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PROLOGUE

This talk is divided into three parts. In part A, I discuss possible experimental hints and theoretical prejudices for extension of the standard model, as well as the form of these extensions. Two major frameworks emerge: the (low-energy scale) compositeness framework, discussed in part B and the supersymmetric framework (including grand unification) discussed in Part C.

While the title of my talk has been assigned to me by the organizers of this Conference, the selection of topics and the form of presentation is all mine. Hopefully, the recent blooming experimental activity will soon tell us if my choices have been the right ones.

A. STANDARD MODEL: NOT ENOUGH

1. - GENERALITIES

We believe today that the standard $SU(3)_C \times SU(2)_L \times U(1)$ gauge model[1)-3)] provides a fair description of all presently established Low Energy Phenomenology[4)] (LEP). The identification card for the standard model reads as follows:

1) Framework : relativistic quantum field theory
2) Dimensions: 4 = 3-space + 1-time
3) Basic symmetries:
   a) Poincaré invariance: "gauge group" O(1,3)
   b) Gauge invariance: gauge group: $SU(3)_C \times SU(2)_L \times U(1)$ (1)
4) Building blocks:

Fundamental fields:
- a) quarks-leptons (FERMIONS)
- b) gauge-Higgs fields (BOSONS)

The credentials of the standard model have been established after many years of hard theoretical and experimental work[4)]. There is a set of very crucial tests that any model has to pass in order to be taken seriously. Among others, these tests include:

1) Absence (to the appropriate order) of Flavour Changing Neutral Currents (FCNC) both in the real part[2)] ($\Delta m_{K_L - K_S}$, $K_L \to \mu\mu$, ...) and in the imaginary part[3)] (CP-violation,...)

2) Appropriate dipole moments either magnetic $[(g-2)_{\mu,e}]$ or electric ($d_{n,e}$). The present experimental values are extremely precise

$$(g-2)_\mu^{5)} = \left(2.331.848 \, {}^{+52}_{-40}\right) \cdot 10^{-9} \tag{2}$$

$$d_n^{6)} \lesssim 4 \cdot 10^{-25} \, e \cdot cm \tag{3}$$

3) Validity of the weak isospin $I_W = \frac{1}{2}$ rule[7)], implying $\rho = 1$, where the ρ-parameter is defined by

$$\rho \equiv \left(\frac{M_W}{M_Z \cos\theta_W}\right)^2 \tag{4}$$

in a self-explanatory notation. The present experimental value[8)]

$$\rho = 1.02 \pm 0.06 \tag{5}$$

already leaves no room for large deviations from unity.

Sure enough, the standard model passes these tests with flying colours. I have concentrated on these tests because they are very sensitive to the low-energy structure (particle as well as interaction content), and they put under severe scrutiny any extension of, or alternative to, the standard model. Needless to say, the recent major discoveries[9)] at CERN of the intermediate vector bosons $W^\pm$, Z with characteristics exactly as predicted by the standard model[4)], signal once more that we are on the right track. Then why do we discuss extensions or alternatives to the standard model ? Well, until recently, we had to invoke only theoretical reasons to justify our lust for extension of the standard model, but as we have heard at this conference[10], we may now have some experimental evidence asking for such an extension. That sounds very exciting, so let me first discuss what I think are the possible experimental signals for extending the standard model, and provide later the rather well-known theoretical reasoning of why the standard model cannot be the whole story. Since these new experimental discoveries have been discussed in considerable detail by others at this conference[10] and elsewhere[11], I will be very brief.

2. - EXPERIMENTAL EVIDENCE (?) FOR EXTENSION OF THE STANDARD MODEL

2.1 CERN S$\bar{p}$pS Collider

The CERN collider has not only provided evidence[9] for the standard model, $M_{W,Z}$; $\sin^2\theta_W$; ρ, ..., all measured at the predicted values, but it probably provides evidence for an extension of the standard model, by observing the following types of events:

i) Radiative Z^0 decays[12] of an unexpectedly high rate and of unusual kinematic configuration. Namely, three events (two with e and one with μ) of the type

$$Z \to \ell^+\ell^-\gamma \qquad (\ell \equiv e,\mu) \tag{6}$$

have been found, where the angle between the photon and one of the leptons is rather small or large depending on one's point of view. At first sight, both the rate (20-25%) and the kinematic configuration make conventional QED bremsstrahlung a rather unlikely candidate (?).

ii) "Zen" events[13]: these are events with one jet (or other evidence of activity) on one side of the interaction point, and nothing (missing neutrino or ?) on the other side. UA(1) reported[13),10] five events containing a "monojet" of energy above 40 GeV with no jet having E_T > 10 GeV recoiling in the opposite direction in the azimuthal plane. Also, the same group reported[13),10] two events with a "photon" of transverse momentum above 40 GeV with large missing transverse energy. It is a general feature of the "monojet" events that their jets are quite "small": the charged multiplicity varies between one and three, and the invariant masses of the charged particles with p_T > 0.5 GeV are generally less than about 2 GeV. Furthermore, in addition to monojet events, there seem to be some multijet events with missing p_T. If true, "Zen" events cannot be accommodated in the standard model.

iii) "Wen" events[14]: the UA(2) group has reported[14),10] the discovery of three events of the type e + jet + missing p_T. It is not inconceivable, but by no means mandatory, to "see" these events as W + jet events, where the W + jet invariant mass is around 160 GeV. It is amusing to notice that the UA(2) group have also a 3σ bump in their multijet mass distribution around 150 GeV, but we still have to wait for further evidence for this bump, not observed by the UA(1) group. Once more, these type of events, if true, remain unexplained in the standard model.

iv) The UA(1) collaboration has reported[15] ten dimuon events, of which seven are opposite-sign (four with jet activity and three without) and three are like-sign (one with jet activity and two without). What is most intriguing about the dimuon events, apart from the presence among them of three like-sign events, is the abundance of strange particles that they contain. For example, there is a $\mu^+\mu^+\Lambda^0$ and a $\mu^-\mu^-\bar{\Lambda}^0$ event. Again, the standard model may have a hard time (?) in explaining these events.

I would like to stress here, that beyond the CERN S$\bar{p}$pS magnificent discoveries which may eventually lead to a drastic modification (extension) of the "standard model", there is an accumulation of other experimental facts that drive us to the same conclusion: the standard model is not enough.

2.2 Other possible experimental hints challenging the standard model

i) As we have repeatedly heard in this conference[10),16], it seems by now a well-established fact that the bottom life-time is quite long, the world average being[10]

$$\tau_b = (1.5 \pm 0.4 \pm 0.3)\cdot 10^{-12}\ \mathrm{sec} \tag{7}$$

which combined with the UA(1) possible value of the top quark mass reported here[10]

$$m_t \simeq (40 \pm 10)\ \mathrm{GeV} \tag{8}$$

and the experimental upper bound[4),10]

$$\bar{R} \equiv \frac{(b \to u)}{(b \to c)} \leq 0.04 \tag{9}$$

may mean trouble for the standard CP-violation Kobayashi-Maskawa (KM) six quark model[3]. Too long a b-lifetime means too small KM angles, which in turn means a rather heavy top quark mass for sufficient CP-violation. But if m_t is given by (8), this is not enough. For more details, see later on.

ii) Another potential problem for the standard KM model[3)] may arise from the new value of ε'/ε reported in this conference[17)]

$$\frac{\varepsilon'}{\varepsilon} = \begin{cases} -0.0046 \pm 0.0053 \pm 0.0024 & \text{(CHICAGO-SACLAY)} \\ +0.0045 \pm 0.0080 & \text{(BNL-YALE)} \end{cases} \qquad (10)$$

If eventually the negative sign persists, then some modification of the standard KM model is needed since it predicts (?) a positive sign[18)]. Also, care should be taken for the absolute value of ε'/ε since its lower bound[18)], related to the mass of the top quark, is rather high for a top quark mass given by (8).

iii) The possible evidence for neutrino masses and/or neutrino oscillations, reported once more in this conference[19)], obviously indicates that some extension of the standard model is demanded.

iv) the uninvited, but most welcome guest: the ζ(8.3 GeV) particle. The Crystal Ball group operating at DORIS II announced at this conference[20)] the discovery of a new particle called ζ which is produced in the radiative decay of Υ ($\Upsilon \to \gamma\zeta$), with the following characteristics

$$\begin{aligned} & M_\zeta = (8322 \pm 8 \pm 24)\ \mathrm{MeV} \\ & \Gamma_\zeta < 80\ \mathrm{MeV} \\ & B[\Upsilon(1S) \to \gamma\zeta] \simeq 0.5\% \\ & \frac{B[\Upsilon(2S) \to \gamma\zeta]}{B[\Upsilon(1S) \to \gamma\zeta]} \lessapprox 0.22 \end{aligned} \qquad (11)$$

True enough, $\Upsilon \to \gamma\zeta$ is the "textbook" way to find a "light" Higgs boson. However, both the large branching ratio of Υ(1S) decay to ζ and the unobserved Υ(2S) decay to ζ, as they are given by (11), may speak against a "traditional" one-Higgs (standard model) interpretation of the ζ(8.3) particle. Furthermore, the last limit in (11) provides also severe problems for multi-Higgs models.

A cynic may observe that all experimental evidence against the standard model stands on shaky grounds and may disappear soon! On the other hand, I find it extremely hard to believe that all the S$\bar{p}$pS "exotic" events and all the other experimental facts that I presented will vanish. Even if a few of them remain valid, they are enough to support my statement that we are in for some changes of the standard model, as is also strongly indicated by the theoretical arguments that follow.

3. - THEORETICAL EVIDENCE FOR EXTENSION OF THE STANDARD MODEL

Theorists always felt uneasy with the standard model. For sure, it gives a fair description of what has been observed (at least till last year!), but true enough, it leaves much to be desired. The heart of the matter is the arbitrariness and uncorrelation of far too many parameters. Either on the mass front (fermion, Higgs, gauge boson masses) or on the coupling constant front (gauge, scalar couplings) everything is at random. This is not exactly what one expects from a fundamental theory. It is convenient to look at the problems raised by the three sectors of the theory, namely the fermion, Higgs and gauge sector.

1) Fermion : sector
- i) quarks-leptons: incommunicado → masses; quantum numbers uncorrelated
- ii) superfluous family repetition
- iii) Mass source: arbitrary Yukawa couplings

↓

2) Higgs sector: Arbitrary structure
- i) Yukawas (h_i)
- ii) Self-couplings (λ_i)
- iii) $-\mu^2$: "funny" mass terms

triggers ↓

3) Gauge sector: The most constrained, but SPONTANEOUS BREAKDOWN (SB) arbitrary and problematic as it leads to the

4) Gauge (scale) hierarchy problem:

As is well known, the raison d'être of the Higgs particles is their ability to get vacuum expectation values (v.e.v) different from zero and thus cause spontaneous breakdown of some gauge symmetry. Clearly, the mass of the gauge bosons, m_G, will be proportional to the v.e.v., v: $m_G \sim v$. On the other hand, v will depend on the parameters of the Higgs potential and it is not surprising that v is proportional to the Higgs mass: $v \sim m_H$ and thus finally reaching the conclusion that

$$m_G \simeq \langle H \rangle \equiv v \sim m_H \qquad (12)$$

And here starts the problem. The weak gauge boson ($W^\pm, Z^0$) masses are much lighter than other known or conjectured mass scales in physics, such as the grand unification[21)-25)] (GUT) mass scale M_X (> 10^{15} GeV) or the Planck scale M_P (≃ 10^{19} GeV):

$$\frac{M_W}{M_X} \lesssim O(10^{-13}) \ , \ \frac{M_W}{M_P} = O(10^{-17}) \tag{13}$$

Scalar boson masses are notoriously unstable with a strong tendency to rise to the largest available mass scale. Elementary scalar bosons propagating through space-time foam[26] at the Planck length scale would acquire masses $O(M_P)$, or propagating through the GUT vacuum also pick up masses $O(M_X)$. That is certainly catastrophic, since (12) would imply that $M_W \sim O(M_X$ or $M_P)$ in violent contradiction with (13)! Even if one tries to cancel out these contributions at the tree level, so as to start off with some "light" Higgs fields with masses $O(M_W)$, then radiative corrections will generate mass shifts $O(\alpha^n M_X$ or $\alpha^n M_P)$ still unbearable. Even if one is willing to forget about space-time foam[26] or grand unification[21)-25)], there are also contributions[27] to scalar boson masses of similar magnitude as before, from the quadratically divergent diagrams of Fig. 1 if one cuts off the loop momenta at $Q = O(M_P)$, namely

$$\delta m_H^2 \left(\sim \delta M_W^2\right) \simeq \Lambda^2 \simeq M_P^2 \tag{14}$$

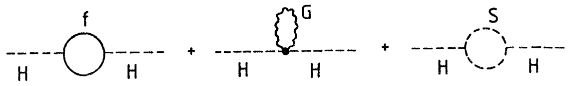


Fig. 1: Quadratically divergent loops contributing to the mass renormalization of elementary scalar fields.

a rather disastrous result. This is the gauge hierarchy problem[25]. Clearly, something goes wrong with our standard model! We are badly in need of some possible mechanism which replaces (14) by something like

$$\delta m_H^2 \left(\sim \delta M_W^2\right) \simeq \Lambda^2 \lesssim O(M_W^2) \tag{15}$$

There are two obvious ways that we can satisfy (15). One scenario is to cut the integration off at momenta O(1) TeV by dissolving the Higgs boson at that scale. Then, Λ is identified with the compositeness scale Λ_c

$$\Lambda_c \lesssim O(1\,\mathrm{TeV}) \tag{16}$$

So, Higgs bosons are not elementary and then by induction, and why not, fermions or even gauge bosons may be composite. This is, in general terms, the COMPOSITENESS framework[28]. An alternative strategy[29] relies on the observation that the boson and fermion diagrams of Fig. 1 have opposite signs. Perhaps they could be made to cancel so that the effective cut-off would be determined by the boson and fermion masses:

$$\Lambda^2 \simeq |m_B^2 - m_F^2| \tag{17}$$

This cancellation persists through the many orders of perturbation theory required to reduce (14) to an acceptable magnitude (15) if the bosons and fermions have identical couplings and similar masses[29),30)]:

$$|m_B^2 - m_F^2| \lesssim O(M_W^2) \tag{18}$$

This requires approximate[29),30)] supersymmetry[31]. This is the SUPERSYMMETRY (SUSY) framework[32].

In either case, the standard model needs some modification, and the very exciting thing is that in either framework, COMPOSITE or SUPERSYMMETRIC, one is bound to find a lot of new "stuff" around M_W, if the gauge hierarchy problem is to be resolved. If we go composite, clearly "excited" states of the known particles cannot be far off because of (16), or if we go supersymmetric, the SUSY partners of the known particles should satisfy (18), thus lurking around and waiting to be discovered. Maybe the new experimental discoveries that I discussed before already hint at new physics around O(100 GeV), as has been expected by theorists for some years now.

Before getting to the physics of the above two general frameworks, let me discuss first the possible extensions of the standard model.

4. - EXTENSION(S) OF THE STANDARD MODEL

The easiest way to classify the possible extensions of the standard model is to look back to its identification card (1) and try to change its content. So, we are naturally led to the following possibilities:

1) Framework: maybe at very short distances like the Planck scale ($\sim 10^{-33}$ cm), quantum gravitational fluctuations cause space-time to become foamy[26] or even fractal, so that the usual notions of relativistic quantum field theory do not apply[33]. One may need modification[33),34)] of quantum mechanics, and such schemes sometimes called[34] unquantum mechanics (UQM) have been already proposed[33),34)]. One of the relevant main characteristics of such schemes[34] is the fact that symmetries do not necessarily and automatically imply conservation laws[34)-36)]; a notion that may eventually be of tantalizing importance.

2) Dimensions: maybe at very short distances like the Planck scale ($\sim 10^{-33}$ cm) space-time dimensions suffer a proliferation. It is not inconceivable that at distances smaller than the Planck scale, there are (d+3)-space + 1-time dimensions, while the extra d-dimensions curl up at the Planck scale so that the "observable" dimensions at large distances are the standard 3+1. Such ideas are realized in the framework of Kaluza-Klein theories[37], which try to unify gravity with all other fundamental interactions[38]. Since others[39] at this conference covered this subject extensively, I will not discuss it further.

3) Basic symmetries: here, one may distinguish two basic philosophies: either at short distances ($\ll M_W^{-1}$), the basic symmetries are reduced or they are extended. In the case of symmetry reduction, one hopes that despite the chaotic behaviour at very short distances, at large distances Poincaré or gauge invariance are progressively established[40]. This is the anti-unification programme[41]. Also here, conservation laws are not valid at very short distances, because there are simply no symmetries present at such small distances. In the more conventional case of symmetry extension, one may distinguish:

a) Extension of the Poincaré algebra, and that brings us uniquely to SUSY[31];

b) Extension of the $SU(3)_C \times SU(2)_L \times U(1)$ gauge symmetry, either by sticking in extra factors like $SU(2)_R$[42],..., or by going to larger gauge groups containing $SU(3)_C \times SU(2)_L \times U(1)$ like SU(5), O(10), E_6,..., i.e., moving to GUTs[21)-25)].

4) Building blocks: maybe at short distances ($< M_W^{-1}$) all or some of the basic fields are composite:

Composite fields:
- a) fermions (quarks-leptons)
- b) vector, scalar bosons (gauge-Higgs).

I believe that the above classification covers more or less all possible attempts for extension of the standard model. I am going to discuss next such attempts in some detail, starting from the end [case 4)] and moving upwards. I should emphasize once more that I have been very liberal with my classification in the following sense.

I take it for granted that we all agree on the excellent performance of the standard model at large distances ($> M_W^{-1}$), so that we all want to keep it valid, at least as an effective theory, at such large distances. Anything that dissolves part or all of the standard model content at distances close to M_W^{-1} is lumped into the COMPOSITENESS framework. Anything that keeps the standard model fundamental, but adds extra group factors or extra elementary particles even at scales close to M_W is lumped into the SUPERSYMMETRY framework. Such a major division emerged naturally from attempts to solve the gauge hierarchy problem. Of course, the possibility remains open that the compositeness framework may need supersymmetry or that the supersymmetric framework may need compositeness but at much smaller distances, like, say, the Planck scale (10^{-33} cm). Still in this case, the first possibility naturally belongs to the compositeness framework while the second one belongs to the supersymmetry framework.

B. COMPOSITENESS

1. - GENERALITIES

It is a historical fact that compositeness has been always invoked when the number of particles once thought "fundamental" become hopelessly large. It is also a sure fact that this procedure has been always highly successful, as it is clearly seen by recalling the happy journey from molecules to atoms to nuclei to hadrons to quarks. All these different stages had a commmon characteristic: the compositeness scale Λ_c was always proportional to the composite particle mass:

$$m_c \gtrsim \mathcal{O}(\Lambda_c) \tag{19}$$

as it is naively expected. Since, once more, we are faced with a plethora of "fundamental" particles (quarks, leptons, gauge bosons, Higgs) it does not take much courage or imagination to invoke once more compositeness. Though this time, things are more complicated; for example, the simple, natural relation (19) is not valid anymore. Available experimental data push the compositeness scale Λ_c at least in the Tev region[43], while the electron or the muon or the photon, etc., do not belong naturally in that mass range, so if they are composite then (19) is wildly violated !

If fermions are composite, one would expect form factor effects, and/or residual interaction effects[44] of the type $(g^2/\Lambda_c^2)(\bar{f}f)(\bar{f}f)$, and/or sizeable contributions[45] to $(g-2)\mu$ from contact interactions of the type $(e/\Lambda_c)(m_\mu/\Lambda_c)\,\bar{f}\sigma_{\mu\nu}fF^{\mu\nu}$. The experimental absence[43] of these effects puts a lower bound on Λ_c:

$$\Lambda_c \gtrsim O(1\,TeV) \qquad (20)$$

barely consistent with (16). Here I have been rather generous by neglecting limits coming from FCNC processes like $\mu \to e\gamma$ or $K_L \to \mu e$, ..., or even proton decay, which are more model dependent. A naive application of compositeness would put a lower bound on Λ_c of 10^6 GeV or 10^{15} GeV, respectively for the case of FCNC or proton decay. We make here the tacit assumption that composite model builders will take care of these problems in a natural way. Still though, there is a large gap between the mass of the electron ($m_e \sim 0.5$ MeV) and the lower bound of Λ_c (>1 TeV) as given by (20). This gross violation of (19) can be interpreted in two ways: either Nature is trying to tell us that there is no further compositeness or that naive compositeness games are over and we should employ more sophisticated schemes. Let us follow the second alternative.

2. - COMPOSITE MODELS

Composite model building[28] varies in the selection of the fundamental blocks of matter, in the dynamics and in the nomenclature! Most people agree to call the fundamental constituents, preons and to use QCD-type dynamics to confine them. Let us follow these semantics. The main problem that all composite model builders face is why $m_f \ll \Lambda_c$? Clearly, some "protection" symmetry should exist that enables some composite fermions to escape virtually massless. Two obvious "weight watchers" recipes come to mind. The first one has to do with the existence of some global, chiral, non-spontaneously broken, anomalous symmetry which, when combined with the assumed preon confinement, leads to massless composite fermions[46]. The need of chiral symmetry for massless fermions is obvious; the need for an anomalous, and thus because of renormalizability, global symmetry is necessary in order to entail the existence of physical massless particles, which cannot be Goldstone bosons since the symmetry is non-spontaneously broken, or cannot be preonic fermions since they are assumed confined, thus they have to be composite massless fermions. Actually, in this case the preonic and composite particle content is related, because they are bound to satisfy the t'Hooft anomaly matching conditions[46] $(\mathrm{Tr}\ T_i^3)_{preon} = (\mathrm{Tr}\ T_i^3)_{compos.}$. These conditions have turned out to be rather cumbersome to satisfy[28]. The second recipe for massless composite fermions has to do with supersymmetry. Here, one invokes[47] the existence of some global continuous symmetry G and supersymmetry which is spontaneously broken to some symmetry $H(\subseteq G)$ while supersymmetry remains unbroken. Then, because of supersymmetry the Goldstone bosons who "live" in G/H should be accompanied by the corresponding quasi-Goldstone fermions, which are then identified as quarks and leptons[47]. Interestingly enough, in most cases, the group H is chiral and all the conditions of the first recipe are met, so that a "double" protection mechanism is at work[48]. It is amusing to notice the "opposite" use of supersymmetry in the compositeness and supersymmetric framework. In composite models, one can easily have Goldstone bosons but needs supersymmetry to ensure the existence of massless spin-1/2 partners of the Goldstone bosons identifiable with the normal quarks and leptons; while in the SUSY framework, one has elementary massless fermions but needs supersymmetry to ensure the existence of "massless" spin 0 partners, identifiable with the electroweak Higgs bosons. Alas, despite the generality of the above mechanisms, and a lot of effort[48),49], no satisfactory model has emerged as yet. There is no model with even a pseudo-realistic mass spectrum because it seems extremely difficult to break eventually the "protection" symmetries such that the composite fermions get some mass[49].

Another serious problem[28] that composite models face is the relation between the Fermi constant G_F and Λ_c. If indeed, $SU(2)_L \times U(1)$ is a fundamental gauge symmetry at the preonic level, then in order to avoid serious residual effects discussed above, we had better have $G_F\Lambda^2 \gg 1$. On the other hand, since most surely the Higgs particle is also composite, one expects $\langle H\rangle^2$ $(\sim G_F^{-1}) \simeq \Lambda^2$ or $G_F\Lambda^2 \simeq 1$, in contradiction with what we just found above. One needs here some rather unnatural arrangements, if at all possible, to satisfy everything. An agnostic point of view[50] would be to identify the residual interactions with the observed $SU(3)_C \times SU(2)_L \times U(1)$ interactions, asserting cynically that everything looks accidentally like a gauge theory ! Any connection between $G_F^{-\frac{1}{2}}$ and Λ_c is lost. In this case, gauge bosons are necessarily composite[50] and our beloved standard model is degraded to the level of the nucleon-nucleon interaction through the exchange of ρ-particles. One now simply "reads" quarks and leptons instead of nucleons and $W^\pm$, Z, γ, G instead of ρ-particles. Personally, I find this particular approach ad-hoc, contrived and ugly and I sincerely hope that Nature has followed a different way.

Till now, I have discussed composite fermions, or composite gauge bosons, so what about composite Higgs bosons ? After all, as I previously emphasized, a possible solution of the gauge hierarchy problem involves only composite Higgs. Well, the situation here is rather bleak. The simplest way[51)] to run the composite Higgs scenario is to replace altogether the Higgs particle by some fermion-antifermion bound state which condenses and gets some v.e.v. of $\langle f\bar{f}\rangle \sim (10^2\text{-}10^3\ \mathrm{GeV})^3$. The fermions that make the condensate may either be some new type of quarks, techniquarks that "feel" a new strong gauge force, called technicolour[52)], or they may be fermions belonging to big representations of the ordinary $SU(3)_{colour}$[53)]. In the first case, the <u>technicolour scenario</u>[52),54)], the technicolour coupling constant becomes unity around $\Lambda_{TC} \simeq$ (100-1000) GeV, so that $\langle\bar{Q}_T Q_T\rangle \simeq \Lambda^3_{TC} \sim$ (100-1000 GeV)^3 as demanded before. Despite its conceptual simplicity and elegance, this picture has suffered fatal blows[54)]. In order to give masses to fermions, one has to introduce further interactions called extended technicolour[55)]. Then, one is drawn into huge problems arising from FCNC[56)]. There is an impasse: either nearly massless fermions or FCNC of at least three to six orders of magnitude larger than experimentally allowed[56)]. FCNC seems to be the Achille's heel of (extended) technicolour.

Furthermore, one predicts[57)] some charged pseudo-Goldstone bosons with masses below 15 GeV, which have not been seen experimentally[43)]. Actually, the observed[10)] canonical decay of $t \to b$ implies that $m_{PG^\pm} > m_t - m_b \simeq 0$ (35 GeV).

In the second case[53)], by using big representations of $SU(3)_C$, one may be able to create condensates with an effective $\Lambda \gg \Lambda_{QCD}$, because of the bigger $SU(3)_C$ charge of the big representations. It is claimed[53)] that in certain cases, there are no problems with either the existence of uncontrollable FCNC or with the existence of undesirable low mass lying charged pseudo-Goldstone bosons, though one may run into problems with the premature breaking of perturbation theory, inability for grand unification and probably the inescapable existence of infra-red fixed points which may jeopardize the whole picture.

Despite the lack at present of any even semi-realistic composite model of fermions and/or gauge bosons and/or Higgs bosons, it is worth discussing some of their immediate experimental consequences.

3. - <u>EXPERIMENTAL HINTS (?) FOR COMPOSITENESS</u>

If indeed the compositeness scale Λ_c is low enough (~ 100 GeV - few TeV), then one naturally expects a rich structure around Λ_c. Composite fermions would imply the existence of excited leptons and quarks, while composite gauge bosons would make highly probable the existence of spin-0 partners or large "anomalous" couplings between themselves. Composite Higgs would support the existence of "Higgs"-like particles but with very different (in principle larger) couplings to other particles from the standard elementary Higgs ones, and/or the existence of very characteristic pseudo-Goldstone bosons. Surely enough, the rich structure of composite models has been vigorously exploited in order to explain the new experimental discoveries mentioned in the beginning. Indeed, if the radiative Z^0 decays (6) are not some caprice of QED bremsstrahlung, then the observed large rate indicates some Z-de-excitation mechanism. The three possible ways are shown in Fig. 2. As can be seen from Fig. 2, the Z-de-excitation may occur through a spin-1, virtual"Z"[58)] (Fig. 2a), or a spin-$\frac{1}{2}$ excited lepton[59)] (Fig. 2b), or finally through a spin-0 particle (Fig. 2c). Obviously, all these processes find their way into composite models. As I discussed above, the anomalously large Z"Z"γ coupling needed in Fig. 2a may easily occur in the composite gauge boson picture[50)], which also may freely provide the X_0 particle needed in Fig. 2c, as a scalar composite partner of the Z^0. On the other hand, the excited lepton needed in Fig. 2b signals for composite fermions. Needless to say, model builders have arranged couplings and masses in such a way as to explain the gross features of the data. Before we get to an euphoric state though, let us have a closer look at these explanations. The ℓ^*[59)] (Fig. 2b) or X_0 (Fig. 2c) explanations ask for bumps in ($\ell\gamma$) or ($\ell^+\ell^-$) invariant masses, which do not seem to be there[12)], while the virtual "Z" explanation[58)] (Fig. 2a) is clearly free from this problem. Furthermore, all three explanations cannot naturally justify the strange kinematic configuration of the events, i.e., the small angle between the photon and one of the leptons. All of them prefer "Mercedes" type configurations. Another problem for all explanations is related to $(g-2)_{\mu,e}$. If in Fig. 2 we "virtualize" the decaying Z particle by suitably hooking its free end and creating a loop, then we get

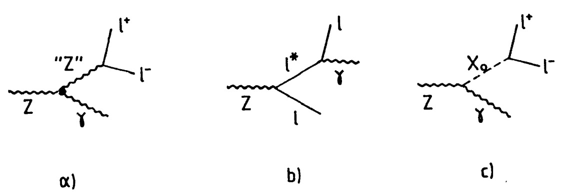


Fig. 2: Z-de-excitation through a) virtual "Z", b) excited lepton ℓ^* and c) scalar boson X_0.

a rather large $(g-2)_{\mu,e}$ contribution[61]. Special care should be taken and the constraints are rather severe[61]. The X_0 inter pretation[60] suffers another blow[62]. The absence of $X_0-\gamma$ interference effects in e^+e^- physics[43] puts a lower bound on its mass[62] $m_{X_0} \gtrsim$ (48 GeV), above one of the observed $(\ell^+\ell^-)$ invariant masses at 42.7 ± 2.4 GeV, and thus making the X_0 interpretation rather improbable. In the case of ℓ^* interpretation[59], extra care should be taken for FCNC effects, which again may become cumbersome. But the excited fermion interpretation still has some further interesting experimental consequences. If $Z \to \ell^{*}\ell^-$, then why not $Z \to \nu^*\nu \to \gamma$ + missing p_T, which may explain the "photon" type of Zen events. Once again, the $\nu\bar{\nu}$ pair should come out almost parallel (extremely small angle), a rather peculiar kinematic configuration[12]. If excited leptons exist, then most probably excited quarks (q^*) should also exist and they may be abundantly produced[63] in $p\bar{p}$ collisions, if they are not too heavy. For example[63], $q+g \to q^* \to qg$; $q\gamma$, which will create bumps in the (jet-jet) or (γ-jet) invariant mass distributions, not unlike the UA(2) ones, if $m_{q^*} \sim$ O(150 GeV). If indeed the q^* is heavier than the $W^\pm$, Z^0, then $q+g \to q^* \to qW$, which will create bumps in the (W-jet), invariant mass distributions not unlike the UA(2) Wen events. Furthermore, the process $q+g \to q^* \to qZ$ may create either bumps in the (Z-jet) invariant mass distribution, or through the frequent $Z \to \nu\bar{\nu}$ decays may create jet + missing p_T type events, not unlike the UA(1) Zen events. So the excited fermion scenario[59],[63], despite its theoretical drawbacks, may provide a phenomenological framework to encompass some of the recent exciting experimental discoveries. In the case of the composite Higgs scenario, one may envisage[64] anomalously large composite Higgs production, through the standard two-gluon fusion mechanism[65]. If one does not bother[64] with the fact that we need an amplitude 1000 times larger, than the elementary Higgs case[65], then things may look fair. The composite Higgs decays mainly to two gluons and that produces the UA(2) jet-jet bump if $m_H \sim$ (150 GeV), while its decays to γZ, Z"Z", W"W" are responsible for the Zen or Wen events. Similarly, one may invoke[66] the production through, say, $q\bar{q}$ or gg annihilation of colour octet mesonic states (pseudo-Goldstone bosons ? Spin-0 coloured composite partners of weak gauge bosons ?) decaying to qq, gg, gZ, gW, $g\gamma$ and thus being responsible for invariant (jet-jet) mass bumps, Zen or Wen or γ + missing p_T events, if $m_8 \simeq$ (150 GeV). Incidentally, it has been recently pointed out[67] that, if the electroweak Higgs mechanism is due to the chiral condensate of big representations of $SU(3)_C$ discussed before[53], then chiral anomalies produce anomalous rates for some weak boson decay modes involving a final-state photon. That might explain[67] either the radiative Z^0 events and/or the "photon" Zen events.

In conclusion, the compositeness framework may provide in principle a phenomenological framework to discuss the recent unconventional experimental discoveries, though there are quite a few bothersome points.

1) There is not even a pseudo-realistic model to provide a theoretical basis for compositeness. The fermion mass spectrum or the family problems are still with us. On the phenomenological front, there are mechanisms here and there to explain bits of experimental data, but no non-problematic unified description has yet emerged.

2) The relation between G_F and Λ_c is either remote or problematic.

3) No good theoretical reason why Λ_c has to be so low ($\Lambda_c \simeq M_W$) and not, say, $\Lambda_c \simeq M_{P\ell}$, which may sound more natural as $M_{P\ell}$ is the ultimate mass scale in microphysics.

Still, for me the clean signal for "low scale" compositeness ($\Lambda_c \sim M_W$) will be the persistence of the high rate (20-25%) of the radiative Z^0 decays (6) in the future experimental runs. Happily enough, we will know very soon.

Next, I come to discuss the supersymmetric framework. Here one takes an "elementary" (anti-composite) point of view: if there is any compositeness at all, it is pushed at superhigh energies, say, $\Lambda_c > O(M_P)$. We have the right to consider quarks, leptons, gauge bosons and Higgs bosons as "elementary" if we are discussing physics up to the Planck scale, which is exactly my modest intention. Let me stress once more that in such a case, supersymmetry is unavoidably needed to cure the gauge hierarchy problem.

C. GRAND UNIFIED THEORIES (GUTs) - SUPERSYMMETRY (SUSY)

GUTs

The spectacular success of the electroweak unification[1] opened the way to more ambitious projects. GUTs[21)-25] that unify electroweak and strong interactions are such a project. The general strategy is expressed schematically by

$$G \underset{M_X}{\longrightarrow} SU(3)_c \times SU(2)_L \times U(1) \underset{M_W}{\longrightarrow} SU(3)_c \times U(1)_{e\text{-}m} \quad (21)$$

One assumes the existence of a large group G [like SU(5), O(10), E_6, ...] which at some superhigh energy scale M_X suffers spontaneous breakdown (SB) to the "standard" group which eventually is spontaneously broken at M_W to $SU(3)_c \times U(1)_{e.m.}$. The consequences of such a natural and simple assumption are rather dramatic:

1) At energy scales above $O(M_X)$, there is only one gauge coupling constant g_G, which at lower energies, because of the renormalization effects[23] which are different for strong and electroweak interactions, give rise to the three observed gauge coupling constants g_3, g_2, g_1. In such a case, one is able to predict[23)-25),68] the value of the electroweak mixing angle θ_W at low energies

$$\sin^2 \hat{\theta}_W (M_W) = 0.214 \begin{smallmatrix} +0.004 \\ -0.003 \end{smallmatrix} \quad (22)$$

in spectacular agreement with the experimental value[4),10]

$$\sin^2 \hat{\theta}_W (M_W) \Big|_{Exp.} = 0.215 \pm 0.010 \pm 0.007 \quad (23)$$

Also, one is able to calculate[23)-25,68] M_X, which is the mass of the superheavy gauge bosons [$G/SU(3)\times SU(2)\times U(1)$] in simplest models[22]

$$M_X = (1.5 \pm 0.5) \cdot 10^{15} \Lambda_{\overline{MS}} \quad (24)$$

with $\Lambda_{\overline{MS}}$ = 0.1 to 0.2 GeV.

2) Quarks and leptons lie necessarily in the same multiplet(s). That means:

i) relations between their internal quantum numbers, e.g., commensurate electric charge (charge quantization) or common weak isospin, etc.

ii) Relations between their masses, which start identical or very similar at superhigh energies but because of renormalization effects become different at low energies. One is able to predict[24),69], for example

$$\frac{m_b}{m_\tau} \simeq 2.8\text{-}2.9 \quad (25)$$

in excellent agreement with experiment.

iii) Baryon and lepton number violating interactions. A clear manifestation of lepton number violating interactions will be the possible existence of neutrino masses. One predicts[70] in GUTs

$$m_\nu \simeq \frac{m^2}{M} \simeq (10^{-5} \to 100)\, eV \quad (26)$$

where m indicates some low energy mass while M stands for some superhigh energy scale (M_X, M_P ?). In principle, the neutrino mass spectrum is such that neutrino oscillations are allowed. A dramatic manifestation of the baryon number violating interactions will be the nucleon instability and/or n-$\bar{n}$ oscillations[71]. It should be categorically stressed that while the nucleon instability is rather compulsory, either the neutrino masses or observable n-$\bar{n}$ oscillations are rather optional. Indeed, the minimal SU(5) model[22] does not predict either ν-masses or observable n-$\bar{n}$ oscillations. On the other hand, the minimal SU(5) model[22] or other minimal GUTs do predict proton decay mainly to $e^+\pi^0$, with a rather "short" lifetime[25]

$$\tau(p \to e^+\pi^0) = \left(\frac{M_X}{4 \cdot 10^{14}\, GeV}\right)^4 \cdot (1\text{-}2) \cdot 10^{29} \text{ years} \quad (27)$$

which by using (24) reads as

$$\tau(p \to e^+\pi^0) \simeq 10^{29 \pm 1} \text{ years} \quad (28)$$

in rather sharp contradiction with the presently available experimental lower bound[72]

$$\tau(p \to e^+\pi^0) \geq 2 \cdot 10^{32} \text{ years} \quad (29)$$

There are several ways out from this grave difficulty suffered by the minimal SU(5)[22] or other minimal GUTs. The central aim of these "gerontologic" attempts[73)-75] is to increase M_X. Since M_X is determined by the (differential) renormalization group equations (RGE), one may try either to change coefficients (parameters) of the equations or the boundary conditions. In the first case, one[73] throws in superfluous "light" stuff coming, say, from some relics of 10 + $\overline{10}$, 45 or whatever you please, with the net effect of slowing down the rate of evolution of the coupling constants, and thus by postponing the "meeting" of the three coupling constants to increase suitably M_X. This mechanism is an imitation what happens in SUSY GUTs, but there[76)-78], as we will see soon, it happens naturally and it does not look so ad-hoc, contrived and unaesthetic as it looks here

In the second case[74], one tries to find ways to modify the boundary conditions of the RGE. One may think[74] that non-renormalizable interactions coming either from gravity or from Kaluza-Klein theories may modify the gauge boson kinetic terms[74]. In such cases, proper normalization of the gauge boson kinetic terms will, in principle, induce changes in the "naive" gauge coupling constants relations at M_X. For example, one may envisage[74] non-renormalizable terms of the form $\varepsilon(\phi/M_P)F^{\alpha}_{\mu\nu}F^{\alpha}_{\mu\nu}$, which after the GUT Higgs field ϕ gets its v.e.v., will necessitate modification of the gauge couplings relations at M_X in order that the GUT gauge field strength $F^{\alpha}_{\mu\nu}$ behaves canonically. In such cases, one finds[74]

$$\sin^2\theta_W(M_X) = \frac{3}{8}(1 - A\varepsilon)$$

and (30)

$$\ln\frac{M_X}{M_W} = (1 + B\varepsilon)\ln\frac{(M_X)_{canon.}}{M_W}$$

with A and B positive, calculable constants. We notice that there is a tendency for M_X to increase with respect to its canonical (ε=0) value $(M_X)_{canon}$, where $\sin^2\theta_W$ has a tendency to decrease its canonical, boundary value of 3/8. Fortunately enough, there are ranges for the "arbitrary" parameter ε, which suitably increase τ_p without a catastrophic decrease in $\sin^2\theta_W$. I find this alternative[74] interesting and worth keeping in mind for other applications as well. Lastly, one may try to exploit[75] the rich structure of big groups like O(10), E_6, ... O(18) in the hope of finding tricks to postpone proton decay. Indeed, several such tricks have been proposed[75], not very appealing and at a rather high price.

All in all, "gerontology" does not look to me very convincing and I do believe that the real elixir for protons has one name

SUPERSYMMETRY

1. Generalities

Supersymmetry is a new kind of symmetry[31] relating bosons to fermions:

$$Q|F\rangle = |B\rangle \quad , \quad Q|B\rangle = |F\rangle \quad . \tag{31}$$

It is generated by spinorial charges Q_α (α a spinor index) which, as you might expect for fermionic operators, obey an anticommutation algebra:

$$\{Q^i_\alpha, Q^{+\dot\alpha}_j\} = -2(\sigma^\mu)^{\dot\alpha}_\alpha P_\mu \delta^i_j \; : \; i = 1, \ldots, N. \tag{32}$$

It has been proven[79] that this is the only possible way of combining internal symmetry [the index i in Eq. (32)] with Lorentz invariance. A theory is said to have simple supersymmetry if N = 1 and extended supersymmetry if N > 1. Particles in renormalizable gauge theories can only have helicities between +1 and −1, and since each charge Q changes helicity by ½ a unit:

$$h = +1 \xrightarrow{Q} +\tfrac{1}{2} \xrightarrow{Q} 0 \xrightarrow{Q} -\tfrac{1}{2} \xrightarrow{Q} -1 \tag{33}$$

there can be at most four extended supersymmetries in a SUSY gauge theory. The analogous argument applied to a theory including gravity, in which case the allowed helicities range from +2 to −2, tells us that up to 8 extended supersymmetries are possible in a supergravity theory[80]. Since general relativity expresses invariance under general co-ordinate transformations which allow different local translations at each point, the supersymmetry transformations must also be local in supergravity theory, just as gauge theories embody local phase transformations. On the other hand, gauge theories can only accommodate global SUSY. In what follows, we will mainly restrict ourselves to simple N = 1 SUSY, in which case the permissible supermultiplets are the graviton smultiplet which includes a spin-3/2 gravitino, and

$$\text{gauge:} \begin{pmatrix} 1 \\ 1/2 \end{pmatrix} \qquad \text{and chiral:} \begin{pmatrix} 1/2 \\ 0 \end{pmatrix} \tag{34}$$

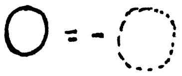
smultiplets. Conventional gauge interactions of spin-1/2 fermions such as quarks are accompanied by interactions involving the supersymmetric partners of fermions and gauge bosons:

$$g(\bar f \gamma f G_\mu) \longleftrightarrow \sqrt{2}\, g\,(\bar f \tilde f \tilde G + h.c.) \tag{35}$$

while Higgs Yukawa interactions are accompanied by interactions involving the fermionic partners of Higgses:

$$\lambda(f f H) \longleftrightarrow \lambda(f \tilde f \tilde H) \tag{36}$$

Such simple relations like (35) and (36) between different sets of coupling constants, as well as similar relations between fermion and boson masses, lie behind the magic properties of SUSY field theories. SUSY field theories are extremely ultraviolent convergent field theories, thanks to the cancellations between suitably "dressed" fermion and boson loops, based on the simple observation

 (37)

which should be the main theme of any SUSY flag. A mathematical expression of these cancellations is given by the no-renormalization theorems[81)], which tell us that if a term is absent from the superpotential at the tree level, then it cannot be generated in any order in perturbation theory. This "set it and forget it" principle is the reason that SUSY has been employed to solve all types of hierarchy problems in particle physics. I discussed in detail before [see Fig. 1; Eqs. (17) and (18)] how SUSY may stabilize "light" Higgs masses and therefore solve[29)], at least in part, the cumbersome gauge hierarchy problem. Another serious hierarchy problem, the strong CP hierarchy problem[82)] may also find its solution[83)] in the SUSY framework. Simply recall that the non-perturbatively created $\theta_{QCD} F^{\alpha}_{\mu\nu} \tilde{F}^{\alpha}_{\mu\nu}$ term, where θ is a free parameter and $F^{\alpha}_{\mu\nu}$ the gluon field strength, violates CP and P and thus contributes to DEMON (d_n) and thus θ_{QCD} should be[84)]

$$\theta_{QCD} \lesssim 10^{-9} \qquad (38)$$

if the bound (3) has to be satisfied. This is a rather small number! Once more, SUSY may rescue[83)] the situation. Starting with $\theta_{QCD} = 0$, the radiative corrections to quark masses, which eventually create catastrophically large contributions to θ_{QCD}, behave themselves and thanks to the fermion-boson loop cancellations any possible $\delta\theta_{QCD}$ always satisfies naturally Eq. (38)[83)]. You may imagine diagrams like the ones in Fig. 1, but with fermions at the external legs instead of bosons. The principle is the same as in the gauge hierarchy problem: "set it and forget it". SUSY <u>respects</u> and <u>stabilizes</u> hierarchies[29),83)].

A remarkable implication of the magic convergent properties of SUSY theories is the existence of non-trivial finite theories. Yes, there are four-dimensional SUSY Yang-Mills (Y-M) field theories, like N = 4 or N = 2 with vanishing one-loop β-function(s) that are <u>finite</u>[85),86)]. The implications, especially for quantum gravity, may be rather dramatic. The following table (Table 1) shows the convergent properties of SUSY theories.

Having established the credentials of SUSY field theories, I move next to discuss their physics applications.

	Y-M Field Theor.	SUSY (Y-M) Field Theor.			Supergravity
		N = 1	N = 2 [86)] $(\beta)_{1\text{-loop}} = 0$	N = 4 [85)]	N = 8
Divergences	$\ln\Lambda$, Λ^2, ...	$\ln\Lambda$	FINITE		?

Table 1

Divergent(less) field theories

2. - SUSY PHENOMENOLOGY

i) Low energy S-physics

It is an unfortunate fact of life that none of the known elementary particles can be the supersymmetric partner of any other. The expected spectrum of supersymmetric particles is illustrated in Table 2.

Particle	Spin	Sparticle	Spin
quark q	½	squark $\tilde{q}$	0
lepton ℓ	½	slepton $\tilde{\ell}$	0
photon γ	1	photino $\tilde{\gamma}$	½
gluon g	1	gluino $\tilde{g}$	½
W	1	Wino $\tilde{W}$	½
Z	1	Zino $\tilde{Z}$	½
Higgs H	0	shiggs $\tilde{H}$	½
graviton	2	gravitino	3/2

Table 2

Supersymmetric particles

As we have repeatedly heard at this conference[10),87),88)], intense experimental searches have put (if not discovered[10)]?) interesting lower bounds on sparticle masses[87),88)]. The absence of any charged particle at PEP and PETRA means that squarks, sleptons, charged winos and shiggses must have masses above O(20) GeV[87),88)]. The small number of the UA(1) "Zen" (monojet) events has been used[13),10)] to set a lower bound of the order of 40 GeV on the squark and gluino masses. This limit is already a factor of O(2) better than the PETRA/PEP bounds for squarks, and a factor of (10-20) better than the beam dump bounds[89)] for gluino masses. Negative results from a) $e^+e^- \to \tilde{\gamma}\tilde{\gamma}\gamma \to \gamma$ + nothing and b) $e^+e^- \to \tilde{\gamma}\tilde{Z}$ searches[88),87)] imply $m_{\tilde{e}} > O(20;30)$ GeV for $m_{\tilde{\gamma}} < O(6,1)$ GeV in case a)[88)] while $m_{\tilde{e}} > O(100,30)$ GeV for $m_{\tilde{Z}} < O(10,40)$ GeV in case b)[87)]. No experimental lower bounds exist on colourless and electrically neutral sparticles. There are several possibilities[90)] for the lightest neutral sparticle, including the sneutrino $\tilde{\nu}$, the photino $\tilde{\gamma}$, a neutral shiggs $\tilde{H}$

and the gravitino $\tilde G$. In most models the sneutrino is not lighter than the charged sleptons, while the gravitino is expected to have a mass comparable to the sleptons in the currently favoured class of supergravity models discussed later on. Plausible possibilities in many models are that the $\tilde\gamma$ and/or the neutral $\tilde H$ may be considerably lighter than the other sparticles. You can see from the forms of the couplings (35) and (36) that sparticles must be pair-produced and there must be another sparticle among any sparticle's decay products. This means that the lightest sparticle is presumably stable, and should be present in the Universe today as a relic from the Big Bang. Conventional Big Bang cosmology and the continued expansion of the Universe impose interesting constraints[91] on the possible light supersymmetric particles. One finds[90),91)] that probably

$$m_{\tilde H} \gtrsim \min(m_{\tilde \ell}, m_{\tilde\gamma}), \quad m_{\tilde\gamma} \geq O(1/2)\,\mathrm{GeV} \tag{39}$$

Interestingly enough, stable, massive photinos are very welcome in cosmology. If indeed "cold", non-baryonic matter, as demanded by viable galaxy formation and clustering models, is responsible for the dark matter then stable, massive photinos are a very plausible candidate[92]. Accretion of photinos into galactic halos is unavoidable in a photino dominated Universe. Then, observable annihilation products of photinos in our halo include γ-rays, cosmic ray positrons and, most significantly low energy cosmic ray antiprotons[93]. For scalar masses O(50) GeV and a photino mass of O(3) GeV, plausible assumptions about the parameters of our dark halo yield[93] an antiproton flux comparable to that observed[94] below the threshold for secondary production in standard cosmic ray propagation models. A very interesting result indeed[93]. Cosmological hints for SUSY ?

In any way, present experimental searches already start probing very interesting mass ranges from a SUSY point of view. We may know soon. Nevertheless, if SUSY indeed resolves the gauge hierarchy problem, then (18) has to be satisfied. The "doubling" of the known particles, as given in Table 3, has to occur at accessible energes $O(M_W)$. In such a case, one should worry about the sparticle effects on the established low energy phenomenology. As advertised in the beginning of the talk and proved later in the case of composite models (technicolour,...), FCNC, g-2, ρ, ..., all act as Damoclios sword on any low energy extension of standard model. Detailed studies of sparticle contributions to "known" physics have put constraints on the form of the sparticle mass spectrum. The most severe constraints come from the study[95] of the K_1-K_2 system. Flavour changing neutral gaugino interactions, if present, would make disastrously large contributions to the K_1-K_2 mass difference, as shown in Fig. 3.

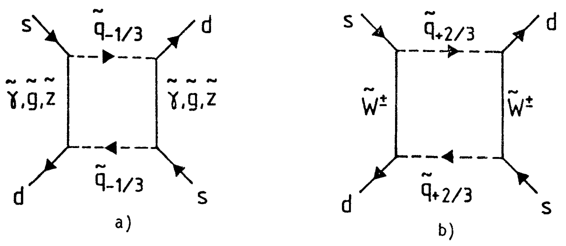


Fig. 3: a) Neutral gaugino and b) charged gaugino SUSY box diagrams contributing to the K_1-K_2 mass difference[95].

To avoid this problem, we must demand that the $m^2_{\tilde q}$ matrix be a simple function of the quark mass matrix m_q, presumably a quadratic function:

$$m^2_{\tilde q} = \tilde m^2 \mathbb{1} + C_1 m_q \tilde m + C_2 m_q^2 \tag{40}$$

By making the ansatz (40), we also avoid any problem[95] with $\tilde W^\pm$ box diagrams. Equation (40) guarantees that

$$m^2_{\tilde c} - m^2_{\tilde u} = O(1)\left(m_c^2 - m_u^2\right) \tag{41}$$

if the $C_i = O(1)$, and therefore, the $\tilde W$ box diagrams will be suppressed by a super-GIM mechanism to

$$O\left(\frac{\alpha G_F \sin^2\theta_c}{4\pi's}\right)\left(\frac{m^2_{\tilde c}-m^2_{\tilde u}}{m^2_{\tilde W,\tilde q}}\right) = O\left(\frac{\alpha G_F \sin^2\theta_c}{4\pi's}\right)\left(\frac{m^2_c-m^2_u}{m^2_{\tilde W,\tilde q}}\right) \tag{42}$$

which is of the same order as the $W^\pm$ box diagrams if m_W and $m_{\tilde q} = O(m_W)$.

Other constraints coming from g-2[96], ρ[97], ..., are rather easily satisfied. For example, a smuon mass above 20 GeV, as already entailed by the PEP/-PETRA limits, suffices[96] to make acceptable contributions to $(g-2)_\mu$, in a large class of models. Furthermore, since some of these constraints are rather model dependent, I will come back to them when I discuss SUSY models. What should not escape our attention is the major fact that despite the plethora of SUSY particles at low masses $O(M_W)$, it is possible to avoid all the low energy physics traps!

ii) SUSY GUTS phenomenology

In minimal SUSY GUTs[30),98)] where the only new particles with masses $O(M_W) \ll M_X$ are the spartners of the conventional $SU(3)_C \times SU(2)_L \times U(1)$ particles, one finds[76)-78)]

$$\sin^2\hat{\theta}_W(M_W) = 0.236 \pm 0.003 \tag{43}$$

and

$$M_X = 6\cdot 10^{16}\,\Lambda_{\overline{MS}} \simeq 10^{16}\ GeV \tag{44}$$

while m_b/m_τ remains numerically unchanged[77),78)] from its value (25) in conventional GUTs. The large value (44) of M_X would yield a very long nucleon lifetime if the decay amplitude were $\propto 1/M_X^2$ and the lifetime $\propto M_X^4$ as in conventional GUTS, (27). However, it has been realized[78),99),100)] that in minimal SUSY GUTs, the exchange of a heavy Higgs colour triplet H_3 smultiplet can give $\Delta B = 1 = \Delta L$ interactions which are $\propto 1/M_{H_3} = 0\,(1/M_X)$, suggesting a dependence $\propto M_X^2$ of the baryon lifetime. The tree level H_3 exchange gives an interaction of dimension 5 involving two quarks or leptons and two of their spin-zero spartners. This interaction must be dressed mainly by $\tilde{W}$ exchange[78)] [$\tilde{g}$ or $\tilde{\gamma}$ exchanges[100)] are usually suppressed[101)] except some pathetic, fine-tuned cases[102)]] to give a $\Delta B = 1 \doteq \Delta L$ four-fermion interaction (see Table 3, first column). Colour symmetry and quark mass factors in the H_3 Yukawa couplings favour the participation of second and third generation particles, so that the dominant nucleon decay modes in minimal SUSY GUTs are[78),100)]

$$p \to \bar{\nu} K^+,\quad n \to \bar{\nu} K^0 \tag{45}$$

Taking the experimental lower limit $\tau(n\to\bar{\nu}K^0) > 0(10^{31}$ years), one finds[78),103),104)] that if $m_{\tilde{q}} \simeq m_{\tilde{\ell}} \simeq m_{\tilde{W}} \simeq 100$ GeV then

$$m_{H_3} \gtrsim 0\,(10^{16}\text{-}10^{17})\ GeV \tag{46}$$

not yet in disagreement with (44). Nevertheless, it does appear that baryon decay <u>should be expected soon</u> in the context of <u>minimal SUSY GUTs</u>.

There is an <u>alternative</u> form[101)] of SUSY <u>d = 5 operator</u> which cannot be simply obtained from superheavy triplet Higgs(ino) exchange, but could be generated, for example, by dynamics at the Planck scale[101)] (see Table 3, second column). In this case, one finds[101)] that the dominant nucleon decay modes are

$$n \to \bar{\nu} K^0,\quad p \to \bar{\nu} K^+;\ \mu^+ K^0 \tag{47}$$

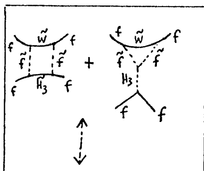
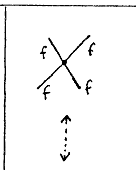
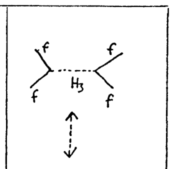
with some non-negligible contamination of e^+K^0 in certain cases.

It is worth noting that non-minimal SUSY GUTs can be made to accommodate radically different baryon lifetimes and decay modes. For example, with a symmetry to suppress the dimension 5 H_3 exchange diagram, one can tolerate[105)] H_3 as light as 10^{10-11} GeV [and may desire them for cosmological baryosynthesis[106)]] in which case the dominant <u>dimension 6</u> $\Delta B = 1 = \Delta l$ interactions due to relatively "light" H_3 exchange (see Table 3, third column) would yield[105)-107)]

$$p \to \bar{\nu} K^+;\ \mu^+ K^0,\quad n \to \bar{\nu} K^0 \tag{48}$$

More exotic baryon decay modes [$\mu^- K^{107)}$, or even $e^+\pi^0$[108)] at an acceptable rate] are also obtainable with ingenuity.

We conclude from the above analysis that while the present lower limits on baryon decays into strange decay modes push very hard on the d = 5 Higgsino exchange or d = 6 Higgs exchange model, there are not yet excluded experimentally. Looking on the optimistic side, most SUSY models suggest that baryon decays into <u>strange particles</u> should be detectable in the <u>near future</u>. Presumably, $N \to K+\ldots$ should be a major theme in any SUSY flag. Furthermore, chiral Lagrangian calculations[109)] which embody current algebra and PCAC predict characteristically different ratios of rates[104)] for $p \to \bar{\nu}K^+$, $n \to \bar{\nu}K^0$ and $p \to \mu^+K^0$ in the different models, as is shown in Table 3. We may know soon if indeed there is a super-elixir for protons.

	 d = 5, $\tilde{H}_3$, H_3 exchange 78),99),100)	 Alternative d = 5 SUSY operator 101)	 d = 6 H_3 exchange 105)-107)
$p \to \bar{\nu}K^+$	0.55	1.75	1
$n \to \bar{\nu}K^0$	1	1	1
$p \to \mu^+K^0$	$(10^{-3} - 10^{-4})$	0.10	0.55

<u>Table 3</u>

Relative ratios for baryon decays into strange particles[104)]

3. - <u>PHENOMENOLOGICAL SUPERGRAVITY MODELS</u>

i) <u>SUSY models</u>

One may wonder if there are at all realistic SUSY models encompassing all different phenomenological constraints previously mentioned. Indeed, realistic SUSY model building is not an easy task[32)]. However, any effort is worthwhile since SUSY models are left as the only candidates for a physical description of the world, at least up to energies of the Planck scale M_P. Any high standard(s) SUSY model should satisfy the following two SUSY golden rules:

1) It should provide naturally an acceptable form of SUSY breaking such that

$$\underbrace{\mathcal{O}\left((20\,\mathrm{GeV})^2\right)}_{\text{Experim.}} \leq m_B^2 - m_F^2 \equiv \tilde{m}^2 \leq \underbrace{\mathcal{O}\left(M_W^2\right)}_{\text{Gauge Hierar.}} \tag{49}$$

where $\tilde{m}^2$ is a typical boson-fermion mass splitting of a supermultiplet. The sparticle mass spectrum should be such that not only all types of low energy constraints are satisfied [e.g., Eq. (40)] but in addition some possible potential problems of the standard model should find a satisfactory resolution.

2) It should provide a complete solution to the three-fold gauge (scale) hierarchy problem: create, stabilize and dynamically explain the scale hierarchy. All (small) mass scales should be determined dynamically in terms of one fundamental one, the super-Planck scale $M \equiv (M_P/\sqrt{8\pi}) \simeq 2.4 \cdot 10^{19}$ GeV:

$$\frac{M_W}{M} \simeq \frac{\tilde{m}}{M} \simeq \mathcal{O}(10^{-16}) \tag{50}$$

Surprisingly enough, such no-scale models[110)-117)] have recently been constructed. Since their construction involves a lot of very interesting physics, it is worth discussing the different steps that lead (uniquely?) to them.

Since $N > 1$ SUSY theories provide vector-like fermion spectra[118)], let us concentrate on $N = 1$ SUSY theories. Let us try first $N = 1$ global SUSY theories. SUSY breaking should be spontaneous or "soft" such that the exorcized-away quadratic divergences do not reappear. In the case of spontaneous global SUSY breaking, one may distinguish between perturbative[119)] and non-perturbative mechanisms[120)]. While perturbative breaking is clearly possible[119)], non-perturbative breaking is still debatable[120)]. In any case, the construction of realistic models in this framework is almost impossible[29),121)]. Such models are drawn into difficulties[121)] with FCNC, ABJ anomalies, out-of-hand proliferation of low mass $O(M_W)$ particles, etc.; in other words, they are hopeless. The case of soft[122)] global SUSY breaking is certainly possible but it requires horrendeous fine tunings[98)] to satisfy the first SUSY golden rule. Furthermore, both spontaneous (perturbative[119)]) or soft[122)] SUSY breaking can never satisfy[111)] the second SUSY golden rule. Because of the non-renormalization theorems[81)], if global SUSY is unbroken at the tree level, it remains unbroken to all orders in perturbation theory[123)]. So, either $\tilde{m}$ is there at the tree level (and thus not dynamically determined) or is not there at all. Non-perturbative[120)] spontaneous global SUSY breaking may in some way partly satisfy the second SUSY golden rule but it seems to me contrived and that it will have all the other problems of perturbative spontaneous SUSY breaking. If so, the only way left is the framework of local SUSY or supergravity[80),124),125)] (SUGAR). It is quite interesting that phenomenological as well as theoretical constraints have led us uniquely into the SUGAR framework[125)], something highly desirable because it involves only local symmetries and provides automatically unification with gravity [see remarks after Eq. (33)]. Before getting to the physics of $N = 1$ SUGAR, we need some "grammar" of the $N = 1$ supergravity language.

ii) N = 1 supergravity primer[124)-125)]

The general couplings of chiral and vector multiplets (34) to $N = 1$ supergravity is specified by two functions of the complex scalar fields ϕ_i contained in the chiral multiplets[126)]. An analytic function $f_{ab}(\phi) = f_{ba}(\phi)$, related to the Y-M part of the Lagrangian, gives for the kinetic terms of the gauge fields,

$$-\frac{1}{4}\left(\mathrm{Re}\, f_{ab}\right) F^a_{\mu\nu} F^{b\mu\nu} + \frac{i}{4}\left(\mathrm{Im}\, f_{ab}\right) F^a_{\mu\nu} \tilde{F}^{b\mu\nu} \tag{51}$$

(a, b are indices of the adjoint representation of the gauge group G). Then, a real gauge invariant function $G(\phi,\phi^*)$, the Kähler potential, defines the scalar kinetic terms, given by

$$G_j^i \left(\partial_\mu \phi^j\right)\left(\partial^\mu \phi_i^*\right) \tag{52}$$

in the notation

$$G_i \equiv \frac{\partial G}{\partial \phi^i}\,, \quad G^i \equiv \frac{\partial G}{\partial \phi_i^*}\,, \quad G_j^i \equiv \frac{\partial^2 G}{\partial \phi_i^* \partial \phi^j}$$

The kinetic terms have a form characteristic of supersymmetric non-linear σ models. The scalar fields ϕ_i in $N = 1$ SUGAR span a Kähler manifold with G^i_j (52) as its metric[127)]. Clearly, the functions $G(\phi,\phi^*)$ and $f_{ab}(\phi)$ largely determine the physics of the $N = 1$ SUGAR Y-M theories[126)]. Indeed, the scalar potential V has two terms[126),128),129)]:

$$V = V_c + V_g \tag{53}$$

The gauge potential V_g reads

$$V_g = \frac{1}{2}\left(\mathrm{Re}\, f_{ab}^{-1}\right) D^a D^b \tag{54}$$

where the real functions D^a are

$$D^a = g_a\, G_j \left(T^a\right)_i^j \varphi^i \tag{55}$$

(g_a is the gauge coupling constant associated to the normalized generator T^a). The "chiral" potential is

$$V_c = \exp G\left[G_i G^j \left(G_i^j\right)^{-1} - 3\right] \tag{56}$$

It is apparent from the potential (56) that, unlike the global case, spontaneously broken SUSY does not imply $\langle V\rangle > 0$. This is fortunate since one can now obtain spontaneous breakdown of local SUSY (the super-Higgs effect) in Minkowski space ($\langle V\rangle = 0$). The theory contains also a gravitino mass term, $m_{3/2}\, \bar{\phi}_{\mu L}\, G^{\mu\nu} \phi_{\nu L}$, with

$$m_{3/2} = \left\langle \exp \frac{G}{2} \right\rangle \tag{57}$$

The most naïve choice for the functions G and f_{ab} would be the ones that provide canonical scalar and vector kinetic terms,

$$G_i^j = \delta_i^j \quad ; \quad f_{ab} = \delta_{ab} \tag{58}$$

corresponding to a flat Kähler manifold. In such cases, one writes

$$G = \varphi_i \varphi^{i*} + \ln \left|f(\varphi)\right|^2 \tag{59}$$

where $f(\phi)$ stands for the gauge invariant superpotential. The "minimal" choice of G and f_{ab} in (58) has some rather unpleasant consequences. The cosmological constant $\langle V\rangle = \Lambda$ is zero due to some unbearable fine-tuning of the parameters in G. Furthermore, scalar boson masses are proportional[130] to the gravitino mass (57), as they are given by the curvature of the potential (56) at the minimum. In this case, the gravitino mass is essentially a free parameter and because of its relation[130] to scalar boson masses or equivalently to $\tilde{m}$ (4?), it has to be chosen by hand $O(M_W)$. Dynamical determination of $\tilde{m}$ is excluded, no way to satisfy the second SUSY golden rule!

There is, however, a very elegant way to circumvent these two unsatisfactory points: there exist non-trivial Kähler potentials for which the chiral potential V_c is identically zero[131),132)]. Supersymmetry is, however, broken. Vacuum expectation values are not determined by the classical theory, ditto for the gravitino mass. The cosmological constant is naturally zero[131),132)]. As we will see later, radiative corrections are then used[110)-114)] to determine the various scales of gauge symmetry breaking, which in general will be closely related to the gravitino mass. Eureka! This is what we are aiming at to satisfy the second SUSY golden rule. In principle, it is sufficient to require zero chiral potential only in the direction of a gauge singlet complex scalar z field, the Polonyi field. In such a case, we may rewrite (57) as[131)]

$$V_c = 9 \exp\left(\frac{4}{3}G\right) G_{zz^*}^{-1} \frac{\partial^2}{\partial z \partial z^*} \exp\left(-\frac{1}{3}G\right) \tag{60}$$

and

$$V_c \equiv 0 \quad \text{implies} \quad \frac{\partial^2}{\partial z \partial z^*} \exp\left(-\frac{1}{3}G\right) = 0$$

with solution[131)]

$$G = -3 \ln\left(z + z^*\right) \tag{61}$$

The scalar kinetic term $G_{zz^*}(\partial_\mu z)(\partial_\mu z^*)$ is never canonical, and the gravitino mass is[131)]

$$m_{3/2} = \left\langle \left(z + z^*\right)^{-3/2} \right\rangle \tag{62}$$

$m_{3/2}$ is undetermined but non-zero since

$$\left\langle G_{zz^*} \right\rangle = 3 \left(m_{3/2}\right)^{4/3} \neq 0 \tag{63}$$

$V_c \equiv 0$ entails a very particular geometry of the Kähler manifold. The Kähler curvature

$$R_{zz^*} \left(\equiv \frac{\partial^2}{\partial z \partial z^*} \ln G_{zz^*} \right)$$

is given by[131)]

$$R_{zz^*} = \frac{2}{3} G_{zz^*} \tag{64}$$

or $R(\equiv (R_{zz^*}/G_{zz^*}) = 2/3$. Equation (64) means that the Kähler manifold is an Einstein space (maximally symmetric space) i.e., that the scalar field z is a co-ordinate of the coset space SU(1,1)/U(1)[131),111)]. The non-compact global SU(1,1) invariance can be checked explicitly in the whole Lagrangian apart from the gravitino mass term[111),133)]. It is very interesting to notice[131)] that the N = 1 Lagrangian for one chiral multiplet with vanishing potential corresponds, up to the gravitino mass term, to a particular truncation of N = 4 supergravity, which is known[134)] to possess an SU(1,1) non-compact global symmetry. Vanishing chiral potentials for an arbitrary number n of chiral multiplets also exist[131),111)-115),133)].

In one case[112], the scalar fields ϕ_i, $i = 1, 2, \ldots, n$, are co-ordinates of an $SU(n,1)/SU(n)\times U(1)$ coset space, which is an Einstein space with curvature given by (64) but with (n+1) replacing 2 in the numerator.

Clearly, in both cases $n = 1$ or $n > 1$ the "flatness" of the potential implies massless scalar bosons, neglecting radiative corrections. There is some kind of "curvature-conservation" between the Kähler manifold spanned from the scalar fields of the chiral multiplets and the chiral potential V_c, as schematically represented by Fig. 4.

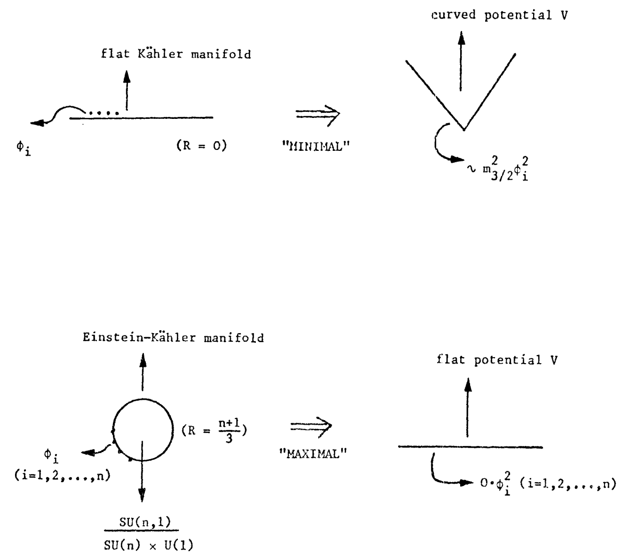


Fig. 4: Relation between the form of the Kähler manifold and the induced scalar (chiral) potential.

It should be stressed that higher than four ungauged extended supergravities also exhibit[135] invariances under non-compact global groups which contain SU(1,1) as a subgroup:

$$N=5:\ SU(5,1);\quad N=6:\ SO^*(12);\quad N=7,8:\ E_{7,7} \tag{65}$$

In addition, extended $N \geq 4$ and gauged $N \geq 2$ supergravities do inevitably contain[134),135] non-minimal kinetic terms for the gauge smultiplets [$f_{ab} \neq \delta_{ab}$ in Eq. (51)]. As a result, tree level gaugino masses are introduced, of the form[126]:

$$m_{\tilde V} = m_{3/2}\,\frac{\langle 0|f'|0\rangle}{\langle 0|f|0\rangle}\cdot\frac{G'}{G} \tag{66}$$

for $f_{ab} = \delta_{ab} f(\phi_i, z)$, and with primes indicating z-differentiation. The usefulness of this remark will become apparent shortly. After this unavoidable digression into the "esoterics" of N = 1 SUGAR, the road is now open to physical applications.

iii) No-scale models

N = 1 supergravity Y-M field theories seem to be the only available framework for physics description below the Planck scale M. Alas, these theories are not renormalizable! This fact should not bother us, since anyway we are going to use the N = 1 SUGAR framework as an effective theory[136] to describe physics for scales below M, according to the general scheme:

$$\mathcal{L}(N=1\ \mathrm{SUGAR}) \xrightarrow[E<M]{} \mathcal{L}(N=1\ \mathrm{SUSY}) + \mathcal{L}_{SOFT} \tag{67}$$

where $\mathcal{L}_{SOFT}$ stands for a highly constrained set of soft SUSY breaking terms. The passage described by (67) is carried out by making a choice for G (52) and f_{ab} (51) such that

1) Supersymmetry is spontaneously broken;
2) Certain fields associated with this breaking (z) decouple (hidden sector);
3) Certain fields become superheavy (X);
4) Remaining fields (ϕ_i) are to be observed in low energy theory (observable sector). After shifting all fields by their v.e.v.s, and discarding terms involving decoupled (z) or superheavy fields (X), $\mathcal{L}_{SOFT}$ is obtained[137),138),130]

$$\mathcal{L}_{SOFT} = m^2\sum_i|\varphi_i|^2 + \Big(m\sum_n (A-3+n) f_n + h.c.\Big) - \frac{1}{2}\Big[m_{\tilde g}\,\tilde g\tilde g + m_{\tilde W}\,\tilde W\tilde W + m_{\tilde B}\,\tilde B\tilde B + h.c.\Big] \tag{68}$$

while more general forms[139] are certainly possible but irrelevant to our discussion here. A is a model-dependent parameter[138] of O(1) and f_n is the nth term in the superpotential $f(\phi_i)$, written schematically as $f(\phi_i) = \Sigma_n f_n = \Sigma_n C_n \phi_i^n$ with $n = 1,2,3,\ldots$, not necessarily terminating at 3, because non-renormalizable terms are allowed in these effective theories. The mass parameters m and $m_{\tilde V}$ in (68) depend on the form of G and f_{ab}. We may distinguish three interesting cases (see also Fig. 4):

i) "Minimal" (58)[137),138]

[All scalar fields $z;\phi_i$ satisfy (59)]:

$$\left\{\begin{array}{l} m = m_{3/2} \\ \qquad\qquad \rightarrow \tilde m = O(m_{3/2}) \\ m_{\tilde V} = 0 \end{array}\right.$$

ii) "Mini-Maxi"[110),111)]

$$\left[\begin{array}{l}\text{All scalar fields } \phi_i \text{ satisfy (59),}\\ \text{while the z (Polonyi) field}\\ \text{satisfies (61)}\end{array}\right] :$$

$$\left|\begin{array}{l} m = m_{3/2} \neq 0 \ (62) \\ \downarrow \\ \text{undetermined at} \rightarrow \tilde{m} = O(m_{3/2}) \\ \text{tree level} \\ \\ m_{\tilde{V}} = \begin{cases} 0 \\ \text{or} \\ \text{given by (66)} \end{cases} \end{array}\right. \qquad (69)$$

iii) "Maximal"[112)-114)]

$$\left[\text{All scalar fields } z, \phi_i \text{ satisfy (61)-like G's}\right]:$$

$$\left|\begin{array}{l} m = 0 \ ; \ A = 0 \\ m_{\tilde{V}} \neq 0 \ (66) \\ \downarrow \qquad\qquad \rightarrow \quad \tilde{m} = O(m_{\tilde{V}}) \\ \text{undetermined at} \\ \text{tree level} \end{array}\right.$$

It is interesting to notice that in the "Maximal" case iii), the emerging low energy theory[112)] (observable sector) is globally supersymmetric ($m = 0$, $A = 0$) and so all the burden of the necessary global SUSY breaking shifts unavoidably to the gaugino mass $m_{\tilde{V}}$[112)-114)]. If the gaugino mass is non-zero (66), then radiative corrections will generate non-zero scalar masses. Thus, we expect the gaugino mass to be $O(M_W)$, but the gravitino mass could a priori be very different[113),114)]. This possible decoupling[113),114)] of the local SUSY breaking parameter ($m_{3/2}$) and the global SUSY breaking parameter ($\tilde{m}$) has some very interesting particle physics and cosmological implications to be discussed later. In order to cover all cases, $\tilde{m}$ will generically stand for either $m_{3/2}$ or $m_{\tilde{V}}$, both originating from supergravity and providing the "seed" for global SUSY breaking. Clearly, the form of $\mathcal{L}_{SOFT}$ as given by (68) is simple enough. Is that good for physics applications ? Certainly, the FCNC constraint, as expressed by (40), is satisfied automatically. That is the good news. The bad news is the fact that all scalar bosons including Higgs seem to get positive m^2 terms, thus making gauge symmetry breaking [e.g., $SU(2)_L \times U(1)$] problematic. Many attempts[137)] to circumvent this problem at the tree level by using superfluous, extra-singlet fields have either failed or became rather involved[140)], because of the induced de-stabilization[141)] (if GUTs exist) of the gauge hierarchy. So, what is left ? Well, I have to remind you that the form of $\mathcal{L}_{SOFT}$ (68) is strictly valid at energy scales just below M (or M_X). At low energies (M_W) $\mathcal{L}_{SOFT}$ should look a bit more complicated, because of the renormalization effects (like fermion mass or coupling constant renormalization in GUTs).

Let us see how this only remaining (?) way out works by first defining the so-called minimal low energy N = 1 supergravity model. It involves the smallest gauge group, $SU(3)_C \times SU(2)_L \times U(1)$, and the fewest particles possible for such a theory. The chiral superfields of the minimal model are $Q_i(3,2,1/6)$, $\bar{U}_i(\bar{3},1,-2/3)$, $\bar{D}_i(\bar{3},1,1/3)$, $L_i(1,2,-1/2)$, $\bar{E}_i(1,1,1)$, $H_1(1,2,-1/2)$ and $H_2(1,2,1/2)$ in an obvious notation ($i = 1,2,3$ is the generation index) and with the necessary and unavoidable appearance, because of chirality reasons, of the two Higgs doublets H_1 and H_2. Given this set of fields the most general gauge invariant superpotential f which conserves matter parity is

$$f = h_U Q\bar{U}H_2 + h_D Q\bar{D}H_1 + h_E L\bar{E}H_1 + m_4 H_1 H_2$$

where h are 3x3 Yuakwa coupling matrices. In fact, many GUTs lead to this low energy theory, so I shall use it as a prototype for an effective low energy theory. We assume the existence of a theory of everything (TOE) - perhaps an extended supergravity theory - which specifies the parameters of the (minimal ?) low energy effective theory at a renormalization scale $\mu = O(M)$. We then evaluate the non-gravitational standard model radiative corrections which cause the soft SUSY breaking parameters to change when one lowers the renormalization scale μ. As long as the soft SUSY breaking scale $\tilde{m}$ is much less than the renormalization scale, the renormalization of the soft SUSY breaking parameters is strictly linear and multiplicative, with a matrix renormalization group equations of the form[142)]

$$\mu \frac{\partial \tilde{\mathcal{M}}^2}{\partial \mu} = [\text{matrix}]\, \tilde{\mathcal{M}}^2 \qquad (70)$$

The pattern of evolution that we seek for m_H^2, $m_{\tilde{q}}^2$, $m_{\tilde{\ell}}^2$ is illustrated in Fig. 5. We want $m_{\tilde{q}}^2$ and $m_{\tilde{\ell}}^2$ to remain positive at all renormalization scales μ, while m_H^2 can become negative. If $m_H^2 < 0$, then weak $SU(2)_L \times U(1)$ gauge symmetry breaking is possible, while the positivity of $m_{\tilde{q}}^2$ and $m_{\tilde{\ell}}^2$ guarantees that $SU(3)_C \times U(1)_{EM}$ is unbroken. Indeed, that is what is happening[136),142)-144)]!

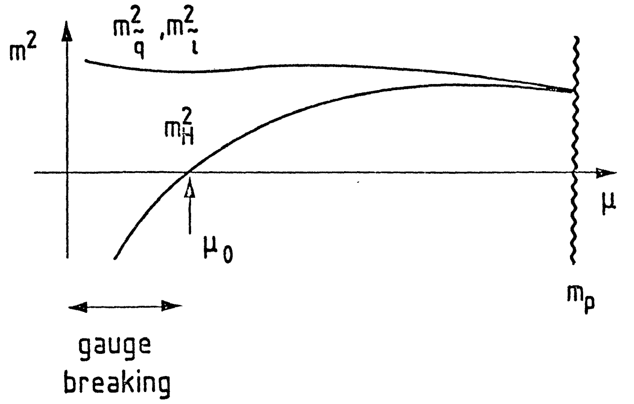


Fig. 5: Renormalization group evolution of scalar mass squared parameters, showing how m_H^2 may be driven below zero at a scale μ_0, triggering weak gauge symmetry breaking, while $m^2_{\tilde q}$ and $m^2_{\tilde l}$ remain positive.

The miraculous spontaneous $SU(2)_L \times U(1)$ breakdown is easily wrought by postulating a heavy quark[136),142)-144)]. Its contribution to the evolution of m_H^2 is positive

$$\mu \frac{\partial m_H^2}{\partial \mu} = (+)\, g^2_{H\bar t t}\, m^2_{\tilde t} - g^2_{2,1}\, m^2_{\tilde W,\tilde B} \qquad (71)$$

It is clear from (71) that if $g^2_{H\bar t t}$ is large enough, the t-quark contribution to the right-hand side will dominate, thus m_H^2 decreases as μ decreases. This is just what we need, since we start with $m_H^2 \geq 0$ at large scale $\mu = O(M)$. Then, as we reduce μ, m_H^2 decreases and may become negative at $\mu_0 = O(M_W)$ if $g_{H\bar t t}$ (and hence m_t) is large enough (see Fig. 5). This scenario[136),142)-144)] can be carried through for the SUSY standard model with

$$20\ \text{GeV} \lesssim m_t \leq O(150)\ \text{GeV} \qquad (72)$$

perfectly consistent with (8).

Furthermore it is easy to understand why $m^2_{\tilde l}$ and $m^2_{\tilde q}$ are positive definite. For sleptons and for the first five squark flavours ($\tilde u, \tilde d, \tilde s, \tilde c, \tilde b$), there is no $g^2_{H\bar t t}$ term in (71), so no danger of decreasing their $(\text{mass})^2$ at low energies; if anything, it would increase. For the $\tilde t$ squark, the $+g^2_{H\bar t t}$ term exists in (71) but there also exists a negative term ($-g_3^2 m^2_{\tilde g}$) which finally wins since $m_{\tilde t} \sim O(m_{\tilde g})$ and clearly $g_3^2 > (g^2_{H\bar t t})_{maxim} \simeq O(g_2^2)$. The net result is that squarks and sleptons have always (at all energy scales) positive m^2. It is highly remarkable that starting from an unbroken gauge symmetry $SU(3)_C \times SU(2)_L \times U(1)$ at superhigh energies M and with the most natural initial conditions (all scalar masses equal), dynamics entail both the energy scale (M_W) and the form $[SU(2)_L \times U(1)_Y]$ of the gauge symmetry breaking.

An interesting possibility in this framework is the determination of the weak interaction scale by dimensional transmutation[143),144)]. The multiplicative renormalization (70) of the soft SUSY breaking mass parameters means that the renormalization scale μ_0 at which m_H^2 goes negative is independent of the magnitude of $\tilde m^2$. The value of μ_0 is determined by the logarithmic rate of evolution specified by the renormalization group equations (70). Hence

$$\mu_0 = M \exp\left(-\frac{O(1)}{\alpha_t}\right) \qquad (73)$$

Once $m_H^2 < 0$, it is possible to have weak gauge symmetry breaking and $M_W = O(\mu_0)$, implying through (73) the highly desirable relation

$$\frac{M_W}{M} \simeq \exp\left(-\frac{O(1)}{\alpha_t}\right) \simeq 10^{-16} \qquad (74)$$

The dynamical determination of M_W has been realized[143),144)]. The first half (50) of the second SUSY golden rule has been satisfied. It should be emphasized that we have tacitly assumed $\tilde m < \mu_0$, otherwise the RGE will be frozen at some renormalization scale $\mu > \mu_0$, implying $m_H^2 > 0$ and thus no $SU(2)_L \times U(1)$ breaking. But who determines $\tilde m$? In the "minimal" case (69i), $\tilde m$ is put in by hand and that is no good. In the "Mini-Maxi" or "Maximal" cases (69), $\tilde m$ is undetermined at the tree level and we should use non-gravitational standard model radiative corrections to determine it, thus finally realizing the no-scale model dream[110)-114)].

To explore the basic mechanism[110)] for this trick, we consider the usual low energy Higgs potential of the SUSY standard model, in an idealized limit where the mixing between the two light Higgs doublets is neglected

$$V = \left(\frac{g_2^2+g'^2}{8}\right)\left(|H_1|^2 - |H_2|^2\right)^2 + m_1^2|H_1|^2 + m_2^2|H_2|^2 \qquad (75)$$

We denote by H_2 the Higgs field coupled to the t quark. Radiative corrections (71) drive $m_2^2 < m_1^2$ and when $m_2^2 < 0$ weak $SU(2)_L \times U(1)_Y$ gauge symmetry is spontaneously broken and the vacuum energy can become negative:

$$V_{min} = \frac{-2\,|m_2^2|^2}{(g_2^2+g'^2)} \propto (-)\,\tilde m^4 \qquad (76)$$

In writing Eq. (76), we have recalled that since the Higgs mass is multiplicatively renormalized (70) for $\mu > O(\tilde{m})$, the Higgs mass is always $\propto \tilde{m}$, and hence $V_{min} \propto |m_2^2|^2 \propto \tilde{m}^4$. The negative coefficient in Eq. (76) means that increasing $|m_2^2|$ and hence $\tilde{m}^2$ is energetically preferred, at least for small values of $\tilde{m}$. However, if $\tilde{m}$ gets to be larger than the scale μ_0 at which m_2^2 falls through zero, then the evolution of m_2 with the renormalization scale μ will become truncated at $\mu = O(\tilde{m}_2) > \mu_0$, m_2^2 will never become negative, and the potential will always be positive semidefinite. The general form of $V_{min}(\tilde{m})$ is therefore as shown in Fig. 6.

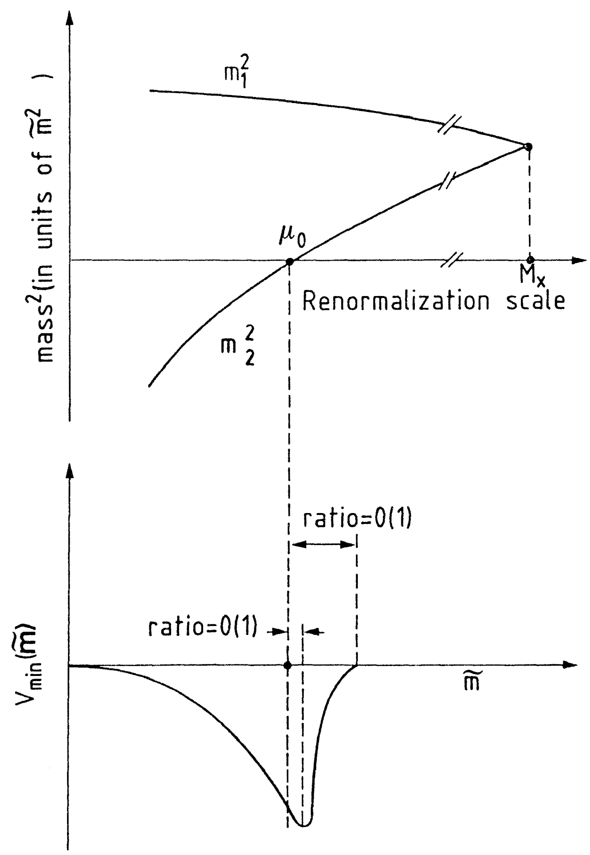


Fig. 6: Sketch of the variation of the SUSY-breaking mass parameters m_1^2, m_2^2 with the renormalization scale μ. The Higgs (mass)^2 $m_2^2 = 0$ at a scale μ_0, which determines the dynamically preferred value of $\tilde{m}$, as seen in the bottom half of the figure.

It decreases from zero to negative values as $\tilde{m}$ increases from zero, but then rises to zero again for some $\tilde{m} = O(\mu_0)$. (The precise value depends on the choice of renormalization scheme, but physical parameters such as particle masses do not depend on this choice.) It is apparent from Fig. 6 that there must[110),111)] be a dynamically preferred value of $\tilde{m}$ in the range $(0,\mu_0)$. Indeed, one finds[110),111)]

$$V_{min}(\tilde{m}) = -c^2 \tilde{m}^4 \ln^2 \frac{\tilde{m}^2}{d(\mu_0^2)} \tag{77}$$

with c and d calculable parameters. Minimizing the potential (77) with respect to $\tilde{m}$, one finds a minimum at

$$\tilde{m} \simeq O(\mu_0) \tag{78}$$

Remember that μ_0 is a dimensional transmutation scale (73), so that the preferred value of $\tilde{m}$ is also fixed by dimensional transmutation, leading through Eqs. (74) and (78) to the golden relation (50),

$$\frac{M_W}{M} \simeq \frac{\tilde{m}}{M} \simeq \exp\left(-\frac{O(1)}{\alpha_t}\right) \simeq 10^{-16} \tag{79}$$

This simultaneous and dynamical determination[110)-114)] of M_W and $\tilde{m}$ combined with the non-renormalization theorems[81)] of global SUSY (stabilization) and with the classic "missing partner" mechanism[145),146)] (natural tree-level GUT Higgs-electroweak Higgs splitting), completely solves the gauge (scale) hierarchy problem.

No-scale SUSY standard models[110),111)] contain three adjustable parameters: possible non-zero gaugino masses $\xi \equiv m_{\tilde{V}}/m_{3/2}$, non-zero H_1H_2 mixing characterized by a mixing parameter m_4 ($\hat{m}_4 \equiv m_4/m_{3/2}$) and the A parameter[138)] (68). As seen in Fig. 7, we find[110),111)] domains of ξ and $\hat{m}_4$ which give phenomenologically acceptable models for which all charged sparticles have masses above 20 GeV (denoted by P), and the cosmological density of the lightest stable neutral sparticle is less than 2×10^{-29} gm/cc (denoted by C). Varying the parameters, we find[110),111)] acceptable solutions for values of m_t in the range indicated (72) and thus consistent with (8).

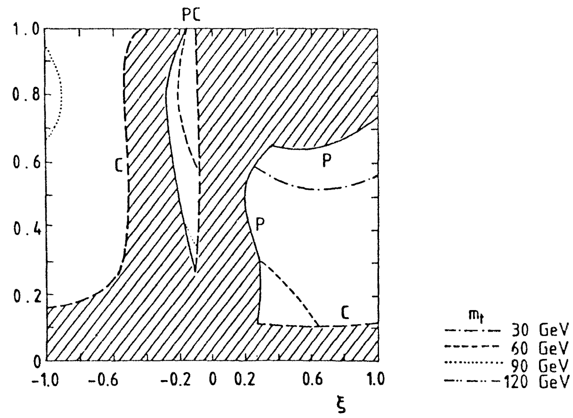


Fig. 7: Values of m_t (in GeV) for general values of the Higgs mixing parameter $\hat{m}_4$, $\xi = m_{\tilde{V}}/m_{3/2}$ and A = 3. Our present vacuum is unstable in the allowed region on the right of the figure. The shaded domain indicates values of $\hat{m}_4$ and ξ disallowed because of the absence of a charged sparticle with mass less than 20 GeV (P), and/or because of an excessive cosmological density of the lightest neutral sparticle (C).

No-scale SUSY GUT models have been also constructed[112]. In most cases, it is obligatory[112] to use the "maximal" alternative (69iii) because global SUSY ($m = 0$, $A = 0$), at the GUT scale M_X is badly needed in order to get rid of the highly undesirable huge GUT vacuum energy. That is good news, since the decoupling entailed in principle between $m_{3/2}$ and $\tilde{m}$ is very welcome. There are two very interesting physical implications of this gravitino mass liberation movement[112)-114], one cosmological and the other hierarchical. As is well known[147], standard cosmological constraints (critical energy density, primordial nucleosynthesis, etc.) exclude gravitino masses in the region 1 keV to 10^4 GeV, thus excluding $m_{3/2} \simeq O(M_W)$. Inflation can save[148] the gravitino by diluting its original number density, but gravitino regeneration processes[149] put a rather low upper bound[149)-151] on the reheating temperature $T_R \leqslant O(10^{10}$ GeV). This is rather unfortunate because then, "standard" baryosynthesis[152] is in conflict with long-lived protons[153),154]. This is the gravitino problem. It seems that we need either superlight ($<$ 1 keV) or (super)heavy ($>$ 10^4 GeV) gravitinos. No-scale models encompassing both possibilities have been constructed[113)-114], thus resolving the gravitino problem in a Gordian-Knot way by "cutting" the $m_{3/2} \sim O(M_W)$ relation. Actually, the superlight gravitino case[114] has a very interesting hierarchical implication[114]: it completely solves the strong CP problem (38). One simply notices from (51) that non-minimal gauge boson kinetic terms, necessarily needed for non-vanishing gaugino masses (66), give rise to $\theta_{QCD} \sim \langle \mathrm{Im}\, f_{ab}/\mathrm{Re}\, f_{ab}\rangle$. Then, since there are two dynamical degrees of freedom (Rez, Imz), it is evident that these will determine dynamically two physical parameters. Since the gravitino is decoupled, we may naturally choose $m_{\tilde{v}}$ and θ_{QCD} as the two dynamically determined quantities. While standard model radiative corrections fix the gaugino mass $m_{\tilde{v}}$, the only non-trivial dynamical dependence on θ_{QCD} comes from non-perturbative QCD effects which favour $\theta_{QCD} = 0$. The dynamical freedom accorded to us by the Im f_{ab} field, enable us to "relax" to $\theta_{QCD} = 0$, thus completely solving[114] the strong CP problem. In this case, in a class of no-scale models[114], the gravitino mass is determined to be

$$m_{3/2} = O\left(\frac{m_{\tilde{v}}^{p}}{M^{p-1}}\right) << m_{\tilde{v}} \tag{80}$$

with $1 < p < 2$, in order to avoid too singular couplings to ordinary matter[155].

I hope I have convinced you by now that (no-scale) phenomenological supergravity models are worth considering and I proceed next to discuss experimental evidence for (or against) them.

4. - (S)-EXPERIMENTAL EVIDENCE (?) FOR SUPERSYMMETRY

Low energy supersymmetric models have a very rich structure that makes them experimentally vulnerable at accessible (present or very near future) energies, if SUSY indeed solves the gauge hierarchy problem [see Eq. (49)]. Since several reports exist[156], covering, in rather lengthy detail, the numerous experimental consequences of SUSY, I shall limit my discussion here to the relation between SUSY and the new experimental results mentioned in the beginning.

Let me start with direct evidence, i.e., with SUSY particle production. As I emphasized before [see discussion above Eq. (39)], SUSY particles in general can only be produced in pairs. This means that among the decay products of every sparticle, there must be another sparticle, and hence the lightest sparticle must be stable. The lightest sparticle is probably[90] neutral and not strongly interacting and the most likely candidate[90] may be the photino $\tilde{\gamma}$ (39). Thus, a characterictic signature for SUSY could be[156] missing energy-momentum carried away by weakly interacting photinos, for example, from gluino or squark-pair production[157]:

$$p + \bar{p} \longrightarrow \tilde{g}\tilde{g} + X \;\; (\tilde{g} \to q\bar{q}\tilde{\gamma}) \quad ; \quad \tilde{q}\tilde{q} + X \;\; (\tilde{q} \to q\tilde{\gamma}) \tag{81}$$

or from supersymmetric W decay[158]:

$$p + \bar{p} \longrightarrow W + X, \quad W \longrightarrow \tilde{W} + \tilde{\gamma}, \quad \tilde{W} \longrightarrow q\bar{q}\tilde{\gamma}\,;\, \ell\nu\tilde{\gamma} \tag{82}$$

It is highly exciting that the observed UA(1,2) "Zen" or "Wen" events, with their characteristic large missing energy-momentum, may be naturally due to production and decay of SUSY particles. Indeed, detailed studies[159)-165] have shown that the UA(1) "Zen" (monojet) events could be due to either squark or gluino production (and decay), according to (81), if either[159] $m_{\tilde{q}}$ or $m_{\tilde{g}} = O(40$ GeV); lower masses down to $O(30$ GeV) cannot be yet firmly (?) excluded[160),163),165]. As is apparent from Fig. 8, for

$m_{\tilde g} < O(40\ \mathrm{GeV})$ and when all experimental cuts are taken into account, the total cross-section for $\tilde g\tilde g$ production is dominated by the one-jet final state[159),160)]. Similar results hold true for the squark case[159),163)]. So, despite the naîve expectations from (81) that $\tilde g\tilde g$ or $\tilde q\tilde q$ production should give four-jet or two-jet final states respectively, surprisingly enough for $m_{\tilde g}$ or $m_{\tilde q} \lesssim O(40\ \mathrm{GeV})$ one-jet final states dominate, in accordance with experiment[13),10)].

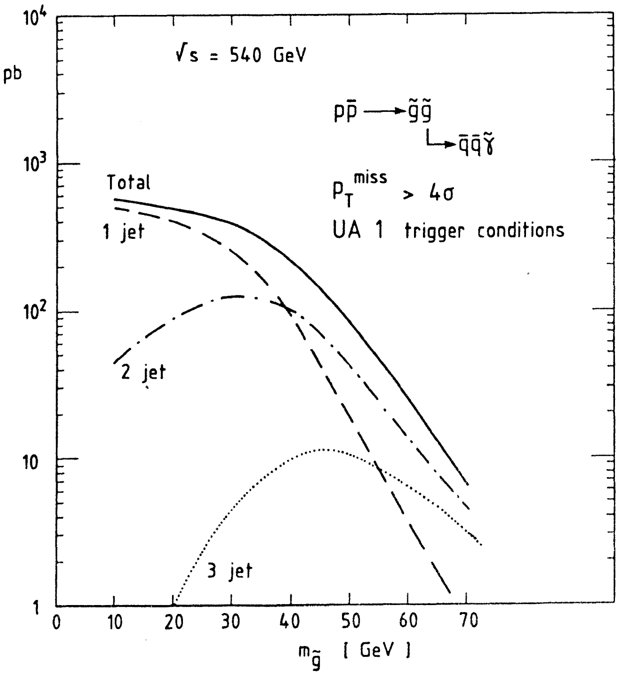


Fig. 8: The total and topological cross-sections for $\tilde g\tilde g$ production followed by $\tilde q \to q\bar q\tilde\gamma$ decay giving one-, two- and three-jet final states with $p_T^{miss} > 4$, fulfilling the UA1 trigger requirements [from Ref. 159)].

Are then the observed events with large p_T^{miss} due to the production of either $\tilde g$ or $\tilde q$ with mass $O(40\ \mathrm{GeV})$? Both are possible interpretations of the UA(1) data[13),10)], but neither can be confirmed or refuted until more data are accumulated. Nevertheless, the squark interpretation has been favoured[159),163)] on two grounds (i) the hardness of the observed missing $p_\perp$ spectrum, which is more naturally explained by two-body $\tilde q \to q + \tilde\gamma$ decays and (ii) the thinness of the observed monojets, which disfavours $\tilde g \to q + \bar q + \tilde\gamma$ decay which yields monojets with invariant masses up to $O(20\ \mathrm{GeV})$ and an average of $O(10\ \mathrm{GeV})$, while the squark decay yields monojet invariant masses $O(2\ \mathrm{GeV})$ consistent[159)] with the observed ones. Other more contrived explanations are still possible. For example, if $\tilde g$ and $\tilde\gamma$ are both very light and approximately degenerate ($\sim$ 3 GeV) to ensure long enough $\tilde g$ lifetime, then $qg \to \tilde q\tilde g \to q(\tilde\gamma\tilde g)$ with $m_{\tilde q} \sim 100$ GeV may also explain[162)] the "monojet" events. Concerning the "photon" UA(1) event(s), it may possibly be a monojet with a large collimated electromagnetic component containing one or more π^0's or η's whose charged multiplicity fluctuated down to zero. This event actually contains some soft charged tracks which are nearby in angle and could perhaps be associated in the "monojet". Another source of photons could be[164)] $\tilde q q \to \tilde\gamma\gamma$, assuming $m_{\tilde q} \leq E_T^{trigger}$ such that the SUSY content of the proton is plausibly excited, but in this case[164)], one should take care of the jet coming from the decay of the left-over spectator $(\bar{\tilde q})$[166)]. Putting everything together, it does not look unreasonable to me to pursue the interpretation[159),163)] of the UA(1) monojet data as squarks with $m_{\tilde q} \simeq 40$ GeV, and infer a lower limit $m_{\tilde g} \geq 40$ GeV on the gluino mass. Furthermore, such an assumption is not in contradiction with possible explanations of the UA(2) "Wen" events[14),10)]. For example[161)], $p\bar p \to \tilde W\tilde g$ or $\tilde W\tilde q$ production, followed by $\tilde W \to e\nu\tilde\gamma$ and $\tilde g \to qq\tilde\gamma$ or $\tilde q \to q\tilde\gamma$, would produce "Wen" events (e + jet + large amounts of p_T^{miss}.). It has been argued[161)] that with $m_{\tilde q} \simeq$ (40-60) GeV, $m_{\tilde g} \simeq$ (70-100) GeV and $m_{\tilde W} \simeq$ (35-40) GeV, one may probably get suitable (?) rates. In addition, such an explanation[161)] is consistent with the observation that in the three UA(2) "Wen" events, the missing "ν" vector is consistently larger than the p_T of the observed electron. They should on average be equal in $W \to e\nu$ decay, but could easily be different in SUSY, where the "ν" is actually a combination of one ν and two photinos $\tilde\gamma$.

If we indeed buy the squark explanation [$m_{\tilde q}$ = 40 GeV; $m_{\tilde g} \geq O(40\ \mathrm{GeV})$] of the Zen-Wen events, then its phenomenological implications are rather dramatic. In a large class of phenomenological supergravity models[142)-144),110)-114)] discussed before, one may write down[144)] convenient approximate formulae for physical squark and slepton masses at relevant renormalization scales $\mu = O(M_W)$:

$$m^2_{\tilde q} \simeq m^2(1+7.6\,\xi^2)\,, \quad \tilde q = \tilde u, \tilde d, \tilde s, \tilde c, \tilde b \tag{83}$$

and

$$\begin{aligned} m^2_{\tilde\ell_L} &\simeq m^2(1+0.15\,\xi^2) \\ m^2_{\tilde\ell_R} &\simeq m^2(1+0.5\,\xi^2) \end{aligned} \tag{84}$$

where m and ξm ($\equiv m_{\tilde v}(M)$) are respectively the scalar boson and gaugino masses at $\mu = M_X$ or M (m is commonly, but not always necessarily, indentified with the gravitino mass). The corresponding formulae for gaugino masses take the simple form[136),142)]

$$m_{\tilde g} = \frac{\alpha_3}{\alpha_G}\, m_{\tilde v}(M) \quad ; \quad m_{\tilde\gamma} = \frac{\alpha_1}{\alpha_G}\, m_{\tilde v}(M) \tag{85}$$

implying

$$\frac{m_{\tilde{\gamma}}}{m_{\tilde{g}}} \simeq \frac{3}{8} \cdot \frac{\alpha}{\alpha_3} \cdot \frac{1}{\sin^2\theta_W} \qquad (86)$$

where it has been assumed that, thanks to grand unification, all gaugino masses are equal at M_X or M, not necessarily an unavoidable assumption[112]. Clearly enough, as it has been repeatedly emphasized in the literature[143),144),110)-114)] for some time now, the supergravity sparticle mass spectrum is rather tight. For example, assuming $m_{\tilde{q}} = 40$ GeV, $m_{\tilde{g}} > O(40 \text{ GeV})$ and $m_{\tilde{\chi}_{L,R}} > 20$ GeV, it is trivial to show[167] that the set of equations (83)-(86) imply:

$$20 \text{ GeV} \leq m_{\tilde{\ell}_{L,R}} \leq 30 \text{ GeV}$$

$$5 \text{ GeV} \lesssim m_{\tilde{\gamma}} \lesssim 10 \text{ GeV} \qquad (87)$$

and

$$40 \text{ GeV} \lesssim m_{\tilde{g}} \lesssim 60 \text{ GeV}$$

a rather "light" and easily accessible experimentally spectrum. In addition, while SUSY in general entails the existence of at least two Higgs-doublets, dimensional transmutation[144] or no-scale type[110)-111)] supergravity models ask for the existence of a "light" [sometimes[144] < 10 GeV] neutral Higgs boson. Actually, if indeed (87) holds, then the O(10 GeV) upper bound is certainly quite firm[144),167)]. The suggestion was then made[144], some time ago, that $\Upsilon \to H^0 + \gamma$ is an excellent place[168),169)] to look for such a "non-standard" light Higgs. As we heard at this conference[20], and I discussed in the beginning, there is already a candidate: the Crystal Ball ζ(8.3 GeV). Though it seems difficult, when the dust settles down, the ζ(8.3) GeV to be identified with such a highly desirable object as the SUSY light Higgs boson. At the moment of writing this looks the least unlikely possibility, while more remote possibilities include glueballons[170], glueballs[171], conventional two Higgs doublets[172]... If indeed (87) holds true, then news (good or bad) should come very soon from almost everywhere: $e^+e^- \to \tilde{\ell}^+\tilde{\ell}^-$; $e^+e^- \to (\tilde{\gamma}\tilde{\gamma})\gamma$[173], [one event has already been reported[88]], $p\bar{p} \to (\tilde{g}\tilde{g} \text{ or } \tilde{q}\tilde{g}) + X$, $p\bar{p} \to (W^+ \to \tilde{\ell}+\tilde{\nu}) + X$ or $p\bar{p} \to (Z^0 \to \tilde{\ell}\tilde{\ell}, \tilde{\nu}\tilde{\nu}) + X$ and $\Upsilon \to H^0 + \gamma$. Wait and see!

In the meanwhile, it is worth discussing the <u>indirect evidence(?)</u> for SUSY, i.e., the possible "virtual" contributions of SUSY particles to low energy processes? The main contribution has to do with FCNC processes. SUSY CP violating effects are of great interest[83),174)-178)]. SUSY models, with their rich structure, contain extra CP-violating phases beyond the K-M phase δ_{K-M}. Assuming that $\mathcal{L}_{SOFT}$ is given by (68), one may recognize a new physical CP-violating phase which may be conveniently defined as the phase difference ϕ between $m_{\tilde{v}}$ and A in (68). Alas, if δ_{K-M} is switched off, the phase ϕ cannot take all the responsibility for the observed CP-violation effects[174),177)]. The reason is simple: there is an extra contribution[83),175)-178)] to d_n (DEMON)

$$(d_n)_{\tilde{g}} = -\frac{2e}{9\pi}\frac{\alpha_3}{m_{\tilde{g}}} f\left(\frac{m^2_{\tilde{q}}}{m^2_{\tilde{g}}}\right) \sin 2\varphi \cdot \text{"}\Theta\text{"} \qquad (88)$$

["Θ" is some product of K-M angles (θ_i)], which puts an upper bound on $\phi < O(10^{-2})$, if (3) has to be obeyed. Then, the ϕ contribution to $(\varepsilon)_{CP}$ becomes negligible. Still, we should keep in mind that a non-zero ϕ may give a big contribution to d_n (88). Let us assume for now $\phi = 0$ and proceed as usual with $\delta_{K-M} \neq 0$.

As I mentioned before, the form of SUSY soft breaking (68) induced by supergravity is such that the FCNC constraint as expressed by (40) is automatically satisfied. However, at low energies because of the renormalization effects (70), the down squark mass matrix becomes[179),175)-176)]

$$M^2_{\tilde{d}} = \begin{pmatrix} \tilde{m}^2_{(L)}\mathbb{1} + M_d M_d^+ + C M_u M_u^+ & A^*\tilde{m} M_d \\ A\tilde{m} M_d^+ & \tilde{m}^2_{(R)}\mathbb{1} + M_d M_d^+ \end{pmatrix} \qquad (89)$$

where $M_{d(u)}$ denotes the down (up) quark mass matrix and C is a model-dependent parameter[179]. Notice that while at energy scales $\mu = O(M)$, $C(M) = 0$, at energy scales $\mu \simeq O(M_W), C(M_W) \simeq O(1)$[179),176)]. From (89), it is apparent that we can no longer diagonalize M_d and $M^2_{\tilde{d}}$ by the same transformations (as in the case of C = 0), and in this way we can expect the manifestation of the K-M CP-violating phase δ_{K-M} in the flavour-violating $\tilde{d}_i - d_j - \tilde{g}$ gaugino vertices. This fact is not without consequences for the kaon[175)-177)] and beon systems[176]. In the case of the kaon system, this extra SUSY CP-violating contribution [Fig. 3a, with $\tilde{g}$ exchange] is maybe needed. As I mentioned in the beginning [see (7)-(9) and subsequent discussion], long bottom lifetime (7), "light" top quark mass (8), minute $b \to u$ decay (9) and $\varepsilon_{CP} \simeq 10^{-3}$ do not fit in the standard model, for acceptable (?) values[180] of the famous B-parameter (< 0.5), which parametrizes the $\langle \bar{K}_0 | \cdots | K_0 \rangle$ matrix element. In the standard model $\varepsilon_{CP} \sim$ "Θ" (m_t^2/M_W^2)[181]

where "Θ" is some product of K-M angles (θ_i), and clearly small "Θ" (long τ_b) implies large m_t, and thus for a given τ_b, one has always a lower bound[182] on m_t, as shown in Fig. 9.

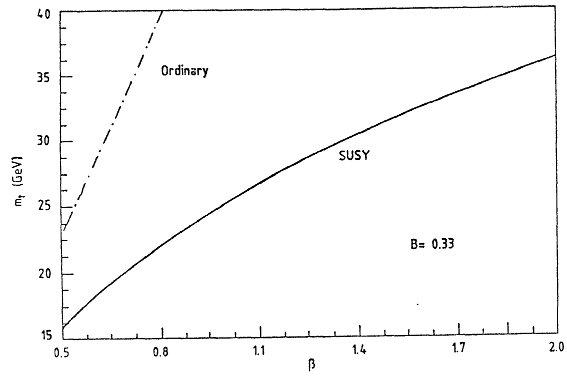


Fig. 9: The minimal allowed top quark mass, m_t, in the ordinary and SUSY case, as a function of β ($\tau_b = \beta \cdot 10^{-12}$ sec). The SUSY inputs can be found in the text [from Ref. 176)].

As can be seen from this figure, there is a problem. However, in the supersymmetric case, the extra gluino contribution (Fig. 3a)

$$(\varepsilon_{CP})_{\tilde g} \sim \text{"}\Theta\text{"}\,\frac{m_t^4}{m^4_{\tilde g,\tilde q}}$$

solves[176] the problem. The appearance of the fourth power of m_t ($\tilde q$ exchange) instead of the second power of m_t (q-exchange) does the trick[176], as is apparently shown in Fig. 9. For the plot of Fig. 9 we have taken[176] $\bar R$ = 0.03 (9), C = -1, $m_{\tilde q,\tilde g} \simeq$ O(40 GeV), $|A|$ = 3 preferred value theoretically[110),111)] and phenomenologically[167], and B $\simeq$ 1/3, the presently favourable value, obtained[180] through PCAC...

It is remarkable that with this rather well-motivated input, the SUSY contribution on $(\varepsilon)_{CP}$ makes such a dramatic difference[176] with respect to the ordinary standard model, as the one indicated in Fig. 9.

There are also extra SUSY contributions[176),177),183)] to ε'/ε. The so-called pengluino diagrams[183] make in certain cases drastic contributions, as indicated in Fig. 10. The very interesting possibility arises[176),183)] of negative or very close to zero values for ε'/ε in accord with (10) and in short contrast with the ordinary standard model result[18),19)] (Fig. 10 solid lines).

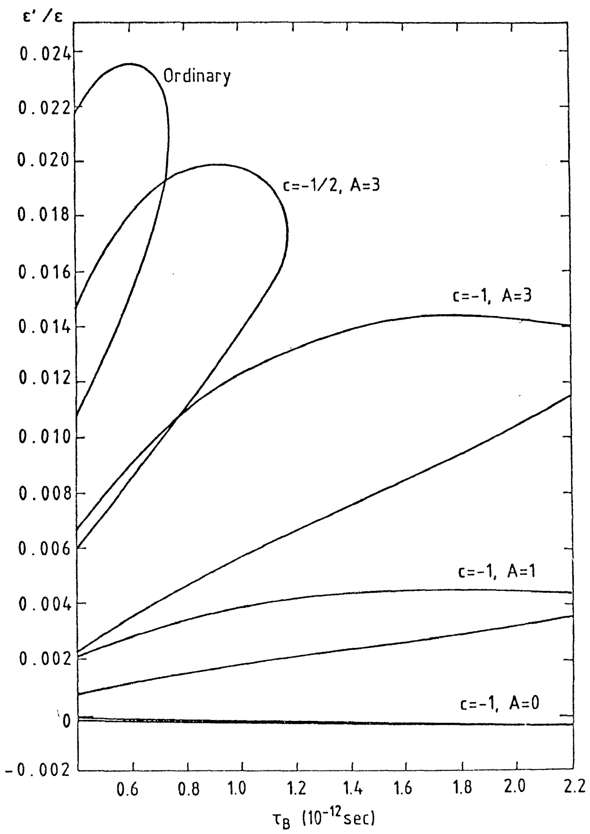
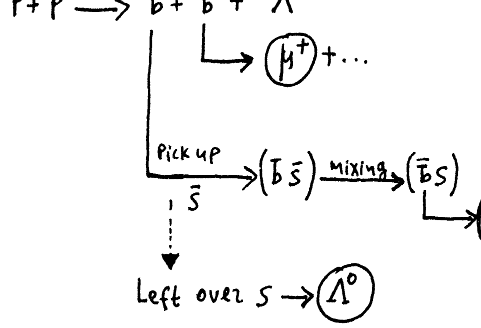


Fig. 10: ε'/ε as a function of the bottom lifetime τ_b for a top quark mass equal to 40 GeV in the SUSY model (for different choices of the parameters C and A) and the standard model [from Refs. 176) and 183)].

Very interesting effects occur also in the beon system[176]. The $B^0_s - \bar B^0_s$ mixing can rise from 0.1-0.2 (ordinary model)[181),184)] to 0.9 - 0.93 while the $B^0_d - \bar B^0_d$ mixing may rise from 10^{-5} (ordinary model)[181),184)] to 10^{-3}, thanks again to the gluino contributions[176](Fig. 3a, but with b instead of d or s). While these differences between ordinary and SUSY models may be testable in the near future, the large same-sign dimuon rate observed[15] from UA(1) may already indicate something[185),186)]. Indeed, it has been rather convincingly argued[186] that the rate, topology and kinematics of these events may be due[185),186)] to a rather large $B^0_s - \bar B^0_s$ mixing, probably unattainable in the ordinary standard model, thus creating another problem for it. For example, the following chain process may take place[186]



(90)

and produce the observed $\mu^+\mu^+\Lambda^0$ or similarly $\mu^-\mu^-\bar{\Lambda}^0$ events. It may be that the SUSY-increased[176] $B^0_s - \bar{B}^0_s$ mixing is needed for a satisfactory explanation[185),186)] of the UA(1) dimuon events. For more details of the SUSY effects in the beon system, see Ref. 176).

SUSY contributions to other low energy phenomena, like $(g-2)_\mu$[96)], ρ[97)], Γ_Z[187)], amazingly enough, satisfy all available data (2), (5), ..., at the present level of accuracy. However, a slight reduction of the errors in these processes may be very revealing. For example, detailed calculations have shown[188)] that in certain cases, even a factor of 3 to 5 reduction of the $(g-2)_\mu$ error may tell us a lot about SUSY models, or a more accurate measurement of Γ_Z can be fatal for certain SUSY models predicting[187)] a Γ_Z increase up to 50% with respect to its ordinary standard model value. Furthermore, it is worth pushing further the limits on d_n (DEMON) in order to find out if we need some $\phi \neq 0$ [see Eq. (88)] or d_n gets its ordinary (or SUSY) standard model value[189)] of $\sim 0(10^{-30}$ e.cm) due only to $\delta_{K-M} \neq 0$.

All in all, it is rather remarkable and highly non-trivial that the simplest (minimal) supergravity models described previously do not only make theoretical sense, by completely solving several hierarchy problems [gauge (scale), strong CP] and by providing for the first time a no-scale type framework[110)-117)], but they also make a lot of phenomenological sense. They do give a rather satisfactory explanation of the new experimental findings in a unified way, without upsetting old results, and they do make well-defined predictions for the near future. Let us see what the future of these SUSY models could be.

5. - SUPER-FUTURE

I have repeatedly emphasized that the N = 1 supergravity framework should be considered only as an "effective theory"[136)] able to describe physics at energy scales below the super-Planck scale M. What is then the fundamental theory ? That is a really tough question and I wish I could give you a straight answer, but we are a bit far (?) from realizing this wish. The apparent success of phenomenological supergravity models may show us the way to get a fundamental theory. Certainly, as we saw before, we need extended supergravities (N > 4), if we are after some natural understanding of the structure of no-scale models[110)-117)]. Global, non-compact symmetries, or non-trivial f_{ab} (66) are contained in the extended supergravity structure (65). The N = 8 extended supergravity[135)] has also been considered as a leading candidate[190)] for providing dynamically[135)] a viable SU(5) GUT theory. It has been speculated[190)] that all fields (spin: 0, 1/2, 1) needed for a realistic SU(5) theory may emerge as composite particles ($\Lambda_{compos.} \simeq M$) made of preons that belong to N = 8 extended supergravity multiplets. It is highly surprising that the maximal allowed value of N (N = 8) is just enough to produce the minimal possible GUT [SU(5)]. This uniqueness seems to me very seductive. It seems that in such a programme, the "magic" properties of SUSY (see Table 2) will play a rather fundamental role. Notice that in this programme, we advocate compositeness and higher N (> 2) supergravities at the Planck scale M and not below (like M_W...) in order to avoid a series of well-known drawbacks. First, as I discussed in the first part, compositeness at a low-scale $\Lambda_c \sim M_W$ is problematic and far from producing an even semi-realistic model. Second, as I mentioned previously, all higher N (> 2) SUSY theories provide vector-like theories[118)], in sharp contrast with the standard model. This is the chirality or "reality" problem. For example, N = 2 SUSY models, where both supersymmetries are broken at M_W, seem to be beset by insurmountable problems[191),192)], e.g., one cannot construct an $SU(3)_C \times SU(2)_L \times U(1)$ model! In our case, by breaking all (local) supersymmetries but one[193)] at M and by advocating[190)] compositeness at M it is hoped that "reality" will be broken.

After these clarifying remarks, let me sketch a possible scenario which, while it encompasses all present facts, also contains the "seeds" of a fundamental theory

$$\begin{array}{ccccc} & & \multicolumn{3}{c}{\text{(Dynamically ?)}} \\ N=8 \ \text{Exten. SUGAR} & \otimes & E_{7,7} & \otimes & SU(8) \\ & & \downarrow & & \downarrow \\ & & \left[\begin{array}{c}\text{non-compact}\\ \text{global}\end{array}\right] & & [\text{local}] \\ \text{broken at } M & \downarrow & & & \\ N=1 \ \text{SUGAR} & \otimes & SU(1,1) & \otimes & SU(5) \\ [\text{Hierarchy problem(s)}]^{193)} & & [\text{no-Scale}]^{110)-114)} & & [\text{GUT}]^{190)} \end{array} \qquad (91)$$

As (91) shows, N = 8 extended supergravity seems to contain all the necessary ingredients, not only to serve as a fundamental theory but also to provide at scales $\mu \ll M$ an effective theory with all the desiderata. Namely, we can get out an effective N = 1 supergravity theory, which while being minimal in the sense that

a) it contains the fewest possible fields, the smallest gauge group $[SU(3)_C \times SU(2)_L \times U(1)]$, the simplest but still most general allowed gauge invariant superpotential (f), and

b) there is a highly constrained set of global SUSY soft breaking terms (68)
it provides

i) a satisfactory solution to the gauge (scale) hierarchy (79) and strong CP hierarchy problems;

ii) a non-trivial extension of the standard model, which not only shares all the successes, but also resolves a few potential problems of the standard model by:

1) providing new particles ($\tilde{g}$, $\tilde{q}$, ..., "light" Higgs) that may be responsible for the recent experimental discoveries;

2) making substantial contributions to CP-violating processes (ε, ε'/ε, B^0-$\bar{B}^0$ mixing,...) thus allowing for long bottom lifetime, "light" top quark, same sign dimuons, ...

3) making nucleons decay predominantly to kaons (K) and not to pions (π)

4) providing freely stable particles [like[90] photinos $\tilde{\gamma}$] with appropriate properties that may serve[92] as the dark matter in the Universe, while their annihilation in our galactic halo may explain[93] some previously mysterious cosmic antiproton ($\bar{p}$) flux[94].

I bet you that it will not take long (by the Berkeley Conference?) to hear that the standard model is dead, long live the SUSY standard model, since we have all reasons to believe that

SUPERSYMMETRY IS SUCH AN OFFER THAT NATURE CANNOT AFFORD TO REFUSE

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