, by splitting the superpotential of the theory into the sum of two terms

$$g(\varphi_i) = h(z_i) + f(x_i) \quad (20)$$

where x_i are the "observable" fields and z_i the "hidden" fields. The f -part of the superpotential has to do with the "observed" sector and contains most of the physics, while the h -part of the superpotential has to do with SB of local SUSY and the vanishing of the cosmological constant. The scalar fields of the hidden sector typically have vev's of M_{Pl} which cause SUSY to be SB at a scale M_S determined by the parameters in the h -part. We may choose these parameters to be such that (11), or equivalently (14), are satisfied, as well as cancelling the cosmological constant at the same time ("superHiggs effect")¹⁸. A celebrated example is the Polonyi potential²³:

$$h(z) = M_S^2 (z + B) \quad (21)$$

which, in the absence of other fields, and for $B = (2-\sqrt{3})M$ and $\langle Z \rangle = (\sqrt{3}-1)M$, [$M \equiv (M_{Pl}/\sqrt{8\pi}) \simeq 2.5 \cdot 10^{18}$ GeV being the appropriate supergravity scale (superPlanck mass)], implies SB of local SUSY at M_S and vanishing cosmological constant. The communication between the two sectors is mediated by the auxiliary fields of the SUGRA multiplet and takes place at tree-level in a model-independent way. Exchange of gravitons or gravitinos plays no role at this level. The elimination of these auxiliary fields produces non-renormalizable interactions between the two sectors. These non-renormalizable interactions include the ones between the hidden fields z_i , usually taken to be neutral under all gauge symmetries, and gauge fields.

All these effects have been summed into an effective scalar potential¹⁸) [(17), but with canonical kinetic energy terms for the chiral superfields].

$$V(\varphi) = \exp\left(\frac{1}{M^2} \sum_i \varphi_i \bar{\varphi}_i\right) \left\{ \sum_i \left| \frac{\partial g(\varphi)}{\partial \varphi_i} + \frac{\bar{\varphi}_i g(\varphi)}{M^2} \right|^2 - \frac{3 |g(\varphi)|^2}{M^2} \right\} + \frac{1}{2} \sum_a D_a^2(\varphi, \bar{\varphi}) \quad (22)$$

This mess involves the scalar fields ϕ_i , the superpotential g and the D -terms, which have their usual global SUSY form. Incidentally, by comparing (22) with the global SUSY potential

$$V(\phi) = \sum_i \left| \frac{\partial g(\phi)}{\partial \phi_i} \right|^2 + \frac{1}{2} \sum_a D_a^2(\phi, \bar{\phi})$$

we notice the richer structure of the SUGRA potential [(22)] and thus we anticipate that more physics will come out from it. When the scalar components of the hidden fields are replaced by their vev's, all superheavy fields are integrated out, and expanding in powers of $1/M$, the resulting effective theory, just below M_X , for the light observable fields, will contain both the usual SUSY terms and a soft SUSY breaking piece^{15),24)-26)}:

$$\begin{aligned} \mathcal{L}_{\text{SOFT}} = & \sum_i m_i^2 |\chi_i|^2 + m_{3/2} \sum_n (A_n f_n + \text{h.c.}) \\ & - \sum_{\alpha=1}^3 \left(\frac{1}{2} M_\alpha \lambda_\alpha \lambda_\alpha + \text{h.c.} \right) \end{aligned} \quad (23)$$

where λ_α is a gauge fermion [gluino ($\alpha=3$), wino ($\alpha=2$) or bino ($\alpha=1$)], f_n is any term in f and χ_i is any scalar (Higgs, squark or slepton). The A_n 's²⁵⁾ are expected to be of order one, while m_i and M_α should be in general of order $m_{3/2}$ ^{15),19),27)}. Actually, when corrections²⁶⁾ from integrating out superheavy fields are ignored, the following relations hold¹⁵⁾:

$$m_i^2 = m_{3/2}^2 \quad (24)$$

for every scalar field, and^{24),25)}

$$A_n = A - 3 + d_n \quad (25)$$

where A is a universal number²⁵⁾ (assumed real) and d_n is the number of fields multiplied together in term n . If we imagine that the low energy theory is embedded in a GUT model at some GUT scale below the Planck mass, then all gaugino masses are equal at the GUT scale, so that only one single parameter M_0 is needed [M_0 in principle may be of $O(m_{3/2})$]:

$$M_\alpha(M_X) = M_0, \quad \alpha=1,2,3 \quad (26)$$

while at lower energies M_α evolves in a manner identical for the gauge couplings:

$$\frac{M_\alpha(\mu)}{M_0} = \frac{\tilde{\alpha}_\alpha(\mu)}{\alpha_G}, \quad \alpha=1,2,3 \quad (26')$$

with $\tilde{\alpha}_\alpha, \alpha_G$ the usual $SU(3), SU(2), U(1),$ GUT fine structure constants.

In the following, we shall assume that corrections²⁶⁾ from integrating out superheavy fields are not very important and proceed with the simplified soft SUSY breaking piece, just below M_X ,

$$\begin{aligned} \mathcal{L}_{\text{SOFT}} = & m_{3/2}^2 \sum_i |X_i|^2 + m_{3/2} \sum_n (A-3+d_n) f_n + h.c.] \\ & - \frac{1}{2} M_0 \sum_\alpha \lambda_\alpha \lambda_\alpha + h.c. \end{aligned} \quad (27)$$

The soft operators in (23) or (27) are determined by the couplings of the hidden fields and by radiative corrections at the Planck scale, including those due to gravity. In spite of this, the very interesting thing is that sometimes it is possible to give the form of these operators without making detailed assumptions about either the hidden sector or the effects of gravitational radiative corrections. For example, in (27), the form of the hidden sector enters only through the three parameters $m_{3/2}, A$ and M_0 . The non-renormalizable interactions of the hidden fields with gauge fields, discussed before, will generally lead to soft Majorana masses, M_0 , for the gauginos [$f_{\alpha\beta} \neq \delta_{\alpha\beta}$ in (15)], while the non-renormalizable interactions between the "observable" and "hidden" sectors, commented before, will produce the soft operators in the scalar potential [first two terms in (23) or (24)].

In a nutshell, the SB of SUSY in SUSY gauge theories coupled to $N = 1$ SUGRA leads to an effective theory below the Planck scale in which global SUSY is explicitly broken by a constrained set of soft operators, at an effective scale:

$$(M_S)_{\text{eff}} \simeq m_{3/2} \quad (28)$$

As discussed before, physics constraints [see (7)] impose the condition $m_{3/2} \sim O(M_W)$ [see (14)], which creates a new hierarchy problem. Since we are dealing with gravitational phenomena, naïvely, the natural mass scale for the gravitino is M_{Pl} and not M_W . I call this the supergravity hierarchy problem, or the SUGAR hierarchy problem. On the other hand, it should be emphasized that the automatic soft breaking of global SUSY that is provided in the SUGAR framework not only splits "low energy" supermultiplets in the right way [all scalars and gauginos getting masses $O(m_{3/2})$], BUT it has the correct form to pass unscathed through all the traps set out by low energy phenomenology.

In the physics applications which follow, we shall make extensive use of two main characteristics of the general framework discussed above. First, since we are dealing with an effective theory (the $N = 1$ SUGAR action is non-renormalizable), the superpotential g is not anymore necessarily constrained by renormalizability to be at most cubic, but it may contain any higher powers, suitably scaled, by inverse powers of M_{Pl} , the natural cut-off of the theory¹⁹⁾. Secondly, because of the non-renormalization theorems⁴⁾ of SUSY (SET IT AND FORGET IT principle), we may set, as we wish, certain parameters equal to zero, even if no symmetry implies that - a very different situation from ordinary gauge theories. Here, no apologies are needed. As explained in detail before, most of the physics is contained in the "observable" sector superpotential $f(\chi_i)$ [see (20)]. Here we shall assume that, in one way or another, the "hidden" sector has played its role, as discussed previously, and we shall concentrate on the form of $f(\chi_i)$. We follow the natural (cosmic) evolution of things starting at energies below M_{Pl} and "coming down" to M_W . So we distinguish physics around the GUT scale (M_X) and physics around the electroweak (E-W) scale (M_W).

All physics from M_{Pl} down to (and including) low energies should emerge from such a programme. We will show next that this is indeed possible.

3. PHYSICS WITH SIMPLE ($N = 1$) SUPERGRAVITY

A. Physics Around the GUT Scale (M_X)

The superhigh energy regime ($\sim 10^{16}$ GeV) is the theorists' paradise. There is a lot of freedom in building models, even though the constraints both from particle physics and cosmology become tighter and tighter. For definiteness, simplicity, and out of habit, we shall take as our prototype GUT an $SU(5)$ type model¹⁾. All GUT physics information will be contained in f_{GUT} , the GUT part of the "observable sector" superpotential. There is no consensus about the definite form of this superpotential, but it should unavoidably contain a piece (f_I) that breaks $SU(5)$ down to $SU(3) \times SU(2) \times U(1)$

and if possible, a piece (f_{II}), providing some explanation about the tree-level gauge hierarchy problem, so we write:

$$f_{GUT} = f_I + f_{II} \quad (29)$$

For example, we may take²⁸⁾

$$f_I = \frac{a_1}{M} X^4 + \frac{a_2}{M^2} X^2 T_\tau(\Sigma^3) \quad (30)$$

and²⁹⁾

$$f_{II} = \bar{\theta} H \left(\lambda_1 \frac{\Sigma^2}{M} + \lambda_2 \frac{\Sigma^3}{M^2} + \dots \right) + \bar{H} \theta \left(\lambda'_1 \frac{\Sigma^2}{M} + \lambda'_2 \frac{\Sigma^3}{M^2} + \dots \right) + M_\theta \bar{\theta} \theta \quad (31)$$

where $X = \underline{1}$, $\Sigma = \underline{24}$, $(\bar{\theta}) = (\underline{\bar{5}})$, $(\bar{H}) = (\underline{\bar{5}})$ are chiral superfields of $SU(5)$. The Higgs fields H and $\bar{H}$ couple to quark and lepton fields in the usual way. All components of θ and $\bar{\theta}$ have a bare mass M_θ (which is taken to be of order M_X or larger), and so remain heavy after $SU(5)$ breaks to $SU(3) \times SU(2) \times U(1)$. After minimizing the potential, obtained by plunging into (22) the sum of f_I and f_{II} as given by (30) and (31), we get zero vev's for H and $\bar{\theta}$ but non-zero ones for

$$\begin{aligned} \langle X \rangle &= \left(\frac{m_{3/2}}{M} \right)^{3/8} M \\ \langle \Sigma \rangle &= \left(\frac{m_{3/2}}{M} \right)^{1/4} M \end{aligned} \quad (32)$$

Furthermore, we find²⁸⁾ that the $SU(3) \times SU(2) \times U(1)$ symmetric minimum is the lowest one for all values of a_1 and a_2 , with a value²⁸⁾

$$V_{eff} \simeq - \left(\frac{m_{3/2}}{M} \right)^{5/2} M^4 \quad (33)$$

What do these results mean? First, since the vev of Σ sets the scale of $SU(5)$ breaking, we find that the GUT scale M_X satisfies²⁸⁾

$$M_X^4 \simeq O(m_{3/2} M^3) \quad (34)$$

which is a highly successful relation. Using as an input the non-hierarchical and easy to explain ratio $(M_X/M) \sim 10^{-2}-10^{-4}$, we obtain that $m_{3/2} \sim O(100 \text{ GeV})$! More generally, relations of the form $M_X^{2p-2} \simeq O(m_{3/2} M^{2p-3})$ with $p \geq 3$, are also possible²⁸⁾ by suitably modifying the exponents in (30). The supergravity hierarchy problem has been solved in a rather simple way.

Secondly, the $SU(3) \times SU(2) \times U(1)$ symmetric minimum is lower in energy density than the $SU(5)$ symmetric minimum $X = \Sigma = 0$ by an amount $(m_{3/2}/M)^{5/2} M^4$. Thirdly, the barrier between these two minima is never larger than $(m_{3/2}/M)^{5/2} M^4$, the same as the splitting between the states. Why this is so can be seen by noting that if we replace X by its vev (32) in (30), the effective renormalizable self-coupling of Σ is $10^{-12} \text{tr}(\Sigma^3)$. Thus we have generated²⁸⁾ a small renormalizable coupling for Σ from our starting point of only non-renormalizable interactions among X and Σ . This small coupling suppresses the barrier between the $SU(5)$ and the $SU(3) \times SU(2) \times U(1)$ phases. The consequences of this suppression for supercosmology³⁰⁾ are difficult to over-estimate. Simply, it now makes possible the transition from the $SU(5)$ to the $SU(3) \times SU(2) \times U(1)$ phase at temperatures $T \sim 10^{10} \text{ GeV}$, which was previously blocked, since the barrier between the two phases was of the order of $(M_X)^4$. Incidentally, in this picture, the number density of GUT monopoles is naturally suppressed³⁰⁾ below its present experimental upper bound.

It should be clear that the basic result - small renormalizable couplings arising from non-renormalizable ones suppressed only by inverse powers of M - is quite general and does not depend on the detailed form of the superpotential (f_I) ²⁸⁾. The main characteristics of these types of models^{28),31)} are that they provide relations of the type (34); they make possible "delayed" $SU(5)$ to $SU(3) \times SU(2) \times U(1)$ phase transitions at $T \sim 10^{10} \text{ GeV}$, and they contain^{28),31)} more "light" particles than the ones in the minimal SUSY $SU(3) \times SU(2) \times U(1)$ model. This last fact may sound dangerous when calculating M_X , $\sin^2 \theta_{E-W}$ and m_b/m_τ , since in general an arbitrary increase of "light" stuff gives an out-of-hand increase^{32),33)}, and thus experimentally unacceptable values for, the above-mentioned quantities. A more careful analysis³⁴⁾ of these cosmological acceptable models (CAM)³⁴⁾ shows that they make predictions as successful (for $\sin^2 \theta_{E-W}$, m_b/m_τ , ...) as at least the ones^{32),33)} of the phenomenologically acceptable minimal type models (MIMO). For a detailed, thorough phenomenological analysis of CAMs, see Ref. 34).

Next, we discuss²⁹⁾ physics related to f_{II} as given in (31). The vev of Σ does not only break $SU(5)$ to $SU(3) \times SU(2) \times U(1)$ but also provides a mass term which mixes the colour triplets in H and $\bar{H}$ with those in θ and $\bar{\theta}$. However, there is no weak doublet in the 50, and so the weak doublets in H and $\bar{H}$ remain massless. The colour triplets will have a mass matrix²⁹⁾

$$\begin{pmatrix} 0 & \sim \frac{M_X^2}{M} \\ \sim \frac{M_X^2}{M} & M_\theta \end{pmatrix} \quad (35)$$

where M_θ should be of order M_X or larger ($\leq M$), to avoid having particles from θ and $\bar{\theta}$ influencing the renormalization group equations at scales below M_X (or even M). The eigenvalues of this mass matrix are $O(M_\theta)$ and $O(M_X^4/M_\theta M^2)$; this latter eigenvalue is about 10^{10} GeV for $M_X \sim O(10^{16}$ GeV) and $M_\theta \sim O(M)$. In this case, the Higgs colour triplet can be used to generate^{28),30)} the baryon number of the Universe after the $SU(5)$ to $SU(3) \times SU(2) \times U(1)$ transition which, as discussed earlier, occurs at temperatures $T \sim 10^{10}$ GeV in CAMs³⁴⁾. It is remarkable that $O(10^{10}$ GeV) is the lower bound³⁵⁾ allowed for colour triplet Higgs masses from present limits on proton decay ($\tau_p > 10^{31}$ years). If indeed there are 10^{10} GeV Higgs triplets, then protons should decay predominantly³⁶⁾ to $\bar{\nu}_\mu K, \mu K$ with a lifetime $\sim O(10^{31})$ years).

The role of supergravity in this natural explanation²⁹⁾ of the Higgs triplet-doublet mass splitting ($\equiv$ tree-level gauge hierarchy problem) is fundamental, in several aspects. The same kind of explanation had been suggested before³⁷⁾ in the framework of renormalizable global SUSY GUTs, where Σ^2 in (31) was replaced by a 75 of $SU(5)$ and higher than two powers of Σ were absent³⁷⁾. Unfortunately, the use of 75 drastically conflicts with cosmological scenarios²⁸⁾⁻³¹⁾ based on SUSY GUTs. The barrier between the $SU(5)$ and $SU(3) \times SU(2) \times U(1)$ phases is impossible to overcome unless most of the 75 is very light ($\sim M_W$). But then, all hell breaks loose. A light 75 makes the gauge coupling in the $SU(5)$ phase decrease at lower energies so there is no phase transition at all^{28),30)}. Furthermore, the presence of these new light particles in the $SU(3) \times SU(2) \times U(1)$ phase changes the renormalization group equations, and prevents perturbative unification. On the contrary, in SUGRA theories, since we may use non-renormalizable terms, we may replace the fundamental 75 by an "effective" 75 contained in Σ^2 . Unlike a light 75, a light 24 neither makes the $SU(5)$ gauge coupling decrease at energies below M_X , nor upsets perturbative unification²⁹⁾. The previously mentioned cosmological scenarios²⁸⁾⁻³¹⁾ can proceed without modification. In addition, SUSY non-renormalization theorems⁴⁾ ensure the stability of the triplet-doublet splitting

to all orders in perturbation theory. Since the only modifications of the theory are at the GUT scale M_X , it seems that we have got²⁹⁾ a harmless and elegant solution of the tree-level, and for that matter, to all orders in perturbation theory, gauge hierarchy problem.

SUGAR models give good physics at the GUT scale - unique, cosmologically acceptable breaking of $SU(5)$ to $SU(3) \times SU(2) \times U(1)$, with an explanation of the smallness of the gravitino mass²⁸⁾ [(34)-like relations], and a natural explanation²⁹⁾ of the Higgs triplet-doublet splitting, cosmologically fitted and general enough. We believe that even if the very specific form of f_I in (30) may change, then f_{II} as given by (31) (or its obvious generalization to other GUT models) will be always a useful part of the f_{GUT} .

After finding plausible explanations for the SUGAR hierarchy problem [gravitino mass $\sim O(100 \text{ GeV})$], the tree-level and higher orders gauge hierarchy problem (triplet-doublet Higgs splitting), it is time to explain the last gauge hierarchy problem (1), i.e., why does $M_W/M_X \leq O(10^{-13})$? This problem brings us naturally to our next subject.

B. Physics Around the Electroweak (E-W) Scale (M_W)

Although there is no consensus on the best way to incorporate grand unification in SUGAR models, a unique minimal low energy model has recently emerged³⁸⁾⁻⁴⁰⁾. In this model, the physics of the TeV scale is described by an effective $SU(3) \times SU(2) \times U(1)$ gauge theory, in which the breaking of weak interaction gauge symmetry is induced by renormalization group scaling of the Higgs (mass)² operators¹⁹⁾. Much of the attractiveness of this model stems from the fact that no gauge symmetries or fields beyond those required in any low energy SUSY theory are included. Sometimes, it may happen, as is the case of Cosmological Acceptable Models³⁴⁾ (CAMs), that there are GUT relics which are light ($\sim M_W$), but they do not seem to play any fundamental role at low energies, so we may neglect them in our present discussion. Furthermore, adding random chiral superfields to the low energy theory may be problematic. For example, the presence of a gauge singlet superfield coupled to the Higgs doublets and added to trigger $SU(2) \times U(1)$ breaking^{41),24),27)}, usually (but not always⁴²⁾) destroys⁴³⁾ any hope of understanding the gauge hierarchy problem; the reason being⁴³⁾ that in a GUT theory, the gauge singlet does not only couple to the Higgs doublets but also to their associate, superheavy colour triplets. Then we have to try hard⁴²⁾ to avoid 10^{10} GeV Higgs doublet masses, generated by⁴³⁾ one-loop effects involving colour triplets. Something smells fishy.

We focus then on the standard low energy $SU(3) \times SU(2) \times U(1)$ gauge group, containing three generations of quarks and leptons, along with two Higgs doublets, as chiral superfields. The low energy effective superpotential (f_{LES}) of the model consists only of the

usual Yukawa couplings of quark and lepton superfields to the Higgs superfields, along, in general, with a mass term coupling the two Higgs doublets, H_1 and H_2 . Explicitly, in a standard notation:

$$f_{LES} = h_{ij} U_i^c Q_j H_2 + \tilde{h}_{ij} D_i^c Q_j H_1 \quad (36)$$

$$+ f_{ij} L_i E_j^c H_1 + m_4 H_1 H_2$$

where a summation over generation indices (i,j) is understood and $Q(U^c)$ denote generically quark doublets (charge $-2/3$ antiquark singlet) superfields, while $L(E^c)$ refers to lepton doublets (charge -1 antilepton singlet) superfields. With the exceptions of the top quark Yukawa coupling and the mass parameter m_4 , which in principle may be of order $O(m_{3/2})$, all other parameters appearing in (36) contribute to the masses of the observed quarks and leptons and are known to be small. Neglecting these small couplings, the effective Low Energy Potential (V_{LEP}) can be written as [see (22), (23) and (27)]:

$$V_{LEP} = \sum_{i=1}^3 \left\{ m_{L_i}^2 |L_i|^2 + m_{E_i^c}^2 |E_i^c|^2 + m_{Q_i}^2 |Q_i|^2 + \right.$$

$$\left. + m_{U_i^c}^2 |U_i^c|^2 + m_{D_i^c}^2 |D_i^c|^2 \right\} \quad (37)$$

$$+ m_1^2 |H_1|^2 + m_2^2 |H_2|^2 + A h_t m_{3/2} (U_3^c Q_3 H_2 + h.c.)$$

$$+ B m_{3/2} m_4 (H_1 H_2 + h.c.) + h_t m_4 (H_1^+ Q_3 U_3^c + h.c.)$$

$$+ h_t^2 \left(|Q_3|^2 |U_3^c|^2 + |Q_3|^2 |H_2|^2 + |U_3^c|^2 |H_2|^2 \right)$$

$$+ \text{"D-terms"}$$

The effective parameters appearing in (37) take, at large scales ($\sim M_X$), the values

$$m_1^2(M_X) = m_2^2(M_X) = m_{3/2}^2 + m_4^2(M_X) \quad (38)$$

$$m_{Q_i}^2(M_X) = m_{U_i}^2(M_X) = m_{D_i}^2(M_X) = m_{L_i}^2(M_X) = m_{E_i}^2(M_X) = m_{3/2}^2$$

$$A(M_X) = A$$

$$B(M_X) = A - 1 \quad (i = 1, 2, 3)$$

as dictated by (27). It should be stressed once more that the boundary conditions (38) are exact, if we only neglect corrections at the Planck scale, ignore the scaling of parameters from M to M_X , and pay no attention to corrections²⁶⁾ at the GUT scale. All these effects are expected to be small and it is assumed that they do not seriously disturb (38) and the picture hereafter.

It is apparent from (37) that SUGAR models can easily succeed in giving weak interaction scale masses ($m_{3/2} \sim M_W$) to squarks, sleptons and gauginos [see (27)]. Alas, SUGAR models also give large positive (mass)² to the Higgs doublets, thus making the breaking of $SU(2) \times U(1)$ difficult. One way to overcome this difficulty is the introduction^{41), 24), 27)} of a gauge singlet coupled to H_1 and H_2 , but, as mentioned above, with disastrous effects⁴³⁾ for the gauge hierarchy. A particularly simple solution to the $SU(2) \times U(1)$ breaking relies upon the fact that the boundary conditions (38) need be satisfied only at M_X (or M), and that large renormalization group scaling effects can produce a negative value for m_H^2 at low energies¹⁹⁾. The full set of renormalization group equations for the parameters in V_{LEP} (37) has been written elsewhere⁴⁴⁾. Here we concentrate on the most interesting equation, the one for the mass-squared of the Higgs (m_2^2), which gives mass to the top quark:

$$\mu \frac{\partial}{\partial \mu} \begin{bmatrix} m_2^2 \\ m_{U_3}^2 \\ m_{Q_3}^2 \end{bmatrix} = \frac{g_t^2}{8\pi^2} \begin{bmatrix} 3 & 3 & 3 \\ 2 & 2 & 2 \\ 1 & 1 & 1 \end{bmatrix} \begin{bmatrix} m_2^2 \\ m_{U_3}^2 \\ m_{Q_3}^2 \end{bmatrix} + \frac{|A|^2 g_t^2 m_{3/2}^2}{8\pi^2} \begin{bmatrix} 3 \\ 2 \\ 1 \end{bmatrix} - \frac{g_t^2 m_4^2}{8\pi^2} \begin{bmatrix} 0 \\ 2 \\ 1 \end{bmatrix} - \frac{8 g_3^2 M_3^2}{3\pi} \begin{bmatrix} 0 \\ 1 \\ 1 \end{bmatrix} \quad (39)$$

where we have neglected gauge couplings other than the "coloured" one, $\alpha_3 (\equiv g_3^2/4\pi)$, M_3 is the gluino mass [see (26) and (27)], and Yukawa couplings other than h_t , for the top, have been dropped. The physics content of (39) is apparent. Since μ is decreasing (we come from high energies down to low energies), the sign of the first two terms in (39) is such as to make all $m_{2,U,Q}^2$ smaller at low energy, with the decrease of m_2^2 becoming more pronounced because of the 3:2:1 weighting. On the other hand, the sign of the last two terms in (39) is such as to make m_U^2 and m_Q^2 (the squark masses) larger at low energy, but have no direct effect on m_2^2 [notice the "zeros" in the corresponding matrices in (39)]. Indirectly though, the net effect on m_2^2 of the last two terms in (39) is to enhance further its decrease at low energies, by increasing m_U^2 and m_Q^2 , which then drive down m_2^2 via the first two terms of (39). This is exactly what we are after! We want large ($\sim M_W^2$) and positive squarks and slepton (masses)², BUT negative Higgs (mass)² to trigger $SU(2) \times U(1)$ breaking. The ways of obtaining negative Higgs (mass)² now become clear [see (39)]. We have to use either a large top Yukawa coupling (h_t), or large A , or large m_4 , or a fourth generation to provide large Yukawa couplings, or some suitable, physically plausible combination of the above possibilities. There are pros and cons for every one of the above situations. In the case of large h_t , a lower bound on the mass of the top quark is set^{19),38)-40)}:

$$m_t > 0 \text{ (60 GeV)} \quad (40)$$

which some people may find uncomfortable. We may avoid a large h_t by moving it into the large $A(>3)$ regime^{38),45)}. The price, though, is high^{38),45)}. The phenomenologically acceptable vacuum becomes unstable against tunnelling into a vacuum in which all gauge symmetries, including colour and electromagnetism, are broken. We must^{38),45)} then arrange things in such a way that the lifetime for this vacuum decay process is greater than the age of the Universe. Some people, not without reason, may find this possibility dreadful. We may avoid large h_t and/or large A by using³⁴⁾ non-vanishing $m_4 (\sim m_{3/2})$ where a rather satisfactory picture then emerges³⁴⁾. Some people may object here to the basic assumption of large $m_4 (\sim m_{3/2})$, since in the case of natural triplet-doublet Higgs splitting-type models²⁹⁾ [see (31)], m_4 has a tendency to be small, if not zero, even though other sources of m_4 may be available.

Finally, we come to the possibility of a fourth generation which, suitably weighed, may help us to avoid large h_t , A , or m_4 . The problem here is that low energy phenomenology (evolution of coupling constants, $m_b/m_\tau, \dots$)³²⁾⁻³⁴⁾ as well as firm cosmological results like nucleosynthesis (especially ^4He abundance)⁴⁶⁾, may suffer almost unacceptable modifications. Furthermore, one has to watch out for the mass of the fourth generation charged lepton, since it

is going to behave like m_2^2 in (39), and thus m_{L_4, E_4}^2 may easily go negative, breaking electromagnetic gauge invariance.

Whatever mechanism (if any) turns out to be correct, it is rather remarkable that in SUGAR-type models, there is a simple explanation of the breaking of $SU(2) \times U(1)$ and of the non-breaking of $SU(3) \times U(1)_{E-M}$. Furthermore, for the first time, we have a simple explanation of why $M_W \lll M_X$ (or M), i.e., a simple solution of the cumbersome gauge hierarchy problem. Starting with a positive Higgs (mass)², of order $m_{3/2}^2$ at M_X , and noticing that [see (39)] the evolution with μ^2 of the Higgs (mass)² is very slow (logarithmic), it is not surprising that we have to come down a long way in the energy scale, before the Higgs (mass)² turns negative and is thus able to trigger $SU(2) \times U(1)$ breaking. For example, in a class of models³⁸⁾ characterized by "small" gravitino masses ($\ll O(M_W)$) and by a Coleman-Weinberg-type⁴⁷⁾ radiative $SU(2) \times U(1)$ breaking, occurring naturally, we get⁴⁸⁾, by dimensional transmutation

$$M_W \simeq \Lambda_{QCD} \cdot \frac{g_2(M_W)}{24} \exp\left(\frac{2\eta}{\alpha_G} F\left(\frac{h_t(M_X)}{\alpha_G}; A\right)\right) \quad (41)$$

where g_2 denotes the $SU(2)$ gauge coupling constant, α_G is the GUT fine structure constant and F is a rather involved function of its indicated variables. Using standard values for $g_2(M_W)$ (~ 0.67), α_G ($\sim 1/25$) and reasonable values for $h_t(M_X)$ ($0.2-0.3$), $A(M_X)$ ($2-3$), (41) gives⁴⁸⁾

$$M_W \simeq (300-600) \Lambda_{QCD} \quad (42)$$

a rather remarkable equation from many points of view. It does not only give $M_W \lll M_X$ as it was required by the gauge hierarchy, but it also provides a new and successful relation between the scale of electroweak unification M_W (~ 80 GeV) and the fundamental scale of strong interactions, Λ_{QCD} ($\sim 0.15-0.3$ GeV) in terms of dimensionless parameters. Clearly, this occurs because the rapid final stages of evolution of m_2^2 are driven by the increases in t -quark Yukawa coupling, and more importantly, in the squark masses which occur when $g_3^2/4\pi$ becomes large. It is only in the SUSY Coleman-Weinberg scenario^{38),34)} that the weak interaction scale is related to that of strong interactions. This contrasts with what usually happens in weak gauge symmetry breaking in SUGAR models³⁹⁾⁻⁴⁰⁾ where M_W

is connected to $m_{3/2}$, but it is not directly related to the strong interaction scale.

Another very amazing fact is that the values of the parameters of the low energy world seem to co-operate with us. Since quarks are feeling strong interactions, (39) tells us that quarks may enjoy large masses (Yukawa couplings) without making squark (masses)² negative, because of the last term ($\sim \alpha_3$), which easily balances off large Yukawa couplings, without any sweat. On the other hand, since leptons are not feeling strong interactions, the balance-off between the weak gauge couplings and large Yukawa coupling becomes extremely delicate and could be problematic. How nice that for all three generations, leptons and down quarks weigh less than 5 GeV and especially for the third generation that the top quark (t) is heavier than the bottom quark (b). An inverse situation would be disastrous, because in any reasonable GUT, a very heavy b quark would mean a very heavy τ lepton, thus making electromagnetic gauge invariance tremble in such SUGAR-type schemes. I will not go any further into the esoterics of this type of $SU(2) \times U(1)$ breaking models, since a rather thorough and detailed exposé of these types of theories and of their phenomenological consequences is now available³⁴⁾. It should be stressed that things are now very constrained, as we see from the Table, taken from Ref. 34), where the whole low energy spectrum is worked out, in terms of very few parameters, $m_{3/2}$, $A(M_X)$, $\xi \equiv M_0/m_{3/2}$ [see (26)] and $m_t(M_X)$. Eventually, with more theoretical insight, we hope to determine even these very few parameters, thus predicting uniquely the low energy spectrum. For example, we have already discussed ways of determining $m_{3/2}$ [see (34)], while some people may favour $A(M_X) = 3$ as a natural solution⁴⁹⁾ to the absence of the cosmological constant problem, etc. Among other interesting things contained in the Table, the existence of a very light ($\sim [3-6]$ GeV) neutral Higgs, with the usual Yukawa couplings to matter, should not escape our attention. Since such a particle is a common feature of a large class of models^{38),34)}, a search in the $Y \rightarrow H^0 + \gamma$ channel, which is expected to be a few per cent of the $Y \rightarrow \mu^+\mu^-$ decay, may turn out to be very fruitful.

There are other, very interesting features concerning low energy phenomenology stemming out from the general form of V_{LEP} [(37) and (38)] in SUGAR models. Very tight constraints coming from natural suppression of flavour changing neutral currents⁵⁰⁾ (FCNC), absence of large corrections to $(g-2)^{51)}$ and ρ ^{39),52)} ($\equiv (M_W/M_Z \cos\theta)^2$) as well as to θ_{QCD} ⁵⁾, which have been the nemesis^{50)-52),5)} of SUSY models with arbitrary and explicit soft SUSY breaking, are satisfied in SUGAR models. The highly-constrained set of soft SUSY operators (27) in SUGAR models fits the bill^{19),39)}. Concerning FCNC, (38) guarantees the super-GIM mechanism, since the mass matrices for the quarks and leptons are diagonalized by the same transformation that renders the mass matrices for their scalar partners and gluino couplings generation diagonal. Despite the fact that this property does

TABLE
PARTICLE SPECTRUM ³⁴⁾

		CAM	MIM	CAM	MIM	CAM	MIM
	$A(M_X)$	3	3	2.8	2.8	2.0	1.6
	$m_{\frac{1}{2}}$	15	15	15	15	15	15
	ξ	2.8	2.2	3.2	3.1	3.5	1.9
	$m_\nu(M_X)$	15	17	16	18	11	7
	top	25	25	35	35	50	50
All families	(sleptons) _L	29	27	32	36	35	25
	(sleptons) _R	21	20	22	23	23	19
1st and 2nd families	(squark) _L	58	77	66	108	72	67
	(squark) _R }	54	74	61	104	67	65
		54	74	60	103	66	65
3rd family	(sbottom) _L	58	76	64	106	68	66
	(sbottom) _R	54	74	60	104	66	65
	(stop) _L	81	96	95	132	112	106
	(stop) _R	26	54	23	78	21	37
	Charged Higgses	96	93	95	94	88	83
	Neutral Higgses }	106	104	105	105	100	95
		3	3	4.4	5.3	6	5
	"Axion"	51	46	49	48	35	19
	Gluinos	42	84	47	118	52	72
	Photino	11	4.6	9.4	7.3	4	3
	HW, WH-inos }	89	87	87	94	84	90
		78	82	79	79	82	75
	HZ, HZ-inos }	99	108	99	116	101	106
		92	85	91	80	88	83
	Axino	26	23	24	24	17	9

Physical mass spectrum of the cosmologically acceptable model (CAM) and the minimal model (MIM) corresponding to the same gravitino mass $m_{\frac{3}{2}} \approx 15$ GeV for top quark masses equal to 25 GeV, 35 GeV and 50 GeV respectively. ξ denotes the ratio of the gaugino to the gravitino mass at M_X . All masses are in GeV units. The light neutral Higgs gets its mass via radiative corrections.

not survive, in general, after renormalization, it has been shown^{53),54)} that these effects are controllable. Furthermore, the Buras stringent upper bound⁵⁵⁾ on the top quark mass [$< 0(40 \text{ GeV})$], coming from kaon phenomenology ($K_L - K_S$ and $K_L \rightarrow \mu^+ \mu^-$ systems), is avoided⁵³⁾ in SUGAR models. There are a lot of cancellations between ordinary and SUSY contributions in K processes⁵⁶⁾, such that the top quark mass may be stretched up to 100 GeV without problem⁵³⁾. That sounds very satisfactory, especially for SUGAR models³⁸⁾⁻⁴⁰⁾ that do need a large top quark mass for $SU(2) \times U(1)$ breaking. It looks like a self-service situation. Similar comments apply in the case of $(g-2)_\mu$ or ρ , where it has been shown that SUGAR model contributions are acceptable^{34),39),52)}. Typical values for SUGAR contributions are³⁴⁾ $|\Delta(g-2)_\mu| \lesssim (3 \cdot 10^{-9})$ and^{39),52)} $\Delta\rho \leq 0.01$, which compare favourably with the present experimental upper bounds of $(4 \cdot 10^{-8})$ and (0.03) respectively, but are large enough to be interesting. Better experimental bounds, especially on $\Delta\rho$, could be revealing.

Concerning θ_{QCD} , it has been shown^{19),57)} that in SUGAR-type models, we not only understand the non-renormalization⁵⁾ of θ , but we also understand^{19),57)} why θ is zero or small ($< 10^{-9}$) to start with. This fact is related with our freedom, discussed before, to use non-minimal ($\neq \delta_{\alpha\beta}$) $f_{\alpha\beta}$ in (15). It has already been observed¹⁸⁾ that the gauge kinetic term in (15), as well as giving rise to the canonical

$$- \frac{1}{4} F_{\mu\nu}^\alpha F_{\mu\nu}^\beta (\text{Re } f_{\alpha\beta}) \quad (43)$$

could also yield the CP-violating θ vacuum term

$$\epsilon^{\mu\nu\rho\sigma} F_{\mu\nu}^\alpha F_{\rho\sigma}^\beta (\text{Im } f_{\alpha\beta}) \quad (44)$$

We know that θ_{QCD} is $< 0(10^{-9})$ experimentally, and that θ is not renormalized in a supersymmetric theory⁵⁾. It is finitely renormalized when supersymmetry is broken, but⁵⁾ this is plausibly only by an amount $\delta\theta = 0(10^{-16})$ in the popular Kobayashi-Maskawa model⁵⁸⁾. Thus we see that θ should be less than $0(10^{-9})$ in a supersymmetric GUT and may be very small. An attractive hypothesis is that $f_{\alpha\beta}$ is a function with only real coefficients as found in extended supergravities⁵⁹⁾. In this case, $\text{Im } f_{\alpha\beta} = 0$ when $\langle 0 | \phi | 0 \rangle = 0$, and the theory is CP-invariant in the gauge sector. If some of the ϕ then acquire complex vacuum expectation values, they will induce a non-zero value of $\text{Im } f_{\alpha\beta}$ and hence violate CP spontaneously in the gauge sector, which is a new twist on an old proposal⁶⁰⁾. If the moduli of

some of these complex $\langle 0|\phi|0\rangle$ were $O(M_{Pl})$, then the effective θ parameter would be $O(1)$ which is phenomenologically unacceptable. However, it is easy to imagine scenarios where θ is much smaller. For example, if $f_{\alpha\beta} = \delta_{\alpha\beta} + O(\phi^2/m_{Pl}^2)$ and the culprit $\langle 0|\phi|0\rangle = O(m_X)$, then

$$\theta = O\left(\left(\frac{m_X}{M_{Pl}}\right)^2 \times \text{some small (?) angle}\right) \leq O(10^{-7}) \quad (45)$$

and the phenomenological constraint on θ_{QCD} could easily be respected. If the only complex $\langle 0|\phi|0\rangle$ were $O(m_W)$, or if all the $\langle 0|\phi|0\rangle$ were real as in all supersymmetric GUTs proposed to date, then the bare $\theta = 0$. Hence supergravity offers the other half of an answer to the θ vacuum problem. It should be stressed that low energy supergravity models have new sources^{39),57),61),62)} of CP violation beyond the standard model. Unfortunately, they do not shed more light on the smallness of the observed CP violation in the K system⁶¹⁾. On the other hand, potential large contributions^{61),57)} to ϵ' , θ_{QCD} and to the Dipole Electric Moment Of the Neutron (DEMON) put, in general, severe constraints^{57),61)} on these new possible CP-violating phases. Actually, it seems almost unavoidable^{5),39),57),62)} that in SUGAR models the DEMON should be near, but not above, the present experimental upper bound⁶³⁾ of $6 \cdot 10^{-25}$ e.cm, a rather drastic and experimentally testable prediction, in sharp contrast with the standard model prediction $O(10^{-30}$ e.cm). We may know soon.

Turning now to baryon decay, an interaction of the form¹⁹⁾

$$f \ni \frac{\lambda}{M} \bar{F} T T T \quad (46)$$

where $\bar{F}$ is a $\bar{5}$ of matter (quark + lepton) chiral superfields in SU(5), T is a $\bar{10}$ of matter superfields and λ is some generic Yukawa coupling, could replace the Higgs exchange in the Weinberg-Sakai-Yanagida⁶⁴⁾ loop diagram for baryon decay. The magnitude of the diagram with (46) relative to the conventional Higgs diagram is:

$$\left(\frac{\lambda}{M}\right) / \left(\frac{\lambda^2}{M_{H_3}}\right) \approx \left(\frac{M_{H_3}}{\lambda M}\right) \quad (47)$$

The ratio (47) could easily be >1 , making a non-renormalizable superpotential interaction the dominant contribution to proton decay. A careful analysis of SUGAR-induced baryon decay shows⁶⁵⁾,

surprisingly enough, that the expected hierarchy of decay modes is similar to that³²⁾ coming from conventional minimal SUSY GUTs. One might have wrongly expected that no hard and fast predictions could be made about gravitationally-induced baryon decay modes. Anyway, this mechanism could give observable baryon decay even if the GUT mass $M_X \approx M$.

Incidentally, similar terms like (46) have been considered⁶⁶⁾ in efforts to explain the "lightness" of the first two generations of quarks and leptons. One replaces⁶⁶⁾ direct Yukawa couplings for the first two generations with (very schematically):

$$f \ni \frac{\tilde{\lambda}_2}{M} \bar{H} \Sigma T_2 \bar{F}_2 + \frac{\tilde{\lambda}'_2}{M} H T_2 T_2 \Sigma$$

$$+ \frac{\tilde{\lambda}_1}{M^2} \bar{H} \Sigma^2 T_1 \bar{F}_1 + \frac{\tilde{\lambda}'_1}{M^2} H T_1 T_1 \Sigma^2 + \dots \quad (48)$$

which not only repairs^{19),66)} wrong relations like $m_d(M_X) \approx m_e(M_X)$, very difficult to correct⁶⁷⁾ in conventional SUSY GUTs, but also provides reasonable masses for the first two generations. Indeed, it follows from (48) that the second generation is getting masses $(M_X/M)M_W \sim (0.1-1 \text{ GeV})$, while the first generation masses are $(M_X/M)^2 M_W \sim (1-10 \text{ MeV})$, exactly what was ordered⁶⁶⁾. It is amazing that in SUGAR models, by increasing $M_X (\sim 10^{16} \text{ GeV})$, relative to its ordinary GUT value (10^{14} GeV), and by decreasing M_{Pl} , what is relevant is the superPlanck scale $M (\sim 10^{18} \text{ GeV})$, the highly-desired ratio $(M_X/M) \sim 10^{-2}$, appears naturally. It seems now, for the first time, that gravitational interactions may be responsible for the masses of at least the first two generations. Once more, non-renormalizable interactions contained in SUGAR models provide a simple solution⁶⁶⁾ to another hierarchy problem, the fermion mass hierarchy problem.

4. FINAL REMARKS

We have shown that gravitational effects, as contained in SUGAR theories, cannot be neglected anymore in the regime of particle physics. On the contrary, it may be that supergravitational effects are really responsible: for the SU(5) breaking at M_X with an automatic triplet-doublet Higgs splitting, for the SU(2) $\times$ U(1) breaking [and SU(3) $\times$ U(1)_{E-M} non-breaking] at M_W , naturally exquisitely smaller than M_X , for the "constrained" soft SUSY breaking at $m_{3/2}$, hierarchically smaller than M in a natural way, for definite, at present experimentally acceptable departures from the

"standard" low energy phenomenology [like the DEMON, or $\Delta\rho$, with values below, but not far from, their present experimental upper bounds, or the existence of very light ($< O(10 \text{ GeV})$) neutral Higgs bosons], as well as a rather well-defined low-energy SUSY spectrum, for observable baryon decay even if $M_X \approx M$, and for the light fermion masses of the first two generations. Furthermore, supergravity theories may provide, for the first time, a problem-free cosmological scenario, from primordial inflation⁶⁸⁾ through GUT phase transitions²⁸⁾⁻³¹⁾ to baryon and nucleosynthesis, ostracizing troublesome particles such as GUT monopoles^{28),30)}, gravitinos⁶⁹⁾, Polonyi fields⁷⁰⁾ or other SUSY relics⁷¹⁾.

Putting the whole thing together, it becomes apparent that spontaneously broken $N = 1$ local SUSY gauge theories, with their prosperous and appropriate structure, may well serve as an effective theory describing all physics from M_{Pl} down to (and including) low energies, with well-defined and rich experimental consequences. What's next, then? Well, we really have to understand where this highly successful theory comes from. There are reasons to believe²¹⁾ that $N = 8$ extended SUPERGRAVITY¹³⁾ may provide the fundamental theory. But this next move asks for a deep understanding of physics at Planck energies, which is as exciting as it is difficult, taking into account that even QUANTUM MECHANICS may need modification⁷²⁾, if quantum gravitational effects have to be considered seriously.

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