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PROTON DECAY

K. Enqvist and D.V. Nanopoulos

CERN - Geneva

ABSTRACT

In this review, the present status of proton decay is discussed in some detail. The theoretical reasons for expecting proton decay are analyzed and different grand unified models with and without supersymmetry are confronted with experiment. While the minimal SU(5) model already seems to have severe problems, supersymmetric grand unified models are in good shape.

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TABLE OF CONTENTS

1. Introduction
2. Grand Unified Theories
 - 2.1 The Standard Model
 - 2.2 Grand Unification
 - 2.3 Supersymmetric GUTs
3. Baryon Number Violation
 - 3.1 Baryon Number Violation in GUTs
 - 3.2 Monopole Catalyzed Nucleon Decay
 - 3.3 Baryon Number Violation in the Early Universe
 - 3.4 Baryon Number Violation without GUTs
4. Proton Decay and Grand Unified Theories
 - 4.1 Proton Decay in Ordinary GUTs
 - 4.2 Proton Decay in SUSY GUTs
5. Proton Decay Experiments
6. Conclusions

1. - INTRODUCTION

Matter around us is remarkably stable against decay. This is evident from the great age of the earth and the moon that has been deduced, e.g., from radioactive dating of rocks, and similar but more indirect indications of the longevity of the whole Universe can be obtained from astronomical observations of stars, clusters of stars, galaxies and clusters of galaxies. The building blocks of atoms, nucleons and electrons, are therefore very long-lived and do not decay, e.g., into photons. The conservation law that rescues the electron from disintegration is the conservation of the electric charge, which we know to be a local gauge symmetry that is exact to a very high precision: the upper limit on the mass of the photon is measured to be 3×10^{-36} GeV (Chibisov 1976). Indeed, we have a good theoretical reason to believe that the photon is exactly massless (i.e., that the electromagnetic gauge invariance is unbroken) and that the electron as the lightest charged particle is therefore absolutely stable.

The fact that protons and neutrons are also very stable is much harder to explain. Apart from invoking a global symmetry leading to baryon number conservation (Stueckelberg 1938; Wigner 1949, 1952), there is no compelling theoretical reason for believing in nucleon stability. Yet there are very severe experimental lower limits on the nucleon lifetime. Surprisingly enough, the very fact that we do not die of internal radioactive radiation already implies a nucleon lifetime $\tau_n > 10^{16}$ years, which is far longer than the present age of the Universe, $\tau \approx (1 \text{ to } 2) \times 10^{10}$ years. Thirty years ago it was possible to set the stringent limit $\tau_n > 10^{21}$ years from the non-observation of spontaneous Th^{232} fission induced by nucleon decay (Reines, Cowan and Goldhaber 1954). Today improved detection techniques and massive efforts to detect nucleon decay have pushed the limit to an incredibly high level of $\tau_n > 10^{32}$ years (for a recent review of nucleon decay experiments, see Koshiya 1984).

Confronted with such stringent limits, one may wonder what, if any, are the theoretical arguments favouring nucleon instability. One reason, albeit not a very good one, for believing in a finite nucleon lifetime is that if the baryon number conservation were associated with a local gauge symmetry, as is the conservation of the electric charge, we would expect to see a massless long-range field coupled to the baryon number. However, the coupling of such fields has long been known to be very weak (Roll, Krothov and Dicke 1964; Braginsky and Panov 1972). It can be measured in the Eötvös experiment which compares the inertial and gravitational masses of bodies which, according to the principle of

equivalence of the general relativity should be equal. The experiments yield a limit on the fractional difference in baryon number $\Delta B/B \lesssim 0(10^{-4})$, which of course by itself does not say anything about nucleon stability, but tells us that the coupling of the long-range field associated with conserved baryon number should be small. However, because of the weakness of the coupling it is tempting to conjecture that no such field exists and accordingly that there is no gauge symmetry protecting the nucleons from decay.

The first theoretical indication for unstable nucleons was pointed out by Sakharov (1967) in the context of the Big Bang cosmology, which explains successfully the observed $T = 2.7$ K microwave background radiation and predicts correctly the abundance of helium and other light elements in the Universe (for a recent review of the Big Bang cosmology, see Olive 1984). What it does not explain is that the baryon to photon ratio, n_B/n_γ , is so large: $n_B/n_\gamma \approx 10^{-9-10}$ or why there is matter at all. Although the above statement may seem paradoxical, it is however true if one assumes that at high temperatures the Universe started from a baryon-antibaryon symmetric state. Then, as the Universe cooled down, baryons and antibaryons were kept in thermal equilibrium by collisions until the annihilation rate started to dominate over pair creation, and as a consequence baryons and antibaryons began to disappear from the Universe. However, this process stops when the Hubble expansion of the Universe overwhelms the annihilation rate, which can be computed to take place when $T \approx 22$ MeV (Steigman 1976). After this the baryon to photon ratio is frozen with a value $n_B/n_\gamma \approx 10^{-18}$ (Steigman 1976), a very small number indeed and in strong conflict with the observed ratio. Moreover, in this case the Universe should contain an equal amount of matter and antimatter, which appears not to be the case.

In the Big Bang cosmology $n_B/n_\gamma \approx 10^{-9}$ must be taken as a boundary condition. However, it can be produced dynamically, as discussed by Sakharov, provided C and CP are violated in the early Universe, if departure from thermal equilibrium is possible (which is easily realized in an expanding Universe) and if the baryon number is not conserved. It is remarkable that these conditions can be met if we turn to a seemingly unrelated topic, the unification of all interactions (for reviews on unification, see Ellis 1983a, 1984b; Langacker 1981), which of course is the driving force behind the present hectic experimental activity for finding nucleon decay (for a review of experimental aspects of the nucleon decay, see Koshiha 1984, Perkins 1984).

2. - GRAND UNIFIED THEORIES

2.1 The Standard Model

Modern high energy physics is founded on gauge theories (for a review, see Abers and Lee 1973). They specify the interactions of matter fermions (quarks and leptons) with radiation (photons, gluons and the electroweak bosons $W^\pm$ and Z^0). If gauge theories are supplemented by scalar particles, the Higgs bosons, one can achieve a spontaneous breakdown of some (or all) of the gauge symmetries, whereby the gauge bosons become massive and the Higgs bosons responsible for the symmetry breakdown provide the longitudinal degrees of freedom of the massive gauge bosons. The correct low energy theory valid at scales $\lesssim 0(100)$ GeV is widely believed to be a gauge theory based on the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$, spontaneously broken down to $SU(3)_C \times U(1)_{em}$, where $SU(3)_C$ refers to the colour group, $SU(2)_L$ to the weak isospin and $U(1)_Y$ to the hypercharge. The generator of $U(1)_{em}$, the electric charge operator, is a linear combination of the hypercharge and of the neutral generator of $SU(2)_L$ (Glashow 1961; Weinberg 1967; Salam 1968).

A gauge theory based on a Lie group G is given once the fermion and Higgs representations are given (the vector gauge fields are in the adjoint representation of the group). The interactions of particles are then fixed by gauge invariance. In the absence of Higgses, the Lagrangian is simply

$$\mathcal{L} = -\bar{f}_i \not{D}^{ij} f_j - \frac{1}{4} F_{\mu\nu}^a F^{a\mu\nu} \quad (1)$$

where

$$F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + if_{abc} A_\mu^b A_\nu^c \quad (2)$$

is the non-Abelian field strength, where A_μ^a is the gauge field and f_{abc} are the structure constants of the Lie algebra which in the Abelian case, for instance for $U(1)_{em}$, vanish. The gauge covariant derivative D_μ^{ij} is

$$D_\mu^{ij} = \delta^{ij} \partial_\mu + ig A_\mu^a (T^a)^{ij} \quad (3)$$

where $(T^a)^{ij}$ is the generator of the group in the representation of the fermions f_i . In the case of QCD the quarks occupy the vector representation of $SU(3)_C$.

Since its gradual emergence, QCD has become accepted as the correct theory of strongly interacting quarks inside baryons and mesons (for a review, see Politzer 1974; Yndurain 1983). Experimentally these interactions have been proved by deep inelastic scattering, which give information about QCD at large Q^2 , where the perturbative approach is valid. Unfortunately, such static properties of hadrons as their masses in terms of basic QCD quantities are at the moment beyond calculations because of the non-perturbative character of QCD at small Q^2 . However, numerical lattice simulations are able to provide some information about this region with very encouraging results (for a review of lattice calculations, see Creutz 1983).

The most successful theory of physics, as far as the numerical accuracy of predictions is concerned, is without doubt quantum electrodynamics. It predicts correctly the anomalous magnetic moment of the muon to an amazing accuracy of $|a_{\text{exp}} - a_{\text{th}}| \approx 4 \times 10^{-9}$, well within experimental errors (see Bailey et al. 1977; Combley, Farley and Picasso 1981). Since the spectacular discoveries of the weak bosons $W^\pm$ and Z^0 at CERN by the UA1 and UA2 collaborations (Arnison et al. 1983a,b; Banner et al. 1983; Bagnaia et al. 1983), there has been no doubt that $U(1)_{\text{em}}$ is the low energy remnant of a spontaneously broken electroweak gauge theory $SU(2) \times U(1)$ (which has been the source of five Nobel prizes during the last five years). The breakdown of $SU(2) \times U(1)$ takes place when an $SU(2)$ doublet Higgs field H obtains a non-vanishing vacuum expectation value $\langle H \rangle$. The charged components of H are eaten up by $W^\pm$, but one real neutral field remains as a physical particle. Such a particle is still waiting to be found experimentally. The recently announced discovery of the top quark by the UA1 collaboration (Arnison et al. 1984b) serves also as an important indirect verification of the $SU(2) \times U(1)$ structure of the electroweak interactions. Without the top, or other additional fermions, the zoo of the observed fermions would have had gauge anomalies in $SU(2) \times U(1)$ (Adler 1969; Bell and Jackiw 1969). These are potentially very dangerous because they destroy gauge invariance and renormalizability; luckily enough, with the top, the Glashow-Weinberg-Salam is free of anomalies. Therefore, it is not surprising that the $SU(3)_C \times SU(2)_L \times U(1)_Y$ model of interactions is the Standard Model of today.

2.2 Grand Unification

For several reasons, the Standard Model cannot be the whole story. It contains three independent coupling constants g_3 , g_2 and g_1 for each group factor and gives no clue why, e.g.,

$$\sin^2 \theta_W = e^2/g^2 = 0.215 \pm 0.014 \quad (4)$$

as measured from νN and ep scatterings (for a recent review, see Winter 1984). This result agrees also with the more recent and independent estimate from the measurements of M_W and M_Z at the CERN $p\bar{p}$ collider, which yields $\sin^2 \theta_W(M_W) = 0.226 \pm 0.010$ (see Arnison et al. 1983; Bagnaia et al. 1983). In the Standard Model there is no explanation for the quantization of the electric charge: why are the positron and the proton electric charges equal to an accuracy that is better than one part in 10^{-20} ? Why do the quark and the lepton masses differ? Why do fermions appear in generations, that is in groups of fifteen two-component fields as depicted in Table 1? Etc., etc.

The quantities like $\sin^2 \theta_W$ are measured, however, at energy scales $\mu \lesssim O(M_W)$, and because of radiative corrections, the values of all parameters of the Standard Model are subject to changes when the energy scale where interactions take place is changed (Georgi, Quinn and Weinberg 1974; Buras, Ellis, Gaillard and Nanopoulos 1978). With three generations of fermions and to one-loop, one can calculate, using the so-called renormalization group equations, that

$$\alpha_k^{-1}(Q) = \alpha_k^{-1}(Q_0) + b_k \ln Q^2/Q_0^2 \quad (5)$$

where $k = 1, 2, 3$ runs over the group factors of the Standard Model [$U(1)$, $SU(2)$ and $SU(3)$, respectively], and for one Higgs doublet

$$\begin{aligned} b_1 &= -4N_g/12\pi \\ b_2 &= (22-4N_g)/12\pi \\ b_3 &= (33-4N_g)/12\pi \end{aligned} \quad (6)$$

where N_g is the number of generations. This fact allows for unification because it makes it possible that the coupling constants, although very different at low energies, actually meet at some higher energy scale M_X . Indeed, one may argue that all these indications strongly point towards the existence of a single grand unified theory of strong and electroweak interactions. Because the running of the coupling constants in energy is logarithmic and thus very slow (which is a reflection of the renormalizability of the theory), it is evident that the scale of unification is bound to be very high, $M_X \gg M_W$. Typically, with the Standard Model fermion content only, one finds $M_X \approx O(10^{15})$ GeV. A schematic picture of

the running of the coupling constants is depicted in Fig. 1. Of course, other quantities, like the masses, are also subject to a similar logarithmic scaling.

Given the scaling of the coupling constants, the grand vision one envisages is the following. Above the unified scale M_X , interactions are unified into a semi-simple compact group G , which at M_X breaks spontaneously down into $SU(3)_C \times SU(2) \times U(1)_Y$ by the Higgs mechanism:

$$G \xrightarrow[\langle \phi \rangle]{} SU(3)_C \times SU(2) \times U(1)_Y \xrightarrow[\langle H \rangle]{} SU(3)_C \times U(1)_{em} \quad (7)$$

Here ϕ is some Higgs field, the representation of which, along the fermion representations, one must specify to build a grand unified model. When it acquires a non-zero vacuum expectation value $\langle \phi \rangle$ as a result of a phase transition which occurs when an energetically more favourable minimum appears in the potential energy of ϕ , the gauge symmetry breaks down.

Georgi and Glashow (1974) showed that the simplest possible grand unified theory (GUT) that contains $SU(3)_C \times SU(2) \times U(1)_Y$ and is free of the gauge anomalies can be based on the group $SU(5)$. There one generation of left-handed fermions occupies a reducible representation $\bar{5} + 10$:

$$\bar{5} \sim \begin{pmatrix} d_1^c \\ d_2^c \\ d_3^c \\ e^- \\ \nu_e \end{pmatrix}_L, \quad 10 \sim \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & u_3^c & -u_2^c & -u_1 & -d_1 \\ -u_3^c & 0 & u_1^c & -u_2 & -d_2 \\ u_2^c & -u_1^c & 0 & -u_3 & -d_3 \\ u_1 & u_2 & u_3 & 0 & -e^+ \\ d_1 & d_2 & d_3 & e^+ & 0 \end{pmatrix}_L \quad (8)$$

where $d^c = C\gamma^0 T d^*$ is the charge conjugated state (and similarly u^c). It is encouraging that this anomaly-free set is just what is needed to accommodate a generation of the known fermions and nothing else, since under $SU(3) \times SU(2)$, $\bar{5} = (\bar{3}, 1) + (1, 2)$ and $10 = (3, 2) + (\bar{3}, 1) + (1, 1)$. For three generations, the representation content of the $SU(5)$ model must be triplicated; this is typical of GUTs, which usually are not able to solve the generation puzzle (see Langacker 1981). The spontaneous breakdown pattern (7) can be achieved by two sets of Higgs field, $\phi \sim 24$ and $H \sim 5$, where the former breaks $SU(5)$ into the wanted direction when

$$\langle \phi \rangle = v \text{diag} \left(1, 1, 1, -\frac{3}{2}, -\frac{3}{2} \right) \quad (9)$$

As a consequence, the gauge bosons $X \sim (3, 2)$ and $Y \sim (\bar{3}, 2)$ get superheavy masses $M_X^2 = M_Y^2 = (25/8) g_{\text{GUT}}^2 v^2$ (Buras, Ellis, Gaillard and Nanopoulos 1978), where g_{GUT} is the grand unified coupling constant. The Higgs H contains a weak doublet and breaks $SU(2) \times U(1)$ to $U(1)_{\text{em}}$ with a vacuum expectation value (VEV) $\langle H \rangle \approx 250$ GeV, giving at the same time masses to $W^\pm$ and Z^0 and to the fermions, which can have gauge invariant Yukawa couplings of the type $f(10)f(10)H^+(\bar{5})$ and $f(\bar{5})f(10)H(\bar{5})$.

In GUTs the choice of representations dictates how the electric charge operator Q is embedded into the GUT group G . Because G is compact, the trace of its generators vanishes. Since Q is some linear combination of these generators, one has also $\text{Tr } Q = 0$, where the trace runs over a given irreducible representation. In particular, in the $SU(5)$ model from, e.g., $\text{Tr } Q_5 = 0$, it follows that $Q_d = 1/3 Q_e$. Therefore, the fact that $|Q_e| = |Q_p|$ is no mystery any longer. Moreover, GUTs provide us with boundary conditions like $g_i(M_X) = g_{\text{GUT}}(M_X)$; the grand unified scale M_X can be calculated from the low energy input, which yields (at one loop)

$$\ln \frac{M_X}{M_W} = \frac{\pi}{11\alpha(M_W)} \left[1 - \frac{8}{3} \frac{\alpha(M_W)}{\alpha_3(M_W)} \right] \quad (10)$$

where $\alpha = e^2/4\pi$. In $SU(5)$, one finds, when taking into account two-loop and threshold effects (Marciano and Sirlin 1981; Llewellyn Smith, Ross and Wheeler 1981; Goldman and Ross 1980; Ellis, Gaillard, Nanopoulos and Rudaz 1980; Binétruy and Schücker 1981; Hall 1981; Antoniadis, Bouchiat and Iliopoulos 1980; Antoniadis, Kounnas and Roiesnel 1982; Chang, Das and Perez-Mercader 1981a,b)

$$M_X = (1.5 \pm 0.5) \times 10^{15} \Lambda_{\overline{\text{MS}}} \quad (11)$$

where $\Lambda_{\overline{\text{MS}}} = (150 \pm 50)$ MeV is the QCD scale parameter evaluated in the modified minimal subtraction scheme (we will return to the determination of $\Lambda_{\overline{\text{MS}}}$ in Section 4.1). Therefore, one is now able to calculate $\sin^2 \theta_W$ at $\mu = M_W$, and in $SU(5)$ one finds (Georgi, Quinn and Weinberg 1974) at one loop

$$\sin^2 \theta_W = \frac{3g_1^2(\mu)}{3g_1^2(\mu) + 5g_2^2(\mu)} \quad (12)$$

$$= \frac{3}{8} \left[1 - \frac{\alpha(M_W)}{4\pi} \left(\frac{110 - N_H}{9} \ln \frac{M_X^2}{\mu^2} \right) \right]$$

where N_H is the number of the Higgs doublets, and one arrives at the theoretical prediction, taking also into account higher order and threshold effects (Marciano and Sirlin 1981; Llewellyn Smith, Ross and Wheater 1981)

$$\sin^2 \theta_W(M_W) = 0.214 \begin{matrix} + 0.004 \\ - 0.003 \end{matrix} \quad (13)$$

where the errors reflect uncertainties in the QCD scale Λ . This should be compared with the experimental quantity (4); note that a priori the theoretical prediction could have given anything between 0 and 1.

Since in GUTs quarks and leptons are usually placed into the same representation [see (8)], their masses must be equal at scales $\mu > M_X$. Typically, one finds that m_q/m_l = some Clebsch-Gordan coefficient $O(1)$. Below M_X , there will arise differences because quarks and leptons have different renormalization group scalings under $SU(3)_C \times SU(2)_L \times U(1)_Y$ (Buras, Ellis, Gaillard and Nanopoulos 1978). In the minimal $SU(5)$ one finds $m_b(M_X) = m_\tau(M_X)$ and

$$\frac{m_b(\mu)}{m_\tau(\mu)} = \left[\frac{\alpha_3(\mu)}{\alpha_{GUT}} \right]^{\frac{12}{33-4N_g}} \quad (14)$$

where $\alpha_3(\mu)$ is the QCD coupling constant with the value $\alpha_3(M_W) \approx 0.12$, $\alpha_{GUT} \approx 1/42$ is the GUT coupling constant as calculated from (5) and N_g is again the number of generations. For three generations, one successfully predicts (Buras, Ellis, Gaillard and Nanopoulos 1978) at the threshold $\mu = 2m_b$ that $m_b/m_\tau \approx 2.7-3$. In fact, for four generations the prediction would be clearly off and therefore one can take this also as an indication for the existence of three generations only (Nanopoulos and Ross 1979, 1982; Hill 1981; Machacek and Vaughan 1981). For lighter generations, the prediction (14) does not work so well, which perhaps is not so worrisome for grand unification, because with light quarks one enters

into the strong coupling regime of QCD where the simple perturbative approach (14) can be expected to fail. In addition, the particle masses might get small contributions from other (non-renormalizable, gravitational) sources that do not affect large ($\gg 100$ MeV) masses, but can change thoroughly the light mass ratios (Ellis and Gaillard 1979; Nanopoulos and Srednicki 1983).

Finally, since quarks and leptons are in GUTs in the same representations, there will be some generators that will transform quarks into leptons and vice versa. As these generators correspond to the gauge bosons, it is obvious that GUTs will in general have both baryon and lepton number violating interactions, which can manifest themselves, e.g., in proton decay. This instability of nucleons is an unavoidable feature of most GUTs, although it is possible to have stable protons also in GUTs by certain devious choices (Gell-Mann, Ramond and Slansky 1978). Baryon number violation can also lead to such $\Delta B = 2$, $\Delta L = 0$ processes as $n-\bar{n}$ oscillations (Kuzmin 1970; Marshak and Mohapatra 1980), which however, need not occur at an observable rate, as is the case of the minimal SU(5) which conserves B-L (for a recent review of the subject, see Baldo-Ceolin 1984). Lepton number violation enables neutrinos to acquire Majorana masses with mass terms of the type $\nu^T C \nu$, and these terms together with neutrino Dirac mass terms may well be responsible for the smallness of the neutrino masses (Gell-Mann, Ramond and Slansky 1979; Yanagida 1979). Such a Majorana character of the neutrino can, in principle, be resolved if one observes the neutrinoless double beta decay $Z \rightarrow Z - 2 + 2e$ (for a review, see Primakoff and Rosen 1981).

Baryon number violation provides an exciting probe of unification and the large energy scales that otherwise would be inaccessible to us. Before turning to this topic we will first, however, review the supersymmetric extensions of GUTs.

2.3 Supersymmetric GUTs

Ordinary GUTs, although in possession of several exciting features, suffer from a malady called the hierarchy problem (for a review, see Ellis 1983a, 1984b). As we saw in the previous section, the scale of unification is necessarily very large, $M_X/M_W \gtrsim 0(10^{13})$, which means that the Higgs masses (or VEVs) that are the sources of these scales should also exhibit the same hierarchy. However, such a hierarchy is not stable under radiative corrections (see Fig. 2), but the Higgs mass will in general be subject to radiatively produced shifts of the order of

$$\delta m_H^2 \simeq \theta(M_X^2) \quad (15)$$

It has also been pointed out (Hawking, Page and Pope 1979, 1980) that Higgs fields propagating in space-time foam would acquire masses of the order of the cut-off (the Planck mass?). This is therefore a problem of the low energy $SU(3)_C \times SU(2)_L \times U(1)_Y$ theory alone and independent of our attempts to unify interactions. There are two very different propositions for solving this dilemma (for a general review, see Nanopoulos 1984b). One way out is to resolve the elementary nature of the Higgs field at some scale typically $O(1)$ TeV by making it a composite object. The early technicoloured attempts, however, failed the test of flavour changing neutral currents, and at the moment there is no dynamically acceptable composite model that would account for the Higgs particle. Therefore, we will turn to the other leading alternative, namely supersymmetry (SUSY) (Gol'fand and Likhtman 1971; Volkov and Akulov 1973; Wess and Zumino 1974a; for an early review, see Fayet and Ferrara 1977). This is a novel symmetry between bosons b and fermions f . Let Q be the generator of SUSY; then

$$Q|f\rangle = |b\rangle \quad ; \quad Q|b\rangle = |f\rangle \quad (16)$$

The generators Q^i ($i = 1, \dots, N$) obey the anticommutation rule

$$\{Q_\alpha^i, Q^{+\dot{\alpha}j}\} = -2(\sigma^\mu)_\alpha^{\dot{\alpha}} P_\mu \delta^{ij} \quad (17)$$

where P_μ is the energy-momentum operator and α is a spin-index; SUSY-generators are therefore spinorial. Particle physics models can be built using one or several SUSY-generators. In the case of $N = 1$ SUSY, one can ascribe particles to the gravity multiplet $(2, 3/2)$, to the gauge multiplet $(1, \frac{1}{2})$ and to the chiral multiplets $(\frac{1}{2}, 0)$, where the numbers refer to the spins of the fields in the above supermultiplets. As models with $N > 2$ do not appear to be wholly realistic at the moment (Del Aguila et al. 1984; Enqvist and Maalampi 1984; Kalara et al. 1984), we will henceforth focus on $N = 1$ SUSY only. There, one finds cancellations amongst the fermion and boson loops in the radiative corrections (see Fig. 3), so that in SUSY one finds

$$\delta m_H^2 = \theta(m_S^2) \quad (18)$$

where $m_S^2 = |m_B^2 - m_F^2|$ is the scale of supersymmetry breaking in the low energy theory. To preserve the gauge hierarchy, we therefore need SUSY GUT models

where $m_S \lesssim 0(1)$ TeV (Witten 1981; Dimopoulos and Georgi 1981; Sakai 1982). The improved ultra-violet behaviour of supersymmetry is due to the high degree of symmetry between bosons and fermions, and in fact, one can prove no-renormalization theorems that are valid to all orders in perturbation theory (Wess and Zumino 1974b; Iliopoulos and Zumino 1974; Ferrara, Iliopoulos and Zumino 1974; Grisaru, Siegel and Roček 1979).

In SUSY matter, quarks, leptons and Higgses are in chiral supermultiplets ϕ_i , where they are accompanied by their SUSY partners. The expected spectrum of supersymmetric particles is illustrated in Table 2. Interactions between chiral multiplets can be expressed in terms of the superpotential $f = f(\phi_i)$, which is an analytic function of ϕ_i and renormalizability forces it to be at most of third order in the fields. This analyticity is the reason why in SUSY, one has to have two weak Higgs doublets, because the u-sector Yukawa coupling would be of the form QH^+U (where capital letters denote chiral supermultiplets), which is not supersymmetric.

In globally supersymmetric theory, the scalar potential can be written as (see Witten 1981 and references therein)

$$V = \sum_i \left| \frac{\partial f}{\partial \phi_i} \right|^2 + \frac{1}{2} D^2 \quad (19)$$

where ϕ_i denote the scalar components of the chiral multiplets, and $D^\alpha = g \phi_i^+ (T^\alpha)_j^i \phi^j$. Supersymmetry must naturally be broken. Early attempts concentrated on models with spontaneously broken supersymmetry which did not turn out to be a feasible alternative (for a recent review, see Nanopoulos 1984a). Dynamical breaking of supersymmetry is problematic (Witten 1981, 1982; Ceccotti and Girardello 1982; Affleck, Harvey and Witten 1982; Affleck, Dine and Seiberg 1983, 1984; Rossi and Veneziano 1984; Amati, Rossi and Veneziano 1984). The only remaining alternative is to add supersymmetry breaking terms that are soft in the sense that they do not affect the no-renormalization properties of SUSY that are the heart of the relation (18). The most general set of such terms is given by

$$\mathcal{L}_{SOFT} = -m_\lambda \bar{\lambda}^\alpha \lambda^\alpha + \sum_i m_i^2 |\phi_i|^2 + (m_{ijk} \phi^i \phi^j \phi^k + h.c.) \quad (20)$$

where λ^α are the SUSY partners of the gauge bosons (gauginos), and the various mass parameters in (20) should again be $\lesssim 0(1)$ TeV. In fact, such soft SUSY

breaking terms can be obtained from $N = 1$ supergravity, where the chiral and gauge supermultiplets have been coupled to gravity (Cremmer et al. 1978, 1979, 1982, 1983). There, the goldstino of the spontaneously broken SUSY is eaten up by the gravitino, the spin $3/2$ superpartner of the graviton, which becomes massive. By now, several phenomenologically viable supergravity models exist (for reviews, see Nanopoulos 1983a; Arnowitt, Chamseddine and Nath 1983a,b; Nilles 1984; Ellis 1984a,b), and although it is conceivable that such soft phenomenological SUSY breaking could eventually be obtained from higher dimensional, Kaluza-Klein type SUSY theories (for a review, see Fré 1984), here we will concentrate only on the already existing alternatives. Supergravity models offer exciting possibilities for explaining dynamically all scales below the Planck scale $M_P \approx 10^{19}$ GeV. There exists a class of models, the no-scale models, which have a naturally vanishing cosmological constant at the tree level and in which the scales of supersymmetry and $SU(2) \times U(1)$ breakdown are determined by radiative corrections (Ellis, Lahanas, Nanopoulos and Tamvakis 1984; Ellis, Kounnas and Nanopoulos 1984a,b,c; Ellis, Enqvist and Nanopoulos 1984a,b). In these models, it may also be possible to solve the notorious strong CP-problem by a dynamical mechanism that relaxes θ_{QCD} to zero (Ellis, Enqvist and Nanopoulos 1984a,b).

The absence of any new charged particles at PETRA and PEP means that the scalar partners of the quarks and the leptons, the squarks and the sleptons, and the fermionic partners of the charged Higgs and $W^\pm$ s, the charged higgsinos and winos, must have masses that lie above $O(20)$ GeV. The small number of the UAL "zen" (monojet) events has been used to set a lower bound of the order of 40 GeV on the squark and the gluino masses (Arnison et al. 1984). The 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 for gluino masses. Negative results from a) $e^+e^- \rightarrow \tilde{\gamma}\tilde{\gamma}\tilde{\gamma} \rightarrow \gamma + \text{nothing}$ and b) $e^+e^- \rightarrow \tilde{\gamma}\tilde{Z}$ searches (Becker 1984; Yamada 1984; Prepost 1984) imply $m_{\tilde{e}} > O(20;30)$ GeV for $m_{\tilde{\gamma}} < O(6,1)$ GeV in the case a), while $m_{\tilde{e}} > O(100,30)$ GeV for $m_{\tilde{Z}} < O(10,40)$ GeV in the case b). No experimental lower bounds exist on colourless and electrically neutral sparticles. There are several candidates for the lightest neutral sparticle, including the sneutrino $\tilde{\nu}$, the photino $\tilde{\gamma}$, a neutral higgsino $\tilde{H}$ 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 many currently popular models. Plausible possibilities in many models are also that $\tilde{\gamma}$ and/or the neutral $\tilde{H}$ may be considerably lighter than the other sparticles.

There are also severe constraints on possible SUSY models that come from the K_1 - K_2 mass difference (Ellis and Nanopoulos 1982; Barbieri and Gatto 1982;

Campbell 1983; Inami and Lim 1982; Lahanas and Nanopoulos 1983; Baulieu, Kaplan and Fayet 1984). The flavour changing neutral gaugino interactions (see Fig. 4) would give too large a contribution unless the splitting between squarks is

$$m_{\tilde{c}}^2 - m_{\tilde{u}}^2 \lesssim \mathcal{O}(1)(m_c^2 - m_u^2) \quad (21)$$

Other constraints, such as the $g-2$ of the muon do not give any essential new limitations.

Supersymmetric GUTs can be built from supermultiplets in the same way as ordinary GUTs. The splitting of the scales M_W and M_X within a particle multiplet (say, the 5 of the Higgs supermultiplet H) can be imposed on the theory at high scales. The no-renormalization theorem then guarantees the perturbative stability of this splitting. In $SU(5)$, one chooses the following superpotential couplings between $\phi \sim \underline{24}$, $H \sim 5$ and $\bar{H} \sim \bar{5}$:

$$\mathcal{P}(\phi, H, \bar{H}) = \lambda \bar{H}^\alpha (\phi_\alpha^\beta + \frac{3}{2} m \delta_\alpha^\beta) H_\beta \quad (22)$$

When ϕ takes a VEV breaking $SU(5)$ to $SU(3) \times SU(2) \times U(1)$ as in (9), the mass of the doublet components $H^{4,5}$, $\bar{H}^{4,5}$ are $3/2(m-v) = 0$ if we choose $m = v$, while the triplet components remain superheavy ($\mathcal{O}(v)$). A more elegant possibility to achieve the doublet-triplet splitting is to use the missing partner mechanism (Grinstein 1982; Masiero, Nanopoulos, Tamvakis and Yanagida 1982). There, one introduces more Higgs supermultiplets which contain additional Higgs triplets that can combine with the triplets in the $5 + \bar{5}$ to acquire large masses, but no extra weak doublets to mix with the light doublets. The lowest suitable representations of $SU(5)$ are $\theta + \bar{\theta} = 50 + \bar{50}$. These can be coupled to the $5 + \bar{5}$ by using a $\Sigma \sim 75$ of Higgses which replaces the more conventional $\Sigma \sim 24$. This may, however, cause problems with the perturbative validity of the model below M_P because with such a large representation the GUT coupling constant runs very rapidly. Here one can invoke the non-renormalizable couplings of supergravity to replace Σ by superpotential couplings of the type $H\Sigma^2\theta/M_P$ (Kounnas, Nanopoulos, Quirós and Srednicki 1983). In both cases, the Higgs triplet mass matrix has the structure

$$(\bar{H}_3, \bar{\theta}_3) \begin{pmatrix} 0 & \theta(M_X) \\ \theta(M_X) & \epsilon \end{pmatrix} \begin{pmatrix} H_3 \\ \theta_3 \end{pmatrix} \quad (23)$$

where the subscript 3 refers to the coloured triplet components, and the entry ϵ may be small or even $O(M_P)$. In the latter case, there will be one relatively light triplet with a mass $O(M_X^2/M_P)$ which will then be of great importance for baryon asymmetry production and also for the proton decay, as we will discuss later.

Because the particle content of the supersymmetrized Standard Model is more than doubled (recall the introduction of an extra Higgs supermultiplet), the renormalization group equations of supersymmetric $SU(3)_C \times SU(2) \times U(1)$ differ from those of the Standard Model. Instead of (6), one has now

$$\begin{aligned} b_1 &= \left(-\frac{3}{5} - 6N_g\right)/12\pi \\ b_2 &= (15 - 6N_g)/12\pi \\ b_3 &= (27 - 6N_g)/12\pi \end{aligned} \quad (24)$$

for two light doublets. These yield a higher GUT scale (Dimopoulos, Raby and Wilczek 1981; Ibáñez and Ross 1982; Einhorn and Jones 1982; Ellis, Nanopoulos and Rudaz 1982):

$$M_X = 6 \times 10^{16} \Lambda_{\overline{MS}} \simeq 10^{16} \text{ GeV} \quad (25)$$

This, as we will see, is important for proton decay. For $\sin^2\theta_W$ one obtains

$$\sin^2\theta_W(M_W) = 0.236 \pm 0.003 \quad (26)$$

Remarkably, the successfully predicted m_b/m_τ -ratio remains in SUSY numerically practically unchanged from the Standard Model ratio $(m_b/m_\tau)_{SM}$. With three generations

$$\frac{(m_b/m_\tau)_{\text{SUSY}}}{(m_b/m_\tau)_{\text{SM}}} = \frac{[\alpha_3(M_W)/\alpha_{\text{GUT}}]^{\frac{8}{27-6N_g}}}{[\alpha_3(M_W)/\alpha_{\text{GUT}}]^{\frac{12}{33-4N_g}}} = 1.0 \quad (27)$$

Therefore, SUSY GUTs preserve all the desirable predictions of the ordinary GUTs. However, baryon decay in SUSY GUTs is very different from that of ordinary GUTs, as we will discuss in the following sections.

3. - BARYON NUMBER VIOLATION

3.1 Baryon number violation in GUTs

In the Standard Model with minimal particle content, baryon and lepton number conservations are consequences of gauge invariance and renormalizability (Nanopoulos 1973; Weinberg 1973). These ensure that B and L are global symmetries of the theory. However, if we enlarge the particle content or extend the gauge group beyond $SU(3) \times SU(2) \times U(1)$, this is no longer necessarily true. The low energy $SU(3) \times SU(2) \times U(1)$ theory is however presumably only an effective theory and may contain effective operators that can cause baryon number violation. Such operators arise from GUTs when superheavy degrees of freedom have been integrated out, to leave behind only the minimal particle content much below the superheavy energy scale M_X . They must naturally be $SU(3) \times SU(2) \times U(1)$ invariant, although they will not be renormalizable but carry dimensions $d > 4$, and must therefore be accompanied by coupling constants with dimensions $(\text{mass})^{4-d}$. All such baryon number violating operators can be classified in terms of their $SU(3) \times SU(2) \times U(1)$ properties (Weinberg 1979, 1980; Wilczek and Zee 1979). In conventional GUTs, bosonic terms in the Lagrangian do not carry B and therefore B-number violation must involve at least four-fermion, i.e., $d = 6$, interactions. Schematically, one finds

$$\begin{aligned} O_6 &= qqq\ell & (d=6) \\ O_7 &= qqq\ell W & (d=7) \\ &\vdots & \end{aligned} \quad (28)$$

where W stands for any gauge boson, Higgs boson or space-time derivative. O_6 has $\Delta B = \Delta L$, whereas O_7 has $\Delta B = -\Delta L$. We expect $d = 7$ operators to be suppressed relative to $d = 6$ operators by a factor of $(\Lambda_{\text{QCD}} \text{ or } M_W)/M_X$ if the GUT has only one mass scale (i.e., no intermediate mass scales). Higher order operators would be even more suppressed. Therefore, we expect B-L to be conserved to a high accuracy in nucleon decay, and protons should not be able to decay into leptons but should go to antileptons instead. Another feature of the operator O_6 is that it does not allow for $\Delta B = \Delta S$ decays. Therefore, the decays $p \rightarrow \bar{K}^0 \lambda^+$ or $n \rightarrow K^- \lambda^+$ are disallowed. If the origin of O_6 is due to one-particle exchanges at the superheavy mass scale, there can be only five kinds of superheavy vectors or scalars having $SU(3) \times SU(2) \times U(1)$ invariant but B-violating interactions with a pair of ordinary fermions. These are the following: $SU(3)$ triplet, $SU(2)$ doublet vectors with the electric charges $(4/3, 1/3)$ and $(2/3, -1/3)$; $SU(3)$ triplet, $SU(2)$ singlet scalars $-1/3$ and $-4/3$; and $SU(3)$ triplet, $SU(2)$ triplet scalars $(2/3, -1/3, -4/3)$. As the coupling constant of O_6 has dimensions $(\text{mass})^{-2}$ and the only relevant scale in GUTs is M_X , one can argue purely on dimensional grounds that the order of magnitude of the proton lifetime must be

$$\tau_p \approx M_X^4 / m_p^5 \quad (29)$$

where m_p is the proton mass, an expectation that is verified by more detailed calculations, as we will see. In $SU(5)$ such a $d = 6$ operator arises from an exchange of a superheavy gauge boson X (Higgs mediated interactions are suppressed relative to this) as depicted in Fig. 5.

In ordinary $SU(5)$, one can write down the effective B-violating $d = 6$ interaction, taking into account mixings between the first and the second families (Ellis, Gaillard, Nanopoulos and Rudaz 1980)

$$\begin{aligned} \frac{1}{4} \mathcal{L}_{GU} = e^{i\phi} \frac{1}{\sqrt{2}} G_{GU} & \left[(\epsilon_{ijk} \bar{u}_{kL}^c \gamma_\mu u_{jL}) \left\{ [(1+\cos^2\theta_c) \bar{e}_L^+ + \right. \right. \\ & + \sin\theta_c \cos\theta_c \bar{\mu}_L^+] \gamma^\mu d_{iL} + [(1+\sin^2\theta_c) \bar{\mu}_L^+ + \sin\theta_c \cos\theta_c \bar{e}_L^+] \gamma^\mu s_{iL} + \\ & + \bar{e}_R^+ \gamma^\mu d_{iR} + \bar{\mu}_R^+ \gamma^\mu s_{iR} \} - [\epsilon_{ijk} \bar{u}_{kL}^c \gamma_\mu (d_{jL} \cos\theta_c + s_{jL} \sin\theta_c)] \\ & \times [\bar{\nu}_{eR}^c \gamma^\mu d_{iR} + \bar{\nu}_{\mu R}^c \gamma^\mu s_{iR}] \Big] + h.c. \end{aligned} \quad (30)$$

where $G_{GU}/\sqrt{2} = g_{GU}^2/8M_X^2 \approx g^2/8M_X^2$. This form is telling us in particular that

baryon decay cannot be Cabibbo-rotated away. It is also evident from (30) that the modes λ^+ + strange and μ^+ + non-strange are disfavoured. In fact, one finds that

$$\frac{\Gamma(N \rightarrow \mu^+ + \text{non-strange})}{\Gamma(N \rightarrow e^+ + \text{non-strange})} = \frac{\sin^2 \theta_c \cos^2 \theta_c}{1 + (1 + \cos^2 \theta_c)^2}$$

$$\frac{\Gamma(N \rightarrow e^+ + \text{strange})}{\Gamma(N \rightarrow \mu^+ + \text{strange})} = \frac{\sin^2 \theta_c \cos^2 \theta_c}{1 + (1 + \sin^2 \theta_c)^2} \quad (31)$$

The above result need not hold in GUTs which contain antisymmetric contributions to fermion mass matrices but is true in, e.g., SU(5) and SO(10) [with only the 10 of SO(10) contributing to the fermion mass matrices].

In supersymmetric GUTs there can appear also baryon number violating $d = 5$ operators (Weinberg 1982; Sakai and Yanagida 1982). This happens because quarks and their scalar partners occupy the same supermultiplet and carry therefore the same global quantum numbers, in particular the baryon number. In the superspace formalism such a $d = 5$ operator is of the form

$$O_5 = \int d^2\theta \frac{\lambda_Y^2}{M_{H_3}} \phi_f \phi_f \phi_f \phi_f \quad (32)$$

where ϕ_f stands for the chiral quark and lepton superfields, λ_Y is a Yukawa coupling and the integration $\int d^2\theta$ is performed over the (half of the) Grassmannian superspace. The product of the chiral fields in (32) must be overall symmetric (the fields are bosonic) and a colour singlet, which necessitates an antisymmetric SU(3) tensor ϵ_{ijk} and implies antisymmetry in the flavour space. Thus the effective interaction (32) involves second and/or third generation particles, and therefore in SUSY one can expect decay modes of the proton that are very different from those of the ordinary GUTs (Ellis, Nanopoulos and Rudaz 1982; Dimopoulos, Raby and Wilczek 1982). We will return to proton decay in supersymmetric theories in more detail in Section 4.2.

3.2 Monopole catalyzed nucleon decay

't Hooft (1974) and Polyakov (1974a,b) were the first to point out that finite energy monopole solutions exist in any gauge theory with a semi-simple gauge group broken spontaneously down to an exact low-energy gauge group containing the electromagnetic $U(1)_{em}$ group. In GUTs like $SU(5)$ or $SO(10)$, where the electromagnetic $U(1)_{em}$ emerges from the breaking at superhigh energies M_X , the mass of the monopole is also superheavy, $m_M \approx M_X/\alpha_{GUT} \approx 0(10^{16})$ GeV. There are stringent phenomenological limits on their abundance and flux, both from direct experimental searches and from astrophysical considerations (for a review, see Preskill 1983; Ellis 1983b). They are expected to be copiously produced at the GUT symmetry breakdown and are therefore rather an embarrassment for the Big Bang cosmology. An early inflationary period can, however, dilute the monopole density exponentially (for a review on inflation, see Nanopoulos 1983b). In fact, in the first approaches to inflation (Guth 1981; Linde 1982; Albrecht and Steinhardt 1982), one would have found typically only one monopole in the whole of our present Universe. However, in the recently proposed two-component inflation, monopole densities, although considerably diluted, may well harbour on the verge of detection (Enqvist and Nanopoulos 1984a,b; Enqvist, Kounnas, Nanopoulos and Quirós 1984).

A grand unified monopole (GUM) can be visualized as a knot in the Higgs field (see Fig. 6) that has the GUT symmetry unbroken inside the knot (Kibble 1976). It is therefore evident that GUM interactions will violate baryon number because GUT interactions do. Naively, one would have guessed that the cross-section should be geometrically small, $\sigma \propto M_X^{-2}$. However, Rubakov (1981a,b; 1982) pointed out that electric and magnetic fields around GUMs can generate, via the axial anomaly, B-violating condensates of fermion pairs within a range that extends out to their Compton wavelengths $m_f^{-1} \gg M_X^{-1}$ and consequently that the $\Delta B \neq 0$ GUM cross-sections might also be large. This was verified by later workers (Callan 1982). Therefore scatterings of monopoles with matter can catalyze proton decay that could be observed in proton decay experiments.

To estimate the cross-section for GUM catalyzed proton decay, one needs to calculate the expectation value of the chiral condensate $F = \epsilon_{\alpha\beta} \epsilon_{ij} \psi_{\alpha i}^{(1)} \psi_{\beta j}^{(2)}$ in the monopole state where α is Lorenz and i, j group are indices. The result is (Ellis, Nanopoulos and Olive 1982; Bais, Ellis, Nanopoulos and Olive 1983)

$$\langle M|F|M \rangle = \frac{1}{4\pi^2 2^{1/4} r^3} + \mathcal{O}(e^2) \quad (33)$$

That is, the cloud of condensates drops only as r^{-3} , and there are no suppression factors. The formation of many fermion condensates is also possible. The resulting picture of the monopole has a multiplet core structure with light fermions on the "surface". The monopoles can catalyze baryon decay through scattering $M + B \rightarrow \ell^+ + M + \text{mesons}$, and the cross-section can be estimated to be

$$\sigma_{\text{GUM}} = \frac{\mathcal{O}(1)}{v} \frac{\hat{\Gamma}(B \rightarrow \ell + \text{mesons})}{m_q^3} \quad (34)$$

where m_q is the effective quark mass inside the baryon B , v is the velocity of the GUM and $\hat{\Gamma}$ is related to the GUT baryon decay width Γ with the replacement $g_{\text{GUT}}^2/2 \rightarrow m_q^2$. One estimates $\hat{\Gamma} = 0(5)/(2\pi)^8 \text{ GeV}$, and altogether

$$\sigma_{\text{GUM}} \simeq \frac{\sigma_0}{v \text{ GeV}^2}, \quad \sigma_0 \simeq 10^{-3} - 10^{-5} \quad (35)$$

Bais et al. (1983) predict a hierarchy of monopole catalyzed decays $p \rightarrow e^+\pi \gg \mu^+K$ with no $p \rightarrow \nu X$, e^+K or $\mu^+\pi$. The passage of a monopole through proton decay detectors might cause, through multiple scatterings, a chain of nucleon decays, although it appears to be quite likely that the second decay would occur during the dead time of the detector which lasts the few milliseconds that it takes to register the previous event.

3.3 Baryon number violation in the early Universe

Besides the nucleon decay, the baryon number violation in GUTs has equally important consequences in the early Universe, as was already discussed in the Introduction. As detailed studies suggest (Steigman 1976) that it is highly probable that there is no antimatter in the whole Universe, there exists a large baryon-antibaryon asymmetry

$$\frac{n_B - n_{\bar{B}}}{n_\gamma} \simeq \frac{n_B}{n_\gamma} \simeq 10^{-9} - 10^{-10} \quad (36)$$

Strictly speaking, one should compare the baryon number density (n_B) of the Universe with its entropy density (s). Most of the entropy of the present Universe

is contained in the microwave radiation background and in the associated neutrino and antineutrino backgrounds. With a microwave background temperature of 3 K, the photon entropy density is $s_\gamma \approx 250/\text{cm}^3$ (for us $k_{\text{Boltzmann}} = 1$ always). The entropy density of each species of neutrino and antineutrino is calculated to be $(7/22)s_\gamma$. The quantity s_γ is very roughly equal to n_γ , but it is better to deal with entropy than photon number densities, because as long as the expansion of the Universe is adiabatic, the total entropy in any co-moving volume remains fixed, while processes in the early Universe, such as e^+e^- annihilation, can affect the number of photons. In the following we will use either n_γ or s_γ (or even s), but the reader should keep in mind the above remarks.

The baryon number density n_B is not so well known from observations. Limits on the deceleration parameter of the cosmic expansion restrict it to be less than about $3 \times 10^{-6}/\text{cm}^3$, while the mass density of visible galaxies provides a lower bound on n_B of about $3 \times 10^{-8} \text{ cm}^3$. The compatibility of the Big Bang nucleosynthesis with the observed Helium, Deuterium and especially Lithium abundances helps further to pinpoint n_B/n_γ to the range indicated in (36).

This ratio remains constant throughout the history of the Universe, as long as the baryon number is conserved and the Universe expands adiabatically. Furthermore, there are no known concentrations of antimatter, and the absence of high-energy photons, which could be the product of $B\bar{B}$ annihilation, indicates that our local cluster of galaxies does not contain any significant fraction of antimatter. Before the advent of GUTs, there was no satisfactory explanation of the ratio (36) or of the apparent matter-antimatter asymmetry.

Grand Unified Theories contain all the necessary ingredients enumerated in the Introduction for successful baryogenesis. As we have seen, baryon number violation is an unavoidable consequence of grand unification, while C and CP-violation, occurring already in the Standard Model, are naturally embedded in GUTs. Furthermore, because we live in an expanding Universe in which the expansion rate (Hubble parameter) H is temperature (or time) dependent (for a radiation dominated Universe)

$$H \approx \sqrt{\frac{8\pi}{3M_P^2}} T^2 \quad (37)$$

It is easy for GUT interactions to go out of equilibrium for a certain period in the very early Universe. It seems likely that the creation of the net baryon number (36) in GUTs is driven by decays and inverse decays of superheavy bosons, either

the gauge vectors, X, Y or the Higgs colour triplets H_3 [which in SU(5) are the gauge partners of the weak Higgs doublets]. It is assumed that the Higgs bosons H_3 while in equilibrium for $T > M_{H_3}$, will drop out of equilibrium when $T_D < T < m_{H_3}$, where T_D is the decay temperature determined by $\Gamma_H \approx H$, or

$$\Gamma_H \approx \alpha_H M_{H_3} \approx \sqrt{\frac{8\pi}{3M_P^2}} T_D^2 \quad (38)$$

Here α_H is the Yukawa fine structure constant of the heaviest available fermions; with a top quark of 40 GeV, $\alpha_H = \alpha_t \approx 3 \times 10^{-3}$. It is crucial that the Higgs colour triplets H_3 decay at $T_D < M_{H_3}$ in such a way that there is a kinematic block for creating back the "parent" particles. (For a review of the cosmological implications of grand unified theories, see Nanopoulos 1980.)

In most GUTs, the superheavy particles responsible for baryon asymmetry generation are expected to decay into pairs of quarks or into an antilepton and an antiquark (see Fig. 7)

$$X \text{ or } H_3 \rightarrow (q + \bar{q}) \text{ or } (\bar{q} + \bar{l}) \quad (39)$$

The total decay rates of the particles and antiparticles must be equal because of CPT. However, partial decay rates into conjugate channels may differ if C and CP are both violated:

$$BR \equiv BR(X \rightarrow \bar{q} \bar{l}) = \frac{\Gamma(X \rightarrow \bar{q} \bar{l})}{\Gamma_{\text{tot}}(X)} \neq \overline{BR} \equiv BR(\bar{X} \rightarrow q l) = \frac{\Gamma(\bar{X} \rightarrow q l)}{\Gamma_{\text{tot}}(\bar{X})} \quad (40)$$

where X stands for a superheavy gauge boson or a Higgs colour triplet. It can be shown that in Higgs decay a non-zero ΔB can first arise in the fourth order in couplings with

$$\Delta B \approx \text{Im Tr}(abcd) / \sum_a \text{Tr} aa^\dagger \quad (41)$$

where a, b, c and d are various coupling matrices of H_3 to fermions and are defined in Fig. 8. Roughly, the baryon to photon ratio is then

$$\frac{n_B}{n_\gamma} \approx \frac{g_X}{g_{\text{tot}}} \Delta B \approx 10^{-2} \Delta B \quad (42)$$

where g_X is the number of the helicity states of the superheavy bosons and g_{tot}

is the total number of helicity states. For instance, in the minimal SU(5) with the usual $\bar{5} + 10$ of fermions, one 24 and one pair $5 + \bar{5}$ of Higgses, one finds that the baryon asymmetry must be produced by high order loop diagrams and that $\Delta B \approx 10^{-14}$, far too small to generate the baryon to photon ratio (36). However, the use of a slightly more complicated Higgs system, e.g., two pairs $5 + \bar{5}$ of Higgses, gives much more heart-warming results (Nanopoulos and Weinberg 1979)

$$\Delta B \approx \frac{1}{32\pi} (\sqrt{2} G_F) m_b^2 \epsilon \approx 10^{-6} \epsilon \quad (43)$$

where ϵ is some CP-violating parameter in the range $(10^{-1} - 10^{-3})$. Clearly, a mere substitution of (43) into (42) shows that the predicted ratio n_B/n_γ in GUTs is in good agreement with the experimental measured one (36).

In supersymmetric GUTs, one might wonder if the large degree of symmetry of these theories would prohibit or greatly suppress the creation of the baryon asymmetry. This apprehension is, however, unnecessary (Ellis et al. 1984). As in ordinary GUTs, where with only one pentaplet of scalars one had to go to higher orders to find a non-vanishing imaginary part in the trace of the Yukawa couplings, in SUSY SU(5) with only one pair of pentaplets the one-loop interference diagram for the colour triplet decay has no imaginary part (Nanopoulos and Tamvakis 1982; Haber 1982). To avoid too large a suppression of n_B/n_γ one must invoke an additional pentaplet pair $H_2 + \bar{H}_2$. The relevant interference supergraph is shown in Fig. 9a, and one of the relevant component diagrams in Fig. 9b. In general, this interference arises from the effective superpotential

$$\begin{aligned} f_{\text{eff}} = & \mu_1 \bar{H}_1 H_1 + \mu_2 \bar{H}_2 H_2 + \mu_3 \bar{H}_1 H_2 + \mu_4 \bar{H}_2 H_1 + \\ & + \lambda_1 T \bar{T} H_1 + \lambda_2 T \bar{T} H_2 + \lambda'_1 T \bar{F} \bar{H}_1 + \lambda'_2 T \bar{F} \bar{H}_2 + \dots \end{aligned} \quad (44)$$

where T and $\bar{F}$ are respectively the 10 and $\bar{5}$ superfields containing the ordinary fermions. The effective Lagrangian relevant to baryosynthesis can be extracted from Fig. 9a and yields, after loop integration, a baryon asymmetry per H_1 decays

$$\Delta B \approx \frac{\mu_1^3 \mu_2}{M_H^4} \theta \left(\frac{\lambda^2 \phi}{16\pi^2} \right) \quad (45)$$

where ϕ denotes some CP violating phase factor arising from the Yukawa couplings,

which could be $O(1)$. M_H is a typical mass eigenstate for the Higgs triplets, and $\lambda \sim \lambda_1 \sim \lambda'_1 \sim \lambda_2 \sim \lambda'_2$. The mass parameters μ_1 and μ_2 are model dependent. For instance, in the model of Ellis et al. (1984) $\mu_1 \sim \mu_2 \sim O(M_X^5/M_P^4)$. Therefore, in that model one can obtain a large enough baryon asymmetry if $M_X \sim O(M_P)$, and in general, the result (45) shows that in SUSY GUTs no unexpected cancellations in the interference diagrams occur.

3.4 Baryon number violation without GUTs

Even without GUTs, it seems likely that baryon number will be violated at some level so that nucleons are ultimately unstable, albeit with a very long lifetime. For instance, 't Hooft (1976 a,b) has pointed out that instantons of the Standard Model can cause $\Delta B = N_G$ reactions which would allow for example

$$N + N \rightarrow \bar{N} + \text{leptons} \quad (46)$$

but the rate for this, being proportional to $\exp(-8\pi \times 137 \times \sin^2 \theta_W)$, is very long. Another possibility involving black holes has been discussed by Zel'dovich (1976) and by Hawking, Page and Pope (1979, 1980). The "no-hair" theorem asserts that the only conserved quantities that can be associated with black holes are those corresponding to long-range fields: mass, angular momentum and electromagnetic charge. Therefore, black holes can in principle catalyze proton decay by

$$p + BH(m, J, Q) \rightarrow BH(m', J + \frac{1}{2}, Q + 1) \rightarrow BH(m, J, Q) + e^+ \quad (47)$$

For black holes of astronomical size, such a process would be exorbitantly slow. However, at distances shorter than the Planck length, protons may be surrounded by virtual black holes with masses of the order of the Planck mass, which could trigger catalyzed proton decay with a lifetime

$$\tau_p \approx (10^{45} \text{ to } 10^{50}) \text{ years} \quad (48)$$

There also exist other, even more speculative possibilities. Because even in GUTs proton decay originates at length scales $M_X^{-1} \sim O(10^{-28})$ cm that have never been probed before, one might ask whether the laws of physics as we know them remain valid at such small distances. For instance, Nielsen and Ninomiya (1978) have conjectured that Lorentz invariance may break down at small distances by Lorentz non-invariant interactions of gauge fields. However, it has been argued that within this framework the Lorentz invariance at scales 10^{-16} cm is

sufficient to guarantee Lorentz invariance also up to scales 10^{-28} cm. If one were willing to sacrifice also the rest of the Poincaré invariance, the translation invariance, which could break down because of granularity of space-time at small distances, this would show up as energy and momentum non-conservation (Ellis, Gaillard, Nanopoulos and Rudaz 1980). Because nucleon decay probes scales $\lambda \lesssim 10^{-28}$ cm where such unconventional phenomena could already be effective, proton decay might be the place to look for events with no energy-momentum balance. However, as many detection techniques rely on such a balance, it is doubtful if energy and/or momentum non-conservation could be observed in the present proton decay experiments.

At very small scales, even quantum mechanics could break down because of fluctuations in the topology of space-time, as has been suggested by Hawking (1982). This, or other possible modifications of quantum mechanics (Ellis, Hagelin, Nanopoulos and Srednicki 1984), can imply also modifications of conservation laws (Banks, Susskind and Peskin 1984; Hawking 1984; Parisi 1984; Gross 1984). However, by now it should be clear that such wild speculations are a Pandora's Box that is better left unopened for the time being. Therefore, for the rest of this review, we will concentrate on baryon decay in GUTs.

4. - PROTON DECAY AND GRAND UNIFIED THEORIES

4.1 Proton decay in ordinary GUTs

As we pointed out in Section 3.1, the proton lifetime in ordinary GUTs is expected on dimensional grounds only to be

$$\tau_p = C M_X^4 / m_p^5 \quad (49)$$

where C is a constant that we are going to calculate in this Section. The lifetime (49) is evidently extremely sensitive to the grand unified mass M_X , and therefore, its exact evaluation is very important in order for the theoretical prediction to be compared with the current lower bound (for a review, see Koshiha 1984):

$$\tau(p \rightarrow e\pi^0) > 2 \times 10^{32} \text{ years} \quad (50)$$

This is a very relevant measure for the minimal SU(5), where the dominant nucleon decay modes are (Machacek 1979)

$$p \rightarrow e^+ \pi^0 \text{ (BR} \approx 30\%) ; n \rightarrow e^+ \pi^- \text{ (BR} \approx 60\%) \quad (51)$$

In principle, one can go back to Eq. (12), use the experimentally measured value of $\sin^2 \theta_W$, Eq. (4), and $\alpha(M_W)$ to calculate M_X . The errors inherent in the experimental value of $\sin^2 \theta_W$ will then be transmitted to uncertainties in M_X with

$$\Delta \sin^2 \theta_W = \frac{3}{8} \left(-\frac{1}{128} \right) \frac{1}{4\pi} \frac{110 - N_H}{9} \Delta \ln M_X^2 \quad (52)$$

If we take the present error in the determination of $\sin^2 \theta_W$ to be ± 0.010 , we find from (52) that this would correspond to a change in M_X by a factor of $\approx 10^{\pm 1}$, a relative uncertainty of 10^4 in the proton lifetime! More reliable estimates can be obtained by starting from Eq. (10) and using the "known" strength of the strong interactions at low energies to obtain the value $M_X = (1.5 \pm 0.5) \times 10^{15} \Lambda_{\overline{MS}}$ of Eq. (11). The experimental determination of $\Lambda_{\overline{MS}}$ has proved to be rather difficult, but fortunately lattice QCD studies have advanced to the point where one is able to derive a more precise value of $\Lambda_{\overline{MS}}$. In the pure SU(3) gauge theory, calculations of the potential between static quark-antiquark pair yield (Barkai, Moriarty and Rebbi 1984a,b)

$$\Lambda_{\text{lattice}} = (9.74 \pm 0.10) \times 10^{-3} \sigma^{1/2} \quad (53)$$

where $\sigma^{1/2}$ is the string tension with a value $\sigma^{1/2} = 420$ MeV. Other authors (Otto and Stack 1984; Campbell, Michael and Rakow 1984) agree with the result (53). The error in (53) does not contain systematic errors which arise, e.g., from the fact that lattice calculations are performed at finite values of the bare coupling constant, and some calculations of meson masses on lattices would agree with Λ_{lattice} 20% larger than (53) (Lipps, Martinelli, Petronzio and Rapuano 1983; Billoire, Marinari and Petronzio 1984). The lattice QCD scale parameter can be related to Λ_{MOM} , the QCD Λ -parameter obtained by momentum space subtraction, by $\Lambda_{\text{lattice}} = 0.0120 \Lambda_{\text{MOM}}$ (Hasenfratz and Hasenfratz 1980; Dashen and Gross 1981), whereas $\Lambda_{\overline{MS}} = 2.16^{-1} \Lambda_{\text{MOM}}$ (with four flavours). From all this we deduce what we feel to be a reasonable estimate for $\Lambda_{\overline{MS}}$

$$\Lambda_{\overline{MS}} = (160 \pm 40) \text{ MeV} \quad (54)$$

which can be rounded off to $\Lambda_{\overline{MS}} = (100 \text{ to } 200) \text{ MeV}$. This yields the following conservative range of values of M_X in the minimal SU(5)

$$M_X = (1 \text{ to } 4) \times 10^{14} \text{ GeV} \quad (55)$$

Having thus established the value of M_X in the minimal SU(5), we now turn to the calculation of the constant C in (49). The strategy to be followed is depicted in Fig. 10 and is shortly the following: (i) write down the effective Lagrangian $\mathcal{L}_{\text{eff}}$ for proton decay; (ii) scale this down to low energies where nucleon decay takes place and (iii) evaluate the hadronic matrix element $\langle 0 | \mathcal{L}_{\text{eff}} | p \rangle$. Of these, step (i) depends on the detail of the GUT in question and has for SU(5) already been done in Section 3.1. Step (ii) depends on the details of the low energy effective theory, which will scale the constant C in (49) through radiative corrections, which are well known. The most important ones are naturally the QCD corrections. Finally, step (iii) is the most uncertain one as it involves "dirty" hadronic physics.

The scaling of $\mathcal{L}_{\text{eff}}$ (30) down to energies $\sim 1 \text{ GeV}$ gives an enhancement factor

$$\mathcal{L}_{\text{eff}}(1 \text{ GeV}) = A \mathcal{L}_{\text{eff}} \quad (56)$$

The enhancement factor A can be calculated using the renormalization group and the anomalous dimension of the $qqql$ operators. To leading order, one finds the enhancement factor due to gluon exchange to be (Buras, Ellis, Gaillard and Nanopoulos 1978)

$$A_3 = \left[\frac{\alpha_3(1 \text{ GeV})}{\alpha_{\text{GUT}}} \right]^{\frac{2}{11 - \frac{4}{3}N_g}} \quad (57a)$$

whereas W^+ , Z^0 and γ exchange yields a factor (Ellis, Gaillard and Nanopoulos 1980)

$$A_2 = \left[\frac{\alpha_2(M_W)}{\alpha_{GUT}} \right]^{\frac{27}{86-16N_g}} \quad (57b)$$

$$A_1 = \left[\frac{\alpha_1(M_W)}{\alpha_{GUT}} \right]^{\frac{69}{6+80N_g}} \quad \text{for } \bar{e}_L^+, \bar{\mu}_L^+ \text{ operators} \quad (57c)$$

$$= \left[\frac{\alpha_1(M_W)}{\alpha_{GUT}} \right]^{\frac{33}{6+80N_g}} \quad \text{for } \bar{e}_R^+, \bar{\mu}_R^+, \nu \text{ operators} \quad (57d)$$

where (57c) and (57d) refer to the different operators appearing in (30). For typical values of the various gauge coupling constants in Eqs. (57), one finds $A_2^2 \approx 5$ and $(A_2 A_1)^2 \approx 2.5$ [for (57c)] or ≈ 2.2 [for (57d)]. The effect of these enhancement factors is to decrease the naive decay rate as calculated from (30) by a factor of 20%.

The most uncertain part of the computation is now to evaluate the hadronic matrix elements of the effective Lagrangian between the proton initial state and the various meson final states. Here several methods have been applied. (For a nice review, see Ellis 1984b). One can assume that the dominant contribution to proton decay comes from the quark-quark annihilation diagram of the type depicted in Fig. 11. The probability of the annihilation is believed to be determined by the overlap of non-relativistic SU(6) wave functions ψ . The value of $|\psi(0)|^2$ can be extracted from Ω and S-wave hyperon decays. Another method is to use a bag model to compute the initial qq overlap in a nucleon bag and then to compute the overlap of the resulting $q\bar{q}$ system with different exclusive mesonic final state bags. These estimates were widely different and led to

$$\tau_p = (0.25 \text{ to } 10) \times 10^{30} \times \left(\frac{M_X}{4 \times 10^{14} \text{ GeV}} \right)^4 \text{ years} \quad (58)$$

More sophisticated estimates are now also available. Brodsky, Ellis, Hagelin and Sachrajda (1984) have used the current algebra and PCAC connection between the baryon decay matrix element and three-quark annihilation matrix elements

$$\langle \bar{l}, \text{pseudoscalar meson } P | \mathcal{L}_{\text{eff}}(\mu) | B \rangle \approx f_B^{-1} \langle \bar{l} | [Q_5^P, \mathcal{L}_{\text{eff}}] | B \rangle + \quad (59)$$

+ pole term

where Q_5^P is the axial charge corresponding to the pseudoscalar meson P , and the pole term is proportional to $\langle \bar{l} | \mathcal{L}_{\text{eff}}(\mu) | B \rangle$. Because the lepton is structureless, current algebra and PCAC can be used to fix both the S and P wave baryon decay amplitudes. Both terms in (59) are proportional to

$$\langle 0 | (d_{iR} u_{jR}) \epsilon_{ijk} u_{kL} | p \rangle \equiv \alpha \quad (60)$$

where α is proportional to the three-quark overlap at a point. Brodsky, Ellis, Hagelin and Sachrajda extract this from the measured baryon form factors at large Q^2 and from $J/\psi \rightarrow p\bar{p}$ decay. These two pieces of data can be fitted by a phenomenological light-cone wave function

$$\psi_{3q} = B \exp \left[-b^2 \sum_{i=1}^3 \left(\frac{k_{i\perp}^2 + m_i^2}{x_i} \right) \right] \chi \quad (61)$$

where $k_{i\perp}$ are quark transverse momenta and

$$\chi = \frac{1}{\sqrt{3}} \left[\frac{u(2)}{\sqrt{6}} (d_{\uparrow}(1) u_{\uparrow}(3) + u_{\uparrow}(1) d_{\uparrow}(3)) - \frac{\sqrt{2}}{3} u_{\uparrow}(1) d_{\downarrow}(2) u_{\uparrow}(3) \right] \quad (62)$$

is SU(6)-style flavour-spin dependence in the infinite momentum frame. Here B and b are parameters which can be fitted to data to yield $B = 0.35 \text{ GeV}^4$ and $b^2 m_q^2 = 0.012$, and m_i are quark masses., With $m_q = 300 \text{ MeV}$

$$\alpha = \frac{4m_q^2}{m_B} \pi \int d^2 k_{i\perp} dx_i \psi \frac{1}{x_1 x_2} \quad (63)$$

is fitted to be $\alpha = 0.03 \text{ GeV}^2$. Perturbative renormalization of the wave function is then used to scale the values down from $Q^2 \approx 20 \text{ GeV}^2$ where the measurements were taken, to the 1 GeV scale, and Brodsky, Ellis, Hagelin and Sachrajda (1984) find

$$\tau(p \rightarrow e^+ \pi^0) = 1.3 \times 10^{29} \left(\frac{M_X}{4 \times 10^{14} \text{ GeV}} \right)^4 \text{ years} \quad (64)$$

in clear conflict with the IMB experimental limit (50). Therefore, on the basis of experimental measurements on the proton lifetime we can assert that we have already obtained a definite result, although a negative one, about grand unification: the simple SU(5) no longer appears to be a feasible candidate for a GUT. Of course, non-minimal extensions of SU(5) with additional Higgses or models based on many other gauge groups such as SO(10) or E_6 (for a review, see Langacker 1981) are not yet ruled out. Another intriguing possibility is that there are non-renormalizable interactions which modify the gauge boson kinetic terms (Hill 1984; Shafi and Wetterich 1984). In such cases, proper normalization of the gauge boson kinetic terms will induce changes in the naive gauge coupling relations at M_X . One may envisage that they are non-renormalizable, e.g., of the form $\epsilon (\phi/M_P) F_{\mu\nu}^a F^{a\mu\nu}$, which after ϕ obtain a VEV, necessitate point transformations in the gauge fields to obtain canonical kinetic energy and at the same time modify the relations between coupling constants. One finds that in SU(5)

$$\begin{aligned} \sin^2 \theta_W(M_X) &= (1 - A\epsilon) \sin^2 \theta_W(M_X) \Big|_{\text{canonical}} \\ \ln \frac{M_X}{M_W} &= (1 + B\epsilon) \ln \left(\frac{M_X}{M_W} \right) \Big|_{\text{canonical}} \end{aligned} \quad (65)$$

where A and B are positive, calculable constants and the canonical values refer to those calculated with canonical kinetic terms. Note that there is a tendency for M_X to increase, and indeed there is a range of values for ϵ for which τ_p can be increased without a catastrophic decrease in $\sin^2 \theta_W$.

4.2 Proton decay in SUSY GUTs

In SUSY GUTs the proton decay is very much different from that of the ordinary GUTs. This is due to the fact that besides ordinary dimension six operators, there also appear dimension five baryon number violating operators which

can compete in mediating proton decay. The superfield form of the latter was already given in Section 3.1. The superspace integration picks out an F-term, and in component form this operator reads

$$\mathcal{L}_{eff}^{(5)} = \frac{\lambda_Y^2}{M_{H_3}} \left[\tilde{q}\tilde{q}q\ell \text{ or } \tilde{q}qq\tilde{\ell} \right] \quad (66)$$

where $\tilde{q}$ and $\tilde{\ell}$ denote scalar quarks and leptons, respectively. Such an operator can be generated by exchanges of colour triplet Higgses and their scalar counterparts as shown in Fig. 12. To obtain an operator capable of mediating proton decay, one must transform the two squarks (a squark and an antilepton) into two quarks (a quark and an antilepton), thus yielding a $d = 6$ $qqql$ operator. This happens by dressing the operator (66) with gaugino exchange as depicted in Fig. 13, and the leading contribution can be shown to arise from exchanges of charged supersymmetric fermions. In general, one expects that the resulting $d = 6$ effective Lagrangian is of the order of (Ellis, Nanopoulos and Rudaz 1982)

$$\mathcal{L}_{eff} \approx \frac{\lambda_Y^2}{M_{H_3}} \left(\frac{g_2^2}{m_{\tilde{W}}^2} \text{ or } \frac{g_1^2}{m_{\tilde{B}}^2} \right) (qqql) \quad (67)$$

where $m_{\tilde{W}}$ and $m_{\tilde{B}}$ are gaugino masses that must be at most $O(M_W)$ because of the stability of the gauge hierarchy. Therefore, the dimensional coupling constant of (67) is much larger than the usual $g_{GUT}^2/2M_X^2$, despite the fact that there is an additional suppression by Yukawa couplings.

Charged SUSY fermions mix in general through the mass matrix of the winos $\tilde{W}^{\pm}$ and the charged Higgsinos $\tilde{H}^{\pm}$:

$$(\tilde{W}^+, \tilde{H}^+)_L \begin{pmatrix} M_2 & g_2 v_2 \\ g_2 v_1 & -\epsilon \end{pmatrix} \begin{pmatrix} \tilde{W}^- \\ \tilde{H}^- \end{pmatrix}_L \quad (68)$$

where M_2 is the SUSY violating SU(2) gaugino mass, $v_{1,2} = \langle \bar{H} \rangle^0$ with $M_W^2 = \frac{1}{2} \times g_2^2 (v_1^2 + v_2^2)$. Let the mass eigenstates of (68) be m_1 and m_2 and let the mixing angle be θ . The effective Lagrangian for nucleon decay arising from the $d = 5$ interaction (66) can (Campbell, Ellis and Nanopoulos 1984) be evaluated in the minimal (two-pentaplet) SUSY SU(5) to be

$$\mathcal{L}_{eff} = - \frac{g_2^2}{16\pi^2} \frac{A}{M_{H_3}} \frac{v_1^2 + v_2^2}{v_1 v_2} K(M_W, m_{\tilde{q}}, m_{\tilde{e}}, m_1, m_2, \theta) O_W \quad (69)$$

where K is a kinematic function of the various mass parameters, and

$$\begin{aligned} \frac{1}{2} O_W = & m_c m_s \cos \theta_c \sin^2 \theta_c [(us)_L (d\nu_\mu)_L + (ud)_L (s\nu_\mu)_L] \\ & + m_c m_s \sin^3 \theta_c [(ud)_L (d\nu_\mu)_L] + m_u m_s \cos \theta_c \times \\ & \times [(su)_L (u\mu - d\nu_\mu \cos \theta_c)_L \cos \theta_c + 2 \cos^2 \theta_c (du)_L (s\nu_\mu)_L] \end{aligned} \quad (70)$$

K is rather sensitive to the mass spectrum of SUSY particles. Small values of K can be obtained if $m_{\tilde{q}} \gg M_W$ and/or $|M_2|, |\epsilon| \ll M_W$. If $m_{\tilde{q}} \approx M_W$, one finds

$$|K| \approx |(M_2 + 2\epsilon) / 3M_W| \quad (71)$$

which has the interesting feature that partial or complete cancellations can occur if M_2 and ϵ have opposite signs.

It should be pointed out that there is an alternative form of SUSY $d = 5$ operator which cannot be simply obtained from superheavy triplet Higgs(ino) exchange, but could be generated at the Planck scale gravito-dynamics (Ellis, Hagelin, Nanopoulos and Tamvakis 1983). One assumes that although particle physics below the Planck scale is described by a renormalizable broken SUSY theory, there are non-renormalizable interactions scaled by the appropriate inverse powers of the Planck mass. To lowest order in M_P^{-1} , such an interaction in the SU(5) superpotential that is relevant to baryon number violation is

$$f_{eff} = \frac{1}{M_P} \bar{F}_a^\alpha T_{\alpha\beta}^b T_{\delta\gamma}^c T_{\lambda\mu}^d \epsilon^{\beta\gamma\delta\lambda\mu} \lambda_{bcd}^a \quad (72)$$

where the superfields transform as $\bar{F} \sim \bar{5}$ and $T \sim 10$, and latin indices denote different generations. This gives rise to a four-fermion operator proportional to

$$O'_W = (su)_L (u\ell - d\nu_\ell) + 2(du)_L (s\nu_\ell)_L \quad (73)$$

The striking feature of the operators (70) and (73) is that they lead to dominant decay modes which will invariably involve strange particles. As such, it is the signature of SUSY that will hopefully be seen in the proton decay experiments. Although precise numbers will depend on the elements of the Kobayashi-Maskawa matrix, one can from (70) assert the hierarchy of the partial rates

$$\begin{aligned} \Gamma(N \rightarrow \bar{\nu}_\tau K) &>> \Gamma(N \rightarrow \bar{\nu}_\mu K) >> \Gamma(N \rightarrow \bar{\nu}\pi; \mu^+ K) >> \\ &>> \Gamma(N \rightarrow \mu^+ \pi) >> \Gamma(N \rightarrow e^+ K) >> \Gamma(N \rightarrow e^+ \pi) \end{aligned} \quad (74)$$

Note especially the low ranking of $N \rightarrow e^+ \pi$, which was the dominant channel in the minimal non-SUSY SU(5). There are some differences in the partial rates for, e.g., $p \rightarrow \nu K^+$ and $n \rightarrow \bar{\nu} K^0$ between the operators (70) and (73), and the latter predicts a considerable rate also for $p \rightarrow \mu^+ K^0$. The relative ratios for both types of $d = 5$ SUSY operators have been presented in Table 3.

To compute the proton lifetime due to the $d = 5$ operators, one should again find the appropriate enhancement factors of the operators (70) and (73) and evaluate the matrix element of the scaled effective Lagrangian between hadronic states. In the case of SUSY GUTs, the renormalization group scaling of the effective Lagrangian produces a net suppression factor, in contrast with the enhancement found in the case of ordinary GUTs. For evaluation of the hadronic matrix elements, one can use techniques based on current algebra and PCAC as outlined in Section 4.1. There is some uncertainty in the value of the three-quark overlap α , Eq. (60), as well as in the value of M_{H_3} , which is rather model dependent. The nucleon decay rate depends on the ratio α/M_{H_3} . Parametrizing this with

$$R^2 = \left(\frac{M_{H_3}}{\alpha} \right)^2 \left(\frac{0.03 \text{ GeV}^3}{10^{16} \text{ GeV}} \right)^2 \quad (75)$$

Campbell, Ellis and Nanopoulos (1984) find the lifetime for $B \rightarrow \bar{\nu} K$ in the minimal SU(5) SUSY GUT to be

$$\tau(B \rightarrow \nu K) = 3.6 \times 10^{26} \text{ yr} \left[\frac{2v_1 v_2}{v_1^2 + v_2^2} \right]^2 \frac{R^2}{K^2} \quad (76)$$

With $M_{H_3} = M_X = 10^{16}$ GeV [see (24)] and with the value of α that was used in Section 4.1 $R^2 = 1$. Taking for simplicity $v_1 = v_2$, a choice that is actually preferred by some currently popular supergravity models (Ellis, Hagelin, Nanopoulos and Tamvakis 1983b; Kounnas, Lahanas, Nanopoulos and Quiros 1983, 1984), as well as interpretations of the UA1 monojet data in terms of supersymmetry (Ellis and Sher 1984; Enqvist, Lahanas and Nanopoulos 1984), the experimental limit (see, e.g., Koshihara 1984)

$$\tau(n \rightarrow \bar{\nu} K^0) > 8 \times 10^{31} \text{ years} \quad (77)$$

requires that $|K| < 2 \times 10^{-3} |R|$. In Fig. 14, we show the experimentally allowed domains in the (M_2, ϵ) space that are consistent with the bound (77) and other constraints for the interesting (Ellis and Kowalski 1984a,b; Reya and Roy 1984a,b) $m_{\tilde{q}} = m_{\tilde{\chi}} = \frac{1}{2} M_W$. These domains would change substantially if M_{H_3} were significantly larger than about 10^{16} GeV assumed in Fig. 14. Therefore, the minimal SUSY SU(5) appears to be consistent with experiments.

Let us now turn to $d = 6$ SUSY operators. These have the following general form

$$\int d^2\theta d^2\bar{\theta} \phi^\dagger \phi \phi^\dagger \phi \quad (78)$$

and can arise through gauge boson exchange or chiral superfield exchange (see Fig. 15). These may be of importance for the proton decay if $d = 5$ operators are forbidden by some selection rules. This is the case in the scenarios which contain "light" Higgs triplets with $M_{H_3} \approx 10^{10}$ GeV (Nanopoulos and Tamvakis 1982a; Srednicki 1982; Nanopoulos, Olive and Tamvakis 1982; Nanopoulos, Olive and Srednicki 1983). Such "light" Higgses have in the past been considered desirable for cosmological reasons. They have been especially needed (Nanopoulos and Tamvakis 1982a) for baryon asymmetry generation after a primordial period of cosmological inflation, whereby the scale factor of the Universe expands by about 60 orders of magnitude, thus giving an explanation for the observed isotropy, homogeneity and flatness of the Universe. After inflation the Universe reheats through oscillations of the inflaton field up to temperatures of the order of

10^{11} GeV, a number essentially fixed by the value of energy density perturbations produced by inflation. Therefore, in this picture one would need "light" Higgs triplets to produce the baryon asymmetry (this is not a general feature, though, see Enqvist and Nanopoulos 1984a,b) and one must then forbid the appearance of the $d = 5$ operators which could otherwise cause catastrophically rapid nucleon decay.

In theories where only two Higgs multiplets contribute significantly to the quark and lepton mass matrices, the dominant $d = 6$ operators are:

$$\begin{aligned} \mathcal{L}_{eff} = & \frac{1}{V_1^2 M_{H_3}^2} \left[\sin^2 \theta_C m_S^2 (su)_R (u\mu)_L - m_d m_S (du)_R (s\nu_\mu)_L \right. \\ & \left. + \dots \right] + \frac{1}{V_2^2 M_{H_3}^2} \left[m_u^2 (ud)_L (ue)_R + \dots \right] \end{aligned} \quad (79)$$

The first two terms in Eq. (79) give essentially equal amplitudes, which result in the relative rates for $p \rightarrow \bar{\nu} K^+$, $n \rightarrow \bar{\nu} K^0$ and $p \rightarrow \mu^+ K^0$ shown in Table 3. We note that the relative rates to be expected are characteristically different from those due to Higgsino or other SUSY $d = 5$ operators. In addition, the $p \rightarrow e^+ \pi^0$ decay rate is given by

$$\frac{\Gamma(p \rightarrow e^+ \pi^0)}{\Gamma(n \rightarrow \bar{\nu}_\mu K^0)} = \frac{1}{2} \left(\frac{V_1}{V_2} \right)^4 \frac{m_u^2}{m_d m_S} \left(\frac{1+D+F}{1+\frac{1}{3}D+F} \right)^2 \quad (80)$$

where D and F are $SU(3)$ breaking factors. Therefore the best bounds on M_{H_3} probably come (almost equivalently) from the experimental lower limits on $n \rightarrow \bar{\nu} K^0$ and on $p \rightarrow \mu^+ K^0$:

$$M_{H_3} \left(g_2 V_1 / M_W \right) \geq 2.3 \times 10^{11} \text{ GeV} \quad (81)$$

This is within the range of values of M_{H_3} , which are suggested by cosmological considerations. The enhancement factors for $d = 6$ operators have also been calculated (Ibáñez and Muñoz 1984) and are practically identical to the corresponding non-supersymmetric enhancement factors. However, without "light" Higgses, the SUSY proton decay rate due to $d = 6$ operators alone would probably be beyond the present experimental capabilities. This is because $M_X \approx 10^{16}$ GeV in the SUSY $SU(5)$ [see (25)] giving an enhancement factor 10^8 with respect to the proton lifetime in ordinary $SU(5)$. Forbidding the appearance of $d = 5$ operators seems to be somewhat unnatural, though, and as recent studies suggest, is not

even demanded by cosmological consistency (Kim, Masiero and Nanopoulos 1984; Ellis et al. 1984). Therefore, SUSY GUT-induced proton decay, if supersymmetry is indeed playing the role usually allotted to it, should be near its experimental discovery with the characteristic and unmistakeable signature $p \rightarrow \text{strange particles}$.

5. - PROTON DECAY EXPERIMENTS

At the moment, there are seven active nucleon decay experiments. (For recent experimental reviews, see Perkins 1984; Litchfield 1984). They use two different techniques for proton decay searches. Three of the experiments use water Cerenkov detectors, which measure the Cerenkov light emitted by relativistic charged particles using photomultipliers. Such a detector was proposed largely because of the dominant $p \rightarrow e^+\pi^0$ mode of the minimal SU(5). The advantages of water Cerenkov detectors are that the Cerenkov light has good directionality, so that two-body decays leaving two-track back-to-back events are easy to recognize. However, they have some difficulties with high multiplicity tracks and are not efficient for detecting charged mesons.

Four of the presently running experiments use iron tracking calorimeters. There are sensitive to all charged particles and can achieve a fine resolution. Instrumenting a calorimeter is difficult, however, and one has to rely on sampling tracks between iron sheets.

The IMB detector is a water Cerenkov detector mounted in the Morton Salt Mine, Ohio (Bionta et al. 1982). It has 3000 tons of fiducial mass and has yielded a total exposure of 1800 ton years. The Kamioka experiment (Arisake et al. 1982) has a 3000 ton cylinder of water located in the Kamioka mine, Japan, with a fiducial mass of about 900 tons. It has reported a total exposure of 660 ton years. The HPW collaboration (Gaidos et al. 1982) has its water detector in a mine in Park City, Utah. A 900 ton cylinder of water had produced an exposure of 500 ton years.

The first iron calorimeter (KGF) is located in Kolar Gold Field, India (Krishnaswamy et al. 1981, 1982). It consists of a sandwich of 1.2 cm iron plates with proportional counters, and has been running for three years giving an exposure of 180 ton years. The NUSEX experiment (Battistoni et al. 1982) uses calorimeter of approximately the same size as the KGF experiment and is being carried out in the Mont Blanc tunnel. It has yielded an exposure of 220 ton

years. The FREJUS detector (Barloutaud et al. 1982) is located in a tunnel between France and Italy and has started operating only recently. The SOUDAN 2 detector is a large, highly instrumented calorimeter that will be installed in the Soudan iron mine in Minnesota. Table 4 shows recent limits on the proton lifetime for different decay modes obtained in some of the above experiments.

The dominant background in proton decay experiments is due to atmospheric neutrinos produced by cascading hadrons in the atmosphere. Their interactions in the detector have characteristic topologies which can be used to differentiate them from proton decay candidates. The flux of such neutrinos has to be evaluated, e.g., by Monte Carlo calculations. However, recently Battistoni et. al (1983) exposed a module similar to that in the calorimeter they employed to search for proton decay, to an accelerator beam from the CERN PS, and can therefore compare directly underground events with neutrino events. (For a detailed description of the background situation, see Perkins 1984).

As can be seen from Table 4, the IMB result excludes the simple SU(5), as was already discussed in previous sections. At the moment, there is no convincing evidence for nucleon instability. The best candidates for proton decay are the $\mu\eta$ event of KAMIOKA and the μK^0 candidate of NUSEX. These are apparently difficult to understand in terms of background which, recalling that in SUSY, one expects $p \rightarrow$ strange to be the dominant mode (see Section 4.2), is perhaps suggestive for those of us who have faith in SUSY. However, we will probably have to wait for at least two years before anything definite about these channels can be said. Meanwhile, with time the IMB experiment should be able to push the limit on $p \rightarrow e^+\pi^0$ to around 10^{33} years.

6. - CONCLUSIONS

Do protons decay ? That is the question ! Actually, it is one of the hottest questions in particle physics today. We have tried to show in this review that there are many good reasons that indicate an affirmative answer to this basic question. There are so many things that fit together if, as seems rather inevitable to us, protons are not for ever. Grand unification is such a natural extension of the Standard Model that avoiding it appears to be pretty hard. Then, it would be almost inconceivable that protons are stable particles. What is of great importance, of course, is the spectrum of the decay products. It seems that the minimal SU(5) model without supersymmetry is, when compared

with experiment, in big trouble. On the other hand, supersymmetric GUTs with an observable proton lifetime and mostly kaons in the final state are not in bad shape when contrasted with preliminary experimental results. A positive signal for proton decay in kaons will be for a lot of us a superhit. Not only the ideas of grand unification will be vindicated but also, albeit indirectly, supersymmetry will be in pretty good shape. It should be stressed once more that deciated proton decay experiments are very difficult and patience with their output is required. Nevertheless, it is not easy to overestimate their importance and we are all eagerly waiting for new experimental results. It is interesting to notice that the baryon number violating interactions responsible for our existence ($n_B/n_\gamma \sim 10^{-10}$) are driving us also to extinction (matter instability).

IF THE PROTON WERE STABLE IT WOULD NOT EXIST

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Table 1

The generation structure of the observed fermions

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

Table 2

Supersymmetric particles

$p \rightarrow \bar{\nu} K^+$	<p>$d = 5$ $\tilde{H}_3, H_3$ exchange</p>	<p>Alternative $d = 5$ SUSY operator</p>	<p>$d = 6$ H_3 exchange</p>	0.55	1.75	1
$n \rightarrow \bar{\nu} K^0$				1	1	1
$p \rightarrow \mu^+ K^0$				$(10^{-3} - 10^{-4})$	0.10	0.55

Table 3
Relative ratios for baryon decays into strange particles

Experiment	KGF	NUSEX	KAMIOKA	IMB
$p \rightarrow e^+\pi^0$	~ 2 (1)	1.5	2.0	12
$\rightarrow \mu^+\pi^0$	~ 2 (1)	1.0	2.0	5.1 (1)
$\rightarrow e^+\rho^0$	-	-	2.5	2.5 (1)
$\rightarrow \mu^+K^0$	~ 2 (1)	0.6 (1)	0.8 (1)	2.9 (2)
$\rightarrow \mu^+\eta^0$	-	-	0.4 (1)	3.1 (2)
$\rightarrow e^+K^0$	-	-	0.9	4.9
$\rightarrow \nu K^+$	$\sim 2(1)$	0.2	0.5 (2)	0.7 (3)
$\rightarrow \nu \pi^+$		0.2 (< 3)	0.3 (2)	-
$n \rightarrow e^+\pi^-$	~ 2 (1)	1.5	1.2	2.5 (4)
$\rightarrow \nu \pi^0$	-	0.7	0.6	0.7 (28)
$\rightarrow \nu K^0$	-	0.5	0.3	1 (2)

Table 4

Recent limits on proton lifetime

In units of 10^{31} yrs $\times$ branching ratio. 90% CL lower limits, unless $\sim$ indicates signal claimed. Number of candidates, if any, given in the brackets. (See Perkins 1984 for references; the IMB column has been updated according to H.S. Park et al. 1985).

FIGURE CAPTIONS

- Fig. 1 : A schematic view of the running of the gauge coupling constants and their unification.
- Fig. 2 : Radiative corrections proportional to $\Lambda_{\text{cut-off}}^2$ to the Higgs mass.
- Fig. 3 : Cancellation of divergencies in SUSY.
- Fig. 4 : a) Neutral gaugino and b) charged gaugino SUSY box diagrams contributing to the K_1 - K_2 mass difference.
- Fig. 5 : Proton decay in SU(5) due to exchanges of superheavy gauge bosons.
- Fig. 6 : Impressions of a) Higgs fields in causally separated domains pointing in different internal group directions causing b) a monopole to be formed where they meet. Note that the value of the Higgs field goes to zero in the centre.
- Fig. 7 : The dominant decay modes of superheavy vector bosons X and Higgs colour triplets H_3 .
- Fig. 8 : An example of a 4th order interference diagram contributing to ΔB .
- Fig. 9 : a) The supergraph responsible for baryogenesis in SUSY GUTs;
b) A component diagram of the supergraph a) showing explicitly the mass insertions on the lines.
- Fig. 10 : The strategy for the evaluation of the proton lifetime.
- Fig. 11 : A schematic picture of nucleon decay.
- Fig. 12 : An effective $d = 5$ baryon number violating operator in SUSY GUTs.
- Fig. 13 : Dressing up the $d = 5$ operator to obtain a $d = 6$ operator capable of mediating proton decay.

Fig. 14 : Allowed domains in (M_2, ϵ) space consistent with cosmology ($- \dots -$), the absence of charged particles lighter than 20 GeV ($- . -$) and the absence of baryon decay with $R = 1$ ($-$), $R = 10$ ($---$) and $R = 100$ ($\dots$) [see (75)]: (a) $v_1 = v_2$, $\epsilon > 0$; (b) $v_1 = v_2$, $\epsilon < 0$; (c) $v_1 = 4v_2$, $\epsilon > 0$; (d) $v_1 = 4v_2$, $\epsilon < 0$; all taking $m_q \approx m_l = \frac{1}{2}M_W$.

Fig. 15 : Supergraphs for $d = 6$ proton decay; a) gauge superfield exchange and b) chiral superfield exchange.

Fig. 3

Fig. 4

Fig. 5

a)

b)

Fig. 6

Fig. 7

Fig. 8

Fig. 9

Fig. 10

Fig. 11

Fig. 12

Fig. 13

Fig. 14

Fig. 15