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PROTON DECAY:
THEORY VERSUS EXPERIMENT

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Contents

Abstract

This review scrutinizes the present status of proton decay in grand unified theories. Baryon and lepton number violation in conventional as well as supersymmetric GUTs is discussed with special emphasis being laid on selection rules and model-independent predictions. The theoretical predictions for nucleon lifetimes and branching ratios, when confronted with experiment, inevitably lead to the conclusion that all great desert GUTs, like the minimal SU(5) model, are definitely ruled out by the experimental non-confirmation of proton decay at the expected rate.

I. Introduction

The most fascinating aspect of grand unification is the fact that it provides a natural framework for baryon and lepton number non-conservation [1.1-1.5]. Grand unified theories (GUTs) predict in general the existence of phenomena like neutrino oscillations, neutron oscillations, and - most spectacular - nucleon instability. In particular, baryon number is no more an absolutely conserved quantity. Its conservation is violated either explicitly by gauge interactions and Yukawa couplings or non-perturbatively by magnetic monopoles. Consequently, the experimental observation of proton decay would be a clear indication for grand unification. Moreover, proton decay might well prove to be the only experimental meansto test the concept of grand unified theories.

On the other hand, there is no particular reason why baryon number should be conserved. If baryon number were the conserved charge of an unbroken gauge symmetry there should be a corresponding massless gauge boson. The Eötvös-Dicke experiment, however, finds no evidence for a long-range force coupling to baryon number.

The intention of the present work is to illustrate the basic concepts of grand unified theories and to examine critically their consequences for proton decay in view of the results of the on-going nucleon decay experiments. Accordingly, I will stress those ideas and aspects which will survive anyway, like selection rules for baryon number violation arising already from low-energy gauge invariance and (eventually) supersymmetry. In contrast to that, I ignore poorly justified assumptions, predictions valid only for a certain narrow range of parameters, and things like that.

This review is organized as follows: Section II introduces the most important features of grand unified theories and presents the B violating sector of conventional GUTs as well as its immediate implications. Section III recalls the primary reasons for believing in supersymmetry and gives the outlines of supersymmetric baryon number violation. Section IV represents a survey of the theoretical expectations for proton (and bound neutron) lifetime, favoured decay modes, and branching ratios, including an account of the attempts to bring theory and experiment into accordance. Section V is a compilation of the upper bounds on the partial decay widths reported by the currently operational nucleon decay experiments. Finally, Section VI contains a brief comment on monopole catalyzed proton decay as well as some concluding remarks.

II. Baryon Number Violation in Conventional GUTs

Nowadays, it is a common belief that the strong and electroweak interactions are described by the standard model, a gauge theory based on the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$. However, although the Glashow-Salam-Weinberg model $SU(2)_L \times U(1)_Y$ correlates the weak and electromagnetic interactions it does not represent a unification of these interactions at all. Its gauge group is a direct product of two factors, consequently there are still two distinct gauge coupling constants having nothing to do with each other. Nevertheless, the Glashow-Salam-Weinberg model is sometimes misleadingly called a unified theory of weak and electromagnetic interactions - a circumstance which forced us to entitle the true unified theories as "grand unified theories".

Grand unified theories attempt a unification of the non-gravitational interactions by embedding $SU(3)_C \times SU(2)_L \times U(1)_Y$ into a simple gauge group [2.1]. Another, by far less elegant possibility would be the embedding into a direct product of isomorphic simple factors related by a discrete symmetry in an irreducible manner [2.2]. In either case one ends up with a more fundamental theory with only a single gauge coupling constant. In order to reproduce the known interactions, the GUT gauge symmetry is assumed to break spontaneously down to the standard model at a GUT mass scale m_X far above the electroweak mass scale m_W,

$$\begin{aligned} \text{GUT} &\xrightarrow[m_X]{} SU(3)_C \times SU(2)_L \times U(1)_Y \\ &\xrightarrow[m_W]{} SU(3)_C \times U(1)_Q \ . \end{aligned} \tag{2.1}$$

Of course, the GUT symmetry breakdown might also take place in several steps, i.e. at several mass scales. In this spirit the very different looking standard-model interactions are only the low-energy manifestation of the basic grand unified theory. Charge quantization already follows from

$$\mathrm{Tr}\,[Q] = 0 , \qquad (2.2)$$

since now the electric charge operator Q has to be a generator of the (semi-) simple GUT gauge group.

Any GUT gauge group coming into question has to fulfil two requirements:

(1) It has to contain the gauge group of the standard model,

$$SU(3) \times SU(2) \times U(1) \subset GUT . \qquad (2.3)$$

This entails that it has to be of at least rank 4 and, in particular, has to contain an SU(3) subgroup.

(2) It must allow for the correct reproduction of the particle content of the observed fermion spectrum. This implies that

(i) it must possess complex representations;

(ii) the representation taken into consideration for the known fermions must decompose under $SU(3)_C$ solely into singlets, triplets, and anti-triplets - and nothing else; and finally

(iii) this representation must not be plagued by Adler-Bardeen-Bell-Jackiw anomalies.

The need of a complex fermion representation arises from two different sources:

(i) The (observed) weak interactions violate parity. In other words, the known fermions of one generation form a complex representation of $SU(3)_C \times SU(2)_L \times U(1)_Y$,

$$f_L = (3,2,\tfrac{1}{6}) + (\bar{3},1,-\tfrac{2}{3}) + (\bar{3},1,\tfrac{1}{3}) + (1,2,-\tfrac{1}{2}) + (1,1,1) . \qquad (2.4)$$

(ii) The experimentally known fermions must not be trapped by the so-called survival hypothesis*: In the course of the spontaneous symmetry

*) For a more detailed discussion see e.g. Ref. [1.2].

breakdown of a group G to a subgroup H at a mass scale m_G,

$$G \xrightarrow[m_G]{} H \, , \tag{2.5}$$

all fermions transforming according to a self-conjugate representation of G will acquire, in general, a mass m_F of the order of magnitude

$$m_F = O(m_G) \, . \tag{2.6}$$

Hence, in a grand unified theory all fermions transforming according to a real representation of the GUT gauge group will become superheavy. Thus, in order to "survive" the first stage of the spontaneous symmetry breakdown (2.1),

$$GUT \xrightarrow[m_X]{} SU(3)_C \times SU(2)_L \times U(1)_Y \, , \tag{2.7}$$

as massless particles all light fermions f_L have to belong to a complex representation of the GUT gauge group.

Requirements (1) and (2) restrict possible candidates for GUTs to the gauge groups SU(n), n $\geq$ 5, SO(4n+2), n $\geq$ 2, and E_6. Mostly the gauge groups SU(5), SO(10), and E_6 have been used in the construction of viable grand unified models [1.1]. Note, that it might well happen that the chain

$$SU(5) \subset SO(10) \subset E_6 \tag{2.8}$$

is more than just a group-theoretical relation, in that it might prove also to indicate the symmetry breaking pattern of the grand unified theory realized in nature.

The grand unification mass scale m_X can be calculated by inspecting the dependence of the standard model coupling constants $\alpha_3(Q^2)$,

$$\alpha_2(Q^2) = \frac{\alpha_{e.m.}(Q^2)}{\sin^2\theta_W(Q^2)} \, , \tag{2.9}$$

and

$$\alpha_1(Q^2) = \frac{5}{3}\frac{\alpha_{e.m.}(Q^2)}{\cos^2\theta_W(Q^2)} \tag{2.10}$$

on the momentum transfer Q^2 (Fig. 1.a). At m_X they converge to the single GUT gauge coupling constant α [2.3]:

$$\alpha_3(m_X^2) = \alpha_2(m_X^2) = \alpha_1(m_X^2) = \alpha \; . \tag{2.11}$$

However, the mass scale m_X and the coupling constant α can be determined unambiguously only in so-called great desert GUTs. Great desert GUTs are characterized by the assumption that no new particle thresholds will appear between the electroweak mass scale m_W and the unification mass scale m_X. This implies that the symmetry breaking (2.7) of the GUT gauge group down to the standard model has to proceed in one single step, which excludes e.g. all left-right symmetric models $SU(2)_L \times SU(2)_R \times U(1)$ for the electroweak interactions. Hence this kind of GUTs is of particularly simple structure, showing an $SU(3)_C \times SU(2)_L \times U(1)_Y$ invariant great desert between m_W and m_X. A representative value for the GUT gauge coupling constant is

$$\alpha = 0.0244 \simeq \frac{1}{41} \; . \tag{2.12}$$

The GUT mass scale m_X turns out to be very accurately proportional to the QCD scale parameter Λ. Starting with the renormalization group extrapolation at

$$\alpha_2(m_W^2) = 0.0372 \simeq \frac{1}{27} \tag{2.13}$$

and

$$\alpha_1(m_W^2) = 0.0164 \simeq \frac{1}{61} \; , \tag{2.14}$$

and taking into account threshold effects and two-loop contributions, one obtains for three light fermion generations and one light Higgs doublet

$$m_X = 1.5\cdot 10^{15}\, \Lambda_{\overline{MS}} \; . \tag{2.15}$$

Thus the nowadays commonly accepted value of the QCD scale parameter [1.6]

$$\Lambda_{\overline{MS}} = 0.16 \begin{smallmatrix} +\ 0.10 \\ -\ 0.08 \end{smallmatrix} \text{ GeV} \tag{2.16}$$

corresponds to a value of the GUT mass of

$$m_X = 2.4 \cdot 10^{14} \text{ GeV} . \tag{2.17}$$

The constant of proportionality in Eq. (2.15) is increased by a factor of about 1.2 for each additional fermion generation while it is reduced by a factor of $(1.5)^{-1}$ for a second light Higgs doublet [1.3]. The major uncertainty in m_X, however, is introduced by the error in $\Lambda_{\overline{MS}}$, i.e. by our poor knowledge of the strong coupling strength.

Great desert GUTs, and only they, allow for a prediction of the Weinberg angle Θ_W,

$$\sin^2\Theta_W(Q^2) = \frac{\alpha_{e.m.}(Q^2)}{\alpha_2(Q^2)} = \frac{\alpha_1(Q^2)}{\alpha_1(Q^2) + \frac{5}{3}\,\alpha_2(Q^2)} . \tag{2.18}$$

Above m_X the value of $\sin^2\Theta_W$ is completely fixed by group theory, in particular by isospin I_3 and electric charge Q of the involved fermions:

$$\sin^2\Theta_W(m_X^2) = \frac{\mathrm{Tr}[I_3^2]}{\mathrm{Tr}[Q^2]} = \frac{3}{8} . \tag{2.19}$$

Renormalization effects reduce this value at the electroweak mass scale to [1.3-1.5]

$$\sin^2\hat{\Theta}_W(m_W^2) = 0.2138 \begin{smallmatrix} +\ 0.0042 \\ -\ 0.0029 \end{smallmatrix} , \tag{2.20}$$

again for the QCD parameter given in Eq. (2.16). This prediction has to be compared with the experimental determination (after including radiative corrections) [1.4]

$$\sin^2\hat{\Theta}_W(m_W^2)_{exp} = 0.219 \pm 0.006 . \tag{2.21}$$

In view of the fact that, a priori, the theoretical value of $\sin^2\theta_W$ could have emerged anywhere inbetween zero and one, it is hard to believe that this remarkable agreement of theory and experiment is purely accidental. Rather this amazing agreement seems to strongly support the great desert hypothesis, i.e. the non-existence of intermediate mass scales between m_W and m_X.

As long as there is no hint from experiment how the GUT chosen by nature looks like, the analysis of baryon number violation should be performed in a way which is entirely independent of any specific grand unified model. This model-independent analysis consists of four main steps:

(i) identification of the possible intermediate bosons,

(ii) enumeration of the resulting four-fermion operators,

(iii) formulation of the effective Lagrangian built up by these operators the coefficients of which are determined by the specific grand unified model,

(iv) renormalization of this Lagrangian from the GUT mass scale down to the relevant hadronic mass scale.

In grand unified theories proton decay is induced by bosons which are allowed by their $SU(3)_C \times SU(2)_L \times U(1)_Y$ quantum numbers to couple simultaneously to two quarks

$$qq\,, \qquad B = \frac{2}{3}\,, \qquad L = 0\,,$$

<u>and</u> to a quark and a lepton

$$\ell q\,, \qquad B = \frac{1}{3}\,, \qquad L = 1\,,$$

or to a quark and an anti-lepton,

$$\ell^c q\,, \qquad B = \frac{1}{3}\,, \qquad L = -1\,.$$

Hence no definite baryon and lepton numbers can be attributed to these bosons. Their interactions respect the $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge invariance but violate the conservation of baryon and lepton number. Considering only couplings to the known light fermions

$$q_L = \begin{pmatrix} u \\ d \end{pmatrix}_L , \quad u_R , \quad d_R , \quad \ell_L = \begin{pmatrix} \nu \\ e^- \end{pmatrix}_L , \quad e^-_R ,$$

there are just five types of bosons of this sort [2.5,2.7,2.9,2.12]:

Type	Notation	$SU(3)_C$	$SU(2)_L$	Y	Q
vector boson	X	$\bar{3}$	2	$\frac{5}{6}$	$\frac{4}{3}, \frac{1}{3}$
	X'	3	2	$\frac{1}{6}$	$\frac{2}{3}, -\frac{1}{3}$
scalar boson	H_X	3	1	$-\frac{1}{3}$	$-\frac{1}{3}$
	H'_X	3	1	$-\frac{4}{3}$	$-\frac{4}{3}$
	H''_X	3	3	$-\frac{1}{3}$	$\frac{2}{3}, -\frac{1}{3}, -\frac{4}{3}$

The experimentally observed longevity of the proton,

$$\tau_{p,exp} \gtrsim 10^{32} \text{ yr} , \tag{2.22}$$

requires

$$m_X \gtrsim O(10^{14}) \text{ GeV} \tag{2.23}$$

for the vector boson masses, which is consistent with the estimate (2.17), and - even for extremely small Yukawa coupling constants of the order $O(10^{-4})$ -

$$m_{H_X} \gtrsim O(10^{10}) \text{ GeV} \tag{2.24}$$

for the Higgs boson masses, i.e. all of these bosons have to be superheavy.

By exchange of one of the superheavy bosons listed above a tree-level interaction between three quarks and one lepton is generated. Integrating out the superheavy degrees of freedom this interaction reduces to an $SU(3)_C \times SU(2)_L \times U(1)_Y$ invariant, B and L violating dimension-six operator of the form

$$O \sim (q\, q\, q\, \ell) \; . \tag{2.25}$$

The coefficients G accompanying these four-fermion interactions in the resulting Lagrangian, however, provide a suppression by two powers of the superheavy mass,

$$G = O(\frac{1}{m_X^2}) \qquad \text{or} \qquad G = O(\frac{1}{m_{H_X}^2}) \; . \tag{2.26}$$

From the known light fermions five Lorentz invariant effective operators of this kind can be formed* [2.7,2.8,2.10,2.11]:

$$O_1 = (\overline{\ell^c{}_R}\, \tilde{q}_L)(\overline{u^c{}_L}\, d_R) \tag{2.27}$$

$$O_2 = (\overline{e^+{}_L}\, u_R)(\overline{q^c{}_R}\, \tilde{q}_L) \tag{2.28}$$

$$O_3 = (\overline{\ell^c{}_R}\, \tilde{q}_L)(\overline{q^c{}_R}\, \tilde{q}_L) \tag{2.29}$$

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*) The operator

$$O_6 = (\overline{e^+{}_L}\, d_R)(\overline{u^c{}_L}\, u_R)$$

contained in the original classifications [2.7,2.8] can be expressed in terms of O_5 with the help of the algebraic identity [2.11]

$$(\overline{\psi^c_{1L}}\, \psi_{2R})(\overline{\psi^c_{3L}}\, \psi_{4R}) + (\overline{\psi^c_{1L}}\, \psi_{3R})(\overline{\psi^c_{4L}}\, \psi_{2R}) + (\overline{\psi^c_{1L}}\, \psi_{4R})(\overline{\psi^c_{2L}}\, \psi_{3R}) = 0 \; .$$

$$O_4 = (\overline{\ell^c_R} \vec{\tau} \tilde{q}_L)(\overline{q^c_R} \vec{\tau} \tilde{q}_L) \tag{2.30}$$

$$O_5 = (\overline{e^+_L} u_R)(\overline{u^c_L} d_R) . \tag{2.31}$$

In this list $\vec{\tau}$ are the three Pauli matrices while $\tilde{q}_L$ denotes the transposed $SU(2)_L$ quark doublet

$$\tilde{q}_L := i \tau^2 q_L = \begin{pmatrix} d \\ -u \end{pmatrix}_L . \tag{2.32}$$

In order to guarantee $SU(3)_C$ invariance these operators have to be antisymmetrized with respect to the colour indices of the quark fields. The generation indices carried by all fermions also have been dropped. Fermi statistics requires more than one generation for the non-vanishing of the operator O_4.

At this point an important observation can be made [2.7,2.8]: Although the effective operators O_i $(i = 1,2,\ldots,5)$ have only been subjected to the requirement of $SU(3)_C \times SU(2)_L \times U(1)_Y$ invariance all of them show a global symmetry (maybe of local origin) represented by the generator $(B-L)^*$. These operators violate B and L separately but conserve the quantum number (B-L). Consequently, the nucleon decay controlled by these operators respects the selection rule

$$\Delta B = \Delta L . \tag{2.33}$$

Thus, nucleon decays into anti-leptons are allowed,

$$N \to \ell^c + M , \qquad \Delta B = \Delta L = -1 , \tag{2.34}$$

but nucleon decays into leptons are forbidden,

$$N \not\to \ell + M , \qquad \Delta B = -\Delta L = -1 . \tag{2.35}$$

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*) Leaving aside negligible instanton effects, $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge invariance also entails baryon number conservation by the standard-model interactions.

Here M represents any mesonic final state, $\ell = e, \mu, \nu_e, \nu_\mu, \nu_\tau$.

Furthermore, for the operators involving a strange quark the additional selection rule [2.6]

$$\frac{\Delta S}{\Delta B} \leq 0 \qquad (2.36)$$

is valid. This restriction allows the decays

$$N \to \ell^c + M(S = +1) \, , \qquad \Delta S = - \Delta B = 1 \, , \qquad (2.37)$$

but forbids the decays

$$N \to \ell^c + M(S = -1) \, , \qquad \Delta S = \Delta B = -1 \quad . \qquad (2.38)$$

All of the effective operators O_1 through O_5 can be obtained by scalar boson exchange. On the other hand, only the operators O_1 and O_2 have the correct chirality structure to be generated by vector boson exchange since - by application of the Fierz transformation

$$(\overline{\psi_{1L}} \gamma_\mu \psi_{2L})(\overline{\psi_{3R}} \gamma^\mu \psi_{4R}) = - 2(\overline{\psi_{1L}} \psi_{4R})(\overline{\psi_{3R}} \psi_{2L}) \qquad (2.39)$$

- only they can be cast into a current-current form:

$$O_1 = \frac{1}{2} [(\overline{e^+_r} \gamma_\mu d_R)(\overline{u^c_L} \gamma^\mu u_L) - (\overline{\nu^c_R} \gamma_\mu d_R)(\overline{u^c_L} \gamma^\mu d_L)] \, , \qquad (2.40)$$

$$O_2 = (\overline{e^+_L} \gamma_\mu d_L)(\overline{u^c_L} \gamma^\mu u_L) \; . \qquad (2.41)$$

However, the Yukawa couplings of colour triplet Higgs bosons to pairs of light fermions are extremely weak. Thus the contributions of Higgs boson exchange to nucleon decay can be estimated to be suppressed compared to the contributions of gauge boson exchange already in amplitude by a factor [2.4]

$$O(\frac{m_X^2}{m_{H_X}^2} \frac{\mu^2}{m_W^2}) \lesssim 10^{-4} \qquad (2.42)$$

for the natural mass relation

$$m_{H_X} \approx m_X , \tag{2.43}$$

μ being the typical hadronic mass scale where the B violating processes take place,

$$\mu = O(1\ \mathrm{GeV}) . \tag{2.44}$$

Consequently, it is usually assumed that the rôle of Higgs bosons in proton decay can be ignored and that one's attention can be restricted to B and L violating gauge interactions.

Introducing for the moment a conceivable but heavy right-handed neutrino ν_R, the most general couplings of the gauge boson $SU(2)_L$ doublets

$$X \equiv \begin{pmatrix} X \\ Y \end{pmatrix} \tag{2.45}$$

and

$$X' \equiv \begin{pmatrix} X'^{\dagger} \\ Y'^{\dagger} \end{pmatrix} \tag{2.46}$$

to fermionic gauge eigenstates read [2.4,2.6]

$$\begin{aligned} L = \frac{g}{\sqrt{2}} \{ & [(\overline{e^{+}_{R}}\, \gamma_{\mu}\, d_{iR}) + (\overline{e^{+}_{L}}\, \gamma_{\mu}\, d_{iL}) + \epsilon_{ijk}(\overline{u_{jL}}\, \gamma_{\mu}\, u^{c}_{k\,L})] X_i^{\mu} + \\ & + [-(\overline{\nu^{c}_{R}}\, \gamma_{\mu}\, d_{iR}) - (\overline{e^{+}_{L}}\, \gamma_{\mu}\, u_{iL}) + \epsilon_{ijk}(\overline{d_{jL}}\, \gamma_{\mu}\, u^{c}_{k\,L})] Y_i^{\mu} - \\ & - [(\overline{\nu^{c}_{R}}\, \gamma_{\mu}\, u_{iR}) + (\overline{\nu^{c}_{L}}\, \gamma_{\mu}\, u_{iL}) + \epsilon_{ijk}(\overline{d_{jL}}\, \gamma_{\mu}\, d^{c}_{k\,L})] X_i'^{\mu} + \\ & + [-(\overline{e^{+}_{R}}\, \gamma_{\mu}\, u_{iR}) - (\overline{\nu^{c}_{L}}\, \gamma_{\mu}\, d_{iL}) + \epsilon_{ijk}(\overline{u_{jL}}\, \gamma_{\mu}\, d^{c}_{k\,L})] Y_i'^{\mu} \} + \\ & + \mathrm{h.c.} \end{aligned} \tag{2.47}$$

The fermion fields have to be understood as generation space vectors, g denotes the GUT gauge coupling constant.

From these renormalizable interactions a B and L violating, $SU(3)_C \times SU(2)_L \times U(1)_Y$ invariant four-fermion interaction valid at low energies can be derived in second order of perturbation expansion (Fig. 2). Obviously, there must exist a formulation of this effective theory entirely in terms of the operators O_1 and O_2. This can be verified by making extensive use of the Fierz transformation

$$(\overline{\psi_{1L}}\, \gamma_\mu\, \psi_{2L})(\overline{\psi_{3L}}\, \gamma^\mu\, \psi_{4L}) = (\overline{\psi_{1L}}\, \gamma_\mu\, \psi_{4L})(\overline{\psi_{3L}}\, \gamma^\mu\, \psi_{2L}) \qquad (2.48)$$

as well as of the operator identity

$$(\overline{\psi_{1R}}\, \gamma_\mu\, \psi_{2R}) = - (\overline{\psi^c_{2\,L}}\, \gamma_\mu\, \psi^c_{1\,L}) \ . \qquad (2.49)$$

However, without detailed knowledge of the underlying grand unified model the rotation of the fermionic gauge eigenstates into the corresponding mass eigenstates cannot be performed. Hence, one has to rely on the assumption that all mixing angles arising in the course of diagonalization of the fermionic mass matrices can be neglected. Under this assumption the most general effective Lagrangian responsible for proton decay is given by

$$\begin{aligned} L = 4 \frac{G}{\sqrt{2}} \epsilon_{ijk} \, [& (1+x)(\overline{e^+_R}\, \gamma_\mu\, d_{iR})(\overline{u^c_{k\,L}}\, \gamma^\mu\, u_{jL}) + \\ & + 2(\overline{e^+_L}\, \gamma_\mu\, d_{iL})(\overline{u^c_{k\,L}}\, \gamma^\mu\, u_{jL}) - \\ & - (1+x)(\overline{\nu^c_{e\,R}}\, \gamma_\mu\, d_{iR})(\overline{u^c_{k\,L}}\, \gamma^\mu\, d_{jL}) + \\ & + (\overline{\mu^+_R}\, \gamma_\mu\, s_{iR})(\overline{u^c_{k\,L}}\, \gamma^\mu\, u_{jL}) + \\ & + (\overline{\mu^+_L}\, \gamma_\mu\, s_{iL})(\overline{u^c_{k\,L}}\, \gamma^\mu\, u_{jL}) - \end{aligned} \qquad (2.50)$$

$$- \quad (\overline{\nu^c_{\mu R}}\, \gamma_\mu\, s_{iR})(\overline{u^c_{k L}}\, \gamma^\mu\, d_{jL})] + \text{h.c.}$$

Unfortunately, even in GUTs some historical ballast is dragged along when defining the four-fermion interaction strength G by analogy to the Fermi coupling constant G_F:

$$\frac{G}{\sqrt{2}} := \frac{g^2}{8 m_X^2} \,. \qquad (2.51)$$

The parameter x denotes the squared ratio of the masses of the gauge bosons X and X',

$$x := \frac{m_X^2}{m_{X'}^2} \,, \qquad (2.52)$$

characterizing by that way the underlying grand unified model. A priori, x can take any positive value,

$$0 \leq x < \infty \,. \qquad (2.53)$$

According to the relative magnitude of the gauge boson masses m_X and $m_{X'}$ there are, however, some special cases which deserve particular interest:

Mass relation	x
$m_X \ll m_{X'}$	0
$m_X \gg m_{X'}$	∞
$m_X = m_{X'}$	1

The case $m_X \ll m_{X'}$ is realized in SU(5) since SU(5) contains only the superheavy gauge bosons X but not X', which is equivalent to $m_{X'} \to \infty$. SO(10) contains among its superheavy particle content both types of B and L violating gauge bosons. Besides the spontaneous breakdown of SO(10) via SU(5) there are, however, symmetry breaking scenarios which either leave X' lighter than X (ideally $m_{X} \to \infty$), or which maintain the left-right symmetry inherent in SO(10) down to low energies, entailing thereby $m_X = m_{X'}$.

Gauge boson exchange in second order of perturbation theory introduces a factor α/m_X^2 in the nucleon decay matrix element. Hence, on dimensional grounds the theoretically expected proton lifetime has to read

$$\tau_p = \frac{1}{\alpha^2}\frac{m_X^4}{M_p^5}C\,, \tag{2.54}$$

where C is a dimensionless constant summarizing the whole hadronic aspect of proton decay, of approximate order of magnitude

$$C \gtrsim O(1)\,. \tag{2.55}$$

Assuming C = 1 one obtains from Eqs. (2.12) and (2.17) a naive estimate of the proton lifetime:

$$\tau_p \approx 1.6\cdot 10^{29}\ \text{yr}\,. \tag{2.56}$$

Basically, this simple consideration has stimulated all of the current searches for proton decay and determined the design of these experiments.

III. Baryon Number Violation in Supersymmetric GUTs

Essentially the only convincing reason for dealing with supersymmetry (SUSY) is a purely theoretical one, namely the so-called Haag-Łopuszański-Sohnius theorem [3.6]: The most general algebra of generators - which are assumed to act additively on initial multi-particle states and to connect only single-particle states of the same mass - of symmetry transformations of a non-trivial S-matrix in a relativistic quantum field theory describing solely massive particles is a graded Lie algebra (known as "supersymmetry algebra" [3.8]) spanned by the energy-momentum operator P_μ, the generator of the homogeneous Lorentz transformations $M_{\mu\nu}$, a finite number of hermitean scalar charges generating a compact Lie group, as well as a set of $N = 1,2,\ldots$ fermionic charges Q^L_α ($L = 1,2,\ldots,N$; $\alpha = 1,2$) and their hermitean conjugates $\bar{Q}^L_{\dot\alpha} = (Q^L_\alpha)^+$, transforming like spinors of rank 1 under the homogeneous Lorentz group and forming an N-dimensional representation of the internal symmetry group. This graded Lie algebra of "N-extended supersymmetry" is the only possible non-trivial unification of internal symmetries and the geometrical space-time symmetries of the Poincaré algebra within a relativistic quantum field theory. Thus it describes the maximal invariance structure of the S-matrix.

In the case that there is only one fermionic generator Q_α present in the theory, i.e. $N = 1$, the SUSY algebra takes a particular simple form

$$\{Q_\alpha,\bar{Q}_{\dot\beta}\} = 2\sigma^\mu_{\alpha\dot\beta} P_\mu , \qquad (3.1)$$

$$\{Q_\alpha,Q_\beta\} = \{\bar{Q}_{\dot\alpha},\bar{Q}_{\dot\beta}\} = 0 , \qquad (3.2)$$

$$[Q_\alpha,P_\mu] = [\bar{Q}_{\dot\alpha},P_\mu] = 0 , \qquad (3.3)$$

$$\begin{aligned} [Q_\alpha,M_{\mu\nu}] &= \frac{1}{2} (\sigma_{\mu\nu})_\alpha{}^\beta Q_\beta \\ [\bar{Q}_{\dot\alpha},M_{\mu\nu}] &= -\frac{1}{2} \bar{Q}_{\dot\beta}(\bar{\sigma}_{\mu\nu})^{\dot\beta}{}_{\dot\alpha} . \end{aligned} \qquad (3.4)$$

Here σ^μ are the generalized Pauli matrices

$$\sigma^\mu_{\alpha\dot\beta} = (1,\sigma^m) , \qquad m = 1,2,3 , \qquad (3.5)$$

in terms of which the generators of the homogeneous Lorentz group in the spinor representation are given by

$$\begin{aligned} \sigma^{\mu\nu} &= \frac{i}{2} (\sigma^\mu \bar{\sigma}^\nu - \sigma^\nu \bar{\sigma}^\mu) \\ \bar{\sigma}^{\mu\nu} &= \frac{i}{2} (\bar{\sigma}^\mu \sigma^\nu - \bar{\sigma}^\nu \sigma^\mu) . \end{aligned} \qquad (3.6)$$

As can be seen from Eq. (3.4), the fermionic generators Q_α and $\bar{Q}_{\dot\alpha}$ change the spin of a given state by half a unit. Hence they turn bosons into fermions and vice versa,

$$Q_\alpha|B\rangle = \sqrt{2E}|F\rangle$$
$$Q_\alpha|F\rangle = \sqrt{2E}|B\rangle \,, \qquad (3.7)$$

without having any effect on the internal quantum numbers of these states. As a consequence of this, supersymmetric theories show a perfect Bose-Fermi invariance: The numbers of bosonic and fermionic degrees of freedom are always identical. All particles appear in pairs of superpartners carrying the same internal quantum numbers but belonging to different statistics. Moreover, as long as supersymmetry is left unbroken, these superpartners have to be degenerate in mass, i.e.

$$m_B^2 = m_F^2 \,. \qquad (3.8)$$

Leaving aside gravity and bearing in mind the troubles of theories which contain spin-3/2 fields but no spin-2 field, one can imagine just two possibilities for pairing particles which differ in spin by half a unit, constituting in this way a so-called supermultiplet [3.3,3.5,3.8,3.9]: The "chiral supermultiplet" contains a two-component Weyl spinor ψ (e.g. a left- or right-handed quark or lepton) and a complex scalar field A ("squark", "slepton"). The "vector supermultiplet" contains a vector field V^μ (e.g. a gluon, W, Z, or photon) and a two-component Weyl spinor λ ("gluino","wino", "zino", "photino"). Re-stated more technically, these two supermultiplets form the massless irreducible representations of lowest spin for (N = 1) supersymmetry.

However, one a little bit unpleasant feature of supersymmetric theories has to be mentioned, that is the doubling of the number of particles in the theory. In a renormalizable theory all vector bosons have to be gauge bosons. Consequently, all vector supermultiplets must transform according to the adjoint representation of the gauge group. On the other hand, the known fermions constitute a fully complex representation of $SU(3)_C \times SU(2)_L \times U(1)_Y$. They thus can only be described by chiral supermultiplets. Thence, no pair of known particles can be regarded as related by supersymmetry.

In addition to the Haag-Łopuszański-Sohnius theorem there is a second somewhat technical but nevertheless grave motivation for supersymmetry, which runs under the heading "(gauge) hierarchy problem"*. What is meant by the term "hierarchy problem" is the instability of light scalar masses against radiative corrections in the presence of a heavy mass scale. In quantum field theory the mass of an elementary scalar boson receives in higher orders of perturbation theory quadratically divergent radiative corrections. (In contrast to that, fermion masses are protected by chiral symmetries.) These radiative corrections tend to renormalize the scalar mass towards the natural mass scale of the theory. Hence, disregarding dubious fine tuning order by order in perturbation expansion, it appears, in general, impossible to maintain two (or more) vastly different mass scales in the theory. The desired hierarchy for the gauge symmetry breakdown (2.1) in grand unified theories,

$$\frac{m_W}{m_X} \approx 10^{-13} \ll 1 \, , \tag{3.9}$$

could not be realized. The problem is even more serious in the standard model. Here the mass of the Glashow-Salam-Weinberg Higgs boson would be increased by the quadratic mass renormalization to the order of magnitude of the Planck mass. In either case, the spontaneous breakdown of $SU(2)_I \times U(1)_Y$ would already take place at ultrahigh energies.

The solution to this problem is provided by supersymmetry. Due to the high degree of symmetry implied by the Bose-Fermi invariance, supersymmetric quantum field theories show a by far less divergent high-energy behaviour, which manifests itself by a drastically reduced number of possible counterterms in the theory. This observation is expressed more precisely by the non-renormalization theorem [3.1,3.2,3.4,3.7,3.10,3.18]. In supersymmetric theories all mass terms, Yukawa couplings, and scalar self-interactions are described by the so-called superpotential. Now, the non-renormalization theorem states that the superpotential is not renormalized at all. Consequently, the parameters of the superpotential (masses, Yukawa coupling constants) are not subject to any renormalization

*) For a more complete discussion of the hierarchy problem see Ref. [1.2].

independent of the wave-function renormalization. In fact, they only are multiplicatively renormalized by appropriate powers of the wave-function renormalization constants. The wave-function renormalization being at most logarithmically divergent at any rate, there is no room for quadratic divergences. Supersymmetric theories are completely free of quadratic divergences*. The reason behind is a mutual cancellation of the contributions of boson and fermion loops to the quadratically divergent parts of all radiative corrections order by order in the loop expansion. Thus, supersymmetry offers a solution to the technical aspect of the hierarchy problem. It does not explain the origin of the tiny mass ratio (3.9). It stabilizes, however, this ratio against radiative corrections.

Supersymmetric theories are most easily formulated in terms of superfields. Superfields [3.3,3.5,3.8,3.9] summarize a finite number of boson and fermion fields in one single object transforming linearly but, in general, reducibly under SUSY transformations. By imposing certain constraints, they can be restricted to irreducible SUSY representations, describing then, for instance, the chiral or the vector supermultiplet. The use of superfields considerably simplifies the construction of supersymmetric models as well as the discussion of their internal symmetries. A remarkable feature of superfields is the behaviour of their highest-dimensional component field under SUSY transformations. The highest component, e.g. the "F component" of a chiral superfield or the "D component" of a vector superfield, is always transformed into a total space-time derivative, the space-time integral of which is usually assumed to vanish. Thus the space-time integral of the highest component is invariant under SUSY transformations.

In a realistic $SU(3)_C \times SU(2)_L \times U(1)_Y \times SUSY$ invariant theory the minimal set of chiral superfields is [3.13-3.15,3.17]

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*) The only exception is the one-loop contribution to the "D term" of a U(1) gauge factor, which is proportional to Tr[Q] and thus, according to the spirit of grand unification, vanishes in GUTs.

Superfield	$SU(3)_C$	$SU(2)_L$	Y	B	L
Q_L	3	2	$\frac{1}{6}$	$\frac{1}{3}$	0
U_R	3	1	$\frac{2}{3}$	$\frac{1}{3}$	0
D_R	3	1	$-\frac{1}{3}$	$\frac{1}{3}$	0
L_L	1	2	$-\frac{1}{2}$	0	1
E_R	1	1	-1	0	1
H_L	1	2	$\frac{1}{2}$	0	0
H'_L	1	2	$-\frac{1}{2}$	0	0

The two left-handed Higgs superfields of opposite hypercharge are required for generating mass terms for both $Q = 2/3$ and $Q = -1/3$ quarks via supersymmetric Yukawa couplings. In contrast to the standard model, supersymmetry forbids the use of one Higgs and its charge conjugate for this purpose because charge conjugation flips the chirality of a superfield.

As has been done in the previous section for conventional GUTs, one can now classify the possible $SU(3)_C \times SU(2)_L \times U(1)_Y$ invariant supersymmetric effective operators which violate baryon number. Due to the existence of squarks and sleptons, i.e. scalar particles carrying non-vanishing baryon or lepton number, B violating operators of dimension less than six can be formed. $SU(3)_C$ gauge invariance requires these operators to be a product of at least three quark superfields. The supersymmetric component of a product of three chiral superfields has at least dimension four. Consequently, the lowest possible dimension for these operators is four [3.13]. As far as proton decay is concerned, it is sufficient to consider only the baryon number violating operators of dimension four and five. The supersymmetric operators of that kind composed by superfields from the minimal set given above are [3.13-3.15,3.17]

Dimension	Operator	ΔB	ΔL
4	$(U_R\ D_R\ D_R)_F$	-1	0
5	$(Q_L\ Q_L\ D_R)_D$	-1	0
	$(Q_L\ Q_L\ Q_L\ L_L)_F$	-1	-1
	$(U_R\ U_R\ D_R\ E_R)_F$	-1	-1
	$(Q_L\ Q_L\ Q_L\ H_L')_F$	-1	0

Again the generation structure has not been exhibited and the obvious contraction of $SU(3)_C$ and $SU(2)_L$ indices in order to form gauge singlets is implicitly understood.

In a low-energy Lagrangian an effective operator of dimension d will appear multiplied by a coefficient G of the order of magnitude

$$G = O(m_X^{4-d}) \ . \tag{3.10}$$

Hence, the B violating operator of dimension four is not suppressed at all by inverse powers of the superheavy mass scale m_X. Proton decay induced by this operator would proceed with a disastrously short lifetime of the order of seconds. It is, however, a simple task to get rid of this dangerous operator. All one has to do is to impose the fermion reflection symmetry [3.14,3.15,3.17]

$$F \to -F\ , \qquad F = Q_L,\ U_R,\ D_R,\ L_L,\ E_R\ . \tag{3.11}$$

This requirement eliminates the operators $(U_R\ D_R\ D_R)_F$, $(Q_L\ Q_L\ D_R)_D$, and $(Q_L\ Q_L\ Q_L\ H_L')$ from the theory*. One is left with the dimension-five operators

$$O_L^{(5)} = (Q_L\ Q_L\ Q_L\ L_L)_F \tag{3.12}$$

and

$$O_R^{(5)} = (U_R\, U_R\, D_R\, E_R)_F \; . \qquad (3.13)$$

These operators respect the selection rule

$$\Delta B = \Delta L \; , \qquad (3.14)$$

as has been the case for the B and L violating four-fermion operators of conventional GUTs. Consequently, irrespective of whether or not the grand unified model is supersymmetric, if nucleon decay is controlled by effective operators of lowest acceptable dimension, the nucleons are expected to decay into anti-leptons, i.e. according to Eq. (2.34). Dimension-six operators are suppressed by an additional power of the superheavy mass scale when compared to dimension-five operators. Thus, the dimension-five operators $O_L^{(5)}$ and $O_R^{(5)}$ are supposed to give the dominant contribution to supersymmetric proton decay. Note the chirality structure of these operators,

$$O_L^{(5)} \sim (L\ L\ L\ L) \; , \qquad (3.15)$$

$$O_R^{(5)} \sim (R\ R\ R\ R) \; , \qquad (3.16)$$

reflecting, of course, nothing else but their $SU(2)_L$ content.

The origin of the dimension-five operators $O_L^{(5)}$ and $O_R^{(5)}$ might be found in Higgs exchange. In a supersymmetric theory the super-Yukawa coupling**

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*) The B violating dimension-four operator poses a more serious problem in superstring theories. There the discrete symmetry needed to exorcize this operator cannot simply be postulated but has to arise in the course of compactification from ten to four dimensions. Until now, no example of a Calabi-Yau manifold possessing an adequate symmetry has been found [3.19].

**) The imposed fermion reflection symmetry guarantees that Higgs superfields couple only to pairs of matter superfields [3.12].

$$(F\ F\ H)_F$$

of a Higgs superfield H to two matter superfields F describes not only the ordinary Yukawa coupling

$$\overline{f^c}\ f\ H$$

of a Higgs boson $H \equiv A_H$ to two fermions $f \equiv \psi_f$ (Fig. 3.a) but also the coupling

$$\overline{f^c}\ \tilde{H}(sf)$$

of the Higgs fermion $\tilde{H} \equiv \psi_H$ to a fermion f and a sfermion $sf \equiv A_F$ (Fig. 3.b) as well as the scalar coupling

$$(sf)(sf)H$$

of the Higgs boson H to two sfermions sf (Fig. 3.c). Hence, in second order of perturbation expansion, the super-Yukawa couplings generate an interaction between two fermions and two sfermions, either by exchange of one of the colour triplet Higgs bosons H_X enumerated in Section II (Fig. 4.a) or by exchange of its supersymmetric counterpart $\tilde{H}_X$ (Fig. 4.b). At low energies an effective theory may be derived by integrating out all superheavy degrees of freedom. The two-fermion-two-sfermion interaction then reduces to the dimension-five operator

$$O^{(5)} \sim (sf)(sf)(f)(f) \qquad (3.17)$$

sketched in Fig. 5. In either case the resulting effective coupling constant G_S which multiplies $O^{(5)}$ proves to be of order

$$G_S = O(\frac{1}{m_{H_X}}) \ . \qquad (3.18)$$

(Recall that the coupling constant of the scalar interaction $(sf)(sf)H_X$ is of the order of magnitude of the superheavy Higgs mass m_{H_X}.)

Of course, in order to be relevant for proton decay, $O^{(5)}$ must involve the appropriate particle combination, in particular contain just one leptonic field. Furthermore, the two sfermions have to be converted into fermions. The exchange of a standard-model gauge fermion, i.e. a gluino $\tilde{g} \equiv \lambda_g$, wino $\tilde{W} \equiv \lambda_W$, or a bino $\tilde{B} \equiv \lambda_B$, between the two sfermions does this job (Fig. 6). A Majorana mass term takes care of the necessary chirality flip of the gauge fermions. In this way the required four-fermion interaction

$$L \sim \frac{G_S}{m_\lambda} (f)(f)(f)(f) \ , \qquad \lambda = \tilde{W}, \tilde{B}, \tilde{g} \ , \qquad (3.19)$$

is generated via the one-loop diagram in Fig. 6.

According to the above sketched mechanism, only $O_L^{(5)}$ can play a rôle in supersymmetric proton decay [3.16]. Due to colour anti-symmetrization the two superfields U_R in $O_R^{(5)}$, Eq. (3.13), have to belong to different generations since otherwise $O_R^{(5)}$ vanishes identically. However, all right-handed fields being $SU(2)_L$ singlets, only $\tilde{B}$ and $\tilde{g}$ can be exchanged between them. Thus, at the gauge vertices the flavour of the sfermions is transfered to the external fermions. Consequently, the four-fermion interaction resulting from $O_R^{(5)}$ unavoidably involves a heavy Q = 2/3 quark. Hence, on kinematical grounds $O_R^{(5)}$ is of no interest for proton decay.

In supergravity theories, however, this statement is no longer true. There the supergravity breaking at the Planck scale by the hidden sector of the model induces soft supersymmetry breaking terms proportional to the gravitino mass $m_{3/2}$ in the effective low-energy theory. These soft breaking terms give rise to off-diagonal entries of the form $Am_{3/2}m_f$ in the sfermion mass matrix, where A denotes the Polonyi constant and m_f is the mass of the corresponding fermion. The off-diagonal matrix elements lead to a mass mixing of the sfermions sf_L and sf_R related to the left- and right-handed fermion components, resp. This left-right mixing may transform the $SU(2)_L$ singlets sf_R contained in $O_R^{(5)}$ into the $SU(2)_L$ doublets sf_L which then can interact via wino exchange in a flavour non-diagonal way (Fig. 7). Thus the resulting four-fermion interaction of mixed chirality structure,

$$O^{(6)}_{LR} \sim (L\ L\ R\ R)\ , \tag{3.20}$$

might well give an energetically allowed contribution to proton decay [4.6].

In conventional GUTs the proton lifetime scales with the fourth power of the GUT mass,

$$\tau_{p,conv} = O(m_X^4)\ . \tag{3.21}$$

In contrast to that, supersymmetric proton decay - mediated by B violating dimension-five operators - proceeds with a lifetime proportional only to the square of the superheavy mass scale,

$$\tau_{p,SUSY} = O(m_X^2)\ . \tag{3.22}$$

One might fear that this lifetime will be by far too short, making thus any model which incorporates B violating dimension-five operators phenomenologically unacceptable. However, in supersymmetric GUTs the contributions of the superpartners of the known light particles to the renormalization group equations for the standard-model gauge coupling constants tend to increase the grand unification mass scale as well as the GUT gauge coupling constant (Fig. 1.b) [3.11]. Including two-loop effects one finds in the minimal supersymmetric SU(5) model [3.16]

$$m_X \simeq 6\cdot 10^{16}\ \Lambda_{\overline{MS}}\ , \tag{3.23}$$

i.e. an enlargement of the value obtained in conventional GUTs by nearly two orders of magnitude,

$$\frac{m_X^{(SUSY)}}{m_X^{(GUT)}} \simeq 40\ . \tag{3.24}$$

In addition, in case the baryon number violation originates in super-Yukawa couplings, there is a lot of mixing angles and small mass ratios $m_f/m_{\tilde{f}}$ providing further suppression of the proton decay rate. The combined effects may be sufficient to prevent any conflict with experiment.

IV. Theoretical Predictions

As has been discussed in some detail in Sections II and III, in one way or other conventional as well as supersymmetric GUTs predict the existence of B and L violating effective four-fermion interactions

$$\mathcal{L} \sim G(q\ q\ q\ \ell) \qquad (4.1)$$

between three quarks and one lepton (Fig. 8). The coupling strength G is fixed in terms of the superheavy mass scale m_X by the type of grand unified model,

$$G = O\left(\frac{1}{m_X^n}\right), \qquad \begin{array}{l} n = 2 \text{ in conventional GUTs,} \\ n = 1 \text{ in supersymmetric GUTs.} \end{array} \qquad (4.2)$$

In order to disintegrate a proton or bound neutron, the four-fermion interaction (4.1) may act upon the quarks inside this nucleon, in principle, in three different ways (Fig. 9):

(i) Two-quark annihilation (Fig. 9.a):

$$q + q \to q^c + \ell^c \ . \qquad (4.3)$$

The anti-quark produced in this process recombines with the remaining spectator quark to a mesonic final state.

(ii) Three-quark fusion (Fig. 9.b):

$$q + q + q \to \ell^c \ . \qquad (4.4)$$

Energy-momentum balance is restored by the emission of a meson or photon before or after this process takes place.

(iii) Quark decay (Fig. 9.c):

$$q \to q^c + q^c + \ell^c \ . \tag{4.5}$$

Quark decay leads to a final state which contains at least two mesons. Consequently it is suppressed by the available phase space. On the other hand, the two-quark and three-quark mechanisms appear to be roughly, i.e. within a factor 3 or so in amplitude, comparable (see Ref. [4.3] and references therein).

The problem which introduces the largest source of uncertainties in the calculations of nucleon decay widths is the evaluation of hadronic matrix elements of interactions formulated in terms of quark field operators. In order to bring this translation from quark level to the hadronic level about, a great variety of hadronic models, like nonrelativistic SU(6), the MIT bag, Bethe-Salpeter amplitudes, or chiral Lagrangians, have been employed [1.2].

The bulk of the theoretical investigations of proton decay concentrated on the minimal SU(5) model. (This model is minimal in the sense that it requires the smallest conceivable Higgs sector.) For conventional GUTs the following picture emerged [1.2]: Splitting off the dependence on m_X, the proton lifetime τ_p is usually parametrized according to

$$\tau_p = a_p \left(\frac{m_X}{10^{14}\ \text{GeV}}\right)^4 . \tag{4.6}$$

The predictions for the parameter a_p cover almost two orders of magnitude,

$$0.2 \le \frac{a_p}{10^{28}\ \text{yr}} \le 14 \ . \tag{4.7}$$

Thus the value (2.17) for the GUT mass, $m_X = 2.4\cdot 10^{14}$ GeV, corresponds to the range

$$0.5 \le \frac{\tau_p}{10^{29}\ \text{yr}} \le 46 \tag{4.8}$$

for the proton lifetime, in accordance with the rough estimate (2.56). The ratio of bound neutron versus proton lifetime appears to be rather close to unity,

$$0.8 \leq \frac{\tau_n}{\tau_p} \leq 1.1 \ . \tag{4.9}$$

Most calculations obtain a slightly larger lifetime for the proton than for the bound neutron [1.2]. The branching ratios

$$B(p \to \ell^c + M) = \frac{\Gamma(p \to \ell^c + M)}{\Gamma(p \to \text{2-body})} \ , \quad M = \pi, \eta, \rho, \omega, K, K^* \ , \quad \ell^c = e^+, \nu_e^{\,c}, \mu^+, \nu_\mu^{\,c} \ , \tag{4.10}$$

found for two-body proton decay within SU(5) are listed in Table 1. The corresponding decay rates for bound neutrons are not independent but related to the ones for proton decay by strong isospin. One observes a still remarkably large spread for the decay modes $p \to e^+\eta$, $e^+\rho^0$, $\nu_e^{\,c}\rho^+$, and μ^+K^0. The decay channels

$$\begin{aligned} &p \to e^+ + \pi^0 \\ &p \to e^+ + \omega \\ &p \to \nu_e^{\,c} + \pi^+ \end{aligned} \tag{4.11}$$

and

$$\begin{aligned} &n \to e^+ + \pi^- \\ &n \to \nu_e^{\,c} + \pi^0 \\ &n \to \nu_e^{\,c} + \omega \end{aligned} \tag{4.12}$$

prove to be dominant all over the whole spectrum of GUTs (see e.g. Ref. [4.1]). Considering the electromagnetic nucleon decay

$$N \to \ell^c + \gamma \ , \tag{4.13}$$

the ratio of the photonic to the pionic decay width,

$$\frac{\Gamma(N \to \ell^c + \gamma)}{\Gamma(N \to \ell^c + \pi^0)} \approx 10^{-3} \ , \tag{4.14}$$

turns out to be independent of the grand unified model [4.1].

In supersymmetric GUTs it is more difficult to arrive at firm conclusions. The proton lifetime is determined by some unknown parameters, mainly by the masses of so far unobserved superpartners. If the dimension-five operators discussed in Section III arise from Higgs exchange and if the exchanged colour-triplet Higgs fields belong to the same supermultiplets which give masses to the known fermions, so that the B violating super-Yukawa couplings have to be proportional to the corresponding fermion masses, one expects proton decay to favour decay channels involving heavy fermions. Accordingly, supersymmetric proton decay should manifest itself by a distinct dominance of strange decay modes,

$$N \to \nu_\ell^{\,c} + K\,, \qquad \ell = \tau,\mu,e\,, \tag{4.15}$$

and

$$N \to \ell^+ + K\,, \qquad \ell = \mu,e\,. \tag{4.16}$$

Neglecting contributions suppressed by small mixing angles or by the mass ratio $m_u/m_c \simeq 0.003$, one obtains, for instance, in the minimal supersymmetric SU(5) model [4.2]

$$\begin{aligned}\Gamma(p \to \nu_\mu^{\,c}K^+) : \Gamma(p \to \nu_\mu^{\,c}\pi^+) : \Gamma(p \to \nu_e^{\,c}K^+) &= \\ = 1 : O(\mathrm{tg}^2\Theta_C) : O(\frac{m_d^2}{m_s^2}\,\mathrm{ctg}^2\Theta_C) &\simeq 1 : 0.11 : 0.048\,,\end{aligned} \tag{4.17}$$

where Θ_C denotes the Cabibbo angle.

The recent activities in the field of proton decay can be divided into two main streams: supersymmetric proton decay and attempts to reconcile the theoretical predictions for the proton lifetime with the experimental findings (2.22). Chadha, Daniel, and varying collaborators have been undefatigably at work in investigating supersymmetric GUTs with the help of chiral Lagrangians [4.4]. The question concerning the amount of the contribution of gluino dressing of dimension-five operators has been settled [4.5].

By now, the experimental lower limit for the proton lifetime has been pushed more than two orders of magnitude above the central value

$$\tau_p \simeq 4.9 \cdot 10^{29 \pm 1}\ \mathrm{yr} \tag{4.18}$$

of the theoretical predictions (4.8) for great desert GUTs (see next section). In view of this, great efforts have been undertaken in order to embed this experimental result in the theoretical framework [1.3-1.5]. These attempts fall into two categories, according to the mass scale where they try, viz. the grand unification mass scale m_X [4.7,4.8] or the hadronic mass scale $\mu = O(1\ \mathrm{GeV})$ [4.9,4.10]. On the one hand, it is always possible to increase the GUT mass m_X by an appropriate enlargement of the grand unified model, that is by introducing new particle thresholds (gauge bosons, fermions, Higgs scalars) in the great desert between m_W and m_X. These new particles will then serve to slow down the rate of approach of the standard-model gauge coupling constants, delaying thereby their grand unification. Any modification of this kind is only constrained by the requirement not to destroy the successful prediction of the Weinberg angle in great desert GUTs. A severe drawback of non-minimal GUTs, however, is their drastically reduced predictive power. – On the other hand, one can try to find the required suppression mechanism for proton decay within the hadronic model one adopts. The most interesting suggestion argues in favour of a Gamow-like barrier penetration factor which might diminish the nucleon decay rate by approximately one order of magnitude [4.10]. However, none of the proposed suppression mechanisms can account for the entire above-mentioned discrepancy. Consequently, leaving aside the possibility of a mysterious conspiracy of all conceivable effects (which, of course, include the uncertainty in the determination (2.16) of $\Lambda_{\overline{MS}}$), non-supersymmetric minimal SU(5) as well as all other conventional great desert GUTs are <u>definitely</u> ruled out by experiment.

V. Experimental Situation

All of the experiments dedicated to the investigation of nucleon instability are designed according to two main technologies differing in their source of decaying nucleons, viz. iron or water: fine grain tracking calorimeters and imaging water Čerenkov detectors. The currently active experiments are briefly characterized in Table 2. In addition, the Soudan 2 experiment of the Argonne-Minnesota-Oxford-Rutherford-Tufts collaboration, using a tracking calorimeter located at the Soudan iron mine (Minnesota) at a depth of 2000 m.w.e., is about to start taking data. Most of the experiments in operation have observed a number of nucleon decay candidates (Table 3). Most important, all of the reported candidates allow for a $\Delta B = \Delta L$ interpretation*. However, only improved statistics supplemented by refined background estimates will decide whether these candidates are true nucleon decays or mere neutrino-induced background. All one can say at the moment is that nucleon decay occurring at a rate of about 5% of the contained events - the overwhelming majority of them being, in any case, due to interactions of atmospheric neutrinos - cannot be excluded.

The lower bounds on the partial nucleon lifetimes

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*) The existence of unambiguously identified candidates for decay modes with charge conjugated final states,

$$n \to \ell^+ + M^- , \qquad M = \pi, \rho , \qquad (\Delta B = \Delta L)$$

as well as

$$n \to \ell^- + M^+ , \qquad M = \pi, \rho , \qquad (\Delta B = - \Delta L) ,$$

i.e. the existence of $\Delta B = - \Delta L$ nucleon decays for which no $\Delta B = \Delta L$ interpretation can be found, would indicate that there is something wrong with these candidate events. Whatever the source of $\Delta B = - \Delta L$ nucleon decays (e.g. dimension-seven operators in conventional GUTs [2.10] or dimension-six operators in supersymmetric GUTs [3.17]) might be, it is highly unlikely that $\Delta B = \Delta L$ and $\Delta B = - \Delta L$ nucleon decay modes are induced at a comparable rate.

$$\tau(N \to \ell^c + M) = \frac{1}{\Gamma(N \to \ell^c + M)} = \frac{\tau_N}{B(N \to \ell^c + M)} \qquad (5.1)$$

for $\Delta B = \Delta L$ two-body decay modes, as reported by NUSEX, IMB, and KAMIOKANDE, are compiled in Table 4. The HPW collaboration has focused on nucleon decay modes involving two or more muons. Their results are

$$\tau(p \to 2\mu + X) \geq 5 \cdot 10^{31}\ \text{yr} \qquad (90\%\ \text{c.l.}) \qquad (5.2)$$

and

$$\tau(p \to 3\mu + X) \geq 2.7 \cdot 10^{31}\ \text{yr} \qquad (90\%\ \text{c.l.})\ . \qquad (5.3)$$

The Fréjus experiment quotes a lower limit of

$$\tau(N \to \ell^{\pm} + X) > 4 \cdot 10^{31}\ \text{yr}\ , \qquad \ell = e,\mu\ , \qquad (90\%\ \text{c.l.}) \qquad (5.4)$$

for nucleon decay into charged leptons.

The comparison of theory and experiment depends, of course, on the grand unified model on which the theoretical discussion is based. For a given grand unified model, according to

$$\frac{\tau(N \to \ell^c + M)_{exp}}{\tau(N \to \ell^c + M)_{th}} = \tau(N \to \ell^c + M)_{exp} \frac{B(N \to \ell^c + M)_{th}}{\tau_{N,th}} \qquad (5.5)$$

the most stringent statement can be made for that decay mode for which the experimental partial lifetime <u>and</u> the theoretical branching ratio are largest. Both requirements are fulfilled for the decay channel

$$p \to e^+ + \pi^0\ , \qquad (5.6)$$

which belongs to the most favoured decay modes in all conventional GUTs, with a representative value for its branching ratio of roughly 1/3 (see e.g. Ref. [4.1]). At present, the best number for the partial lifetime, hence the actual lower bound, is provided by the Kamioka group [5.12],

$$\tau(p \to e^+ + \pi^0) > 3.3 \cdot 10^{32}\ \text{yr} \qquad (90\%\ \text{c.l.})\ . \qquad (5.7)$$

Consequently, adopting the central value (4.18) for the theoretical proton lifetime and a branching ratio of

$$B(p \to e^+ + \pi^0) \simeq 40\% \tag{5.8}$$

taken from Table 1, one ends up with a discrepancy of more than two orders of magnitude,

$$\frac{\tau(p \to e^+ + \pi^0)_{exp}}{\tau(p \to e^+ + \pi^0)_{th}} \gtrsim 3 \cdot 10^2 , \tag{5.9}$$

for conventional minimal SU(5). This discrepancy between theory and experiment is bound to get increased in all other great desert GUTs where additional $\not{B}$ violating gauge bosons contribute to the nucleon decay rate.

The results obtained by IMB and KAMIOKANDE clearly point out that further progress demands nucleon decay detectors of a fiducial mass minimum of about 1000 tons. In view of this, a further operation of NUSEX and the Fréjus experiment has to be challenged very seriously. On the other hand, most promising in this respect is the proposal for a Čerenkov detector of 32 kilotons total mass (22 kilotons fiducial mass) to be installed in the Kamioka mine. This detector should be able to continue the search for proton decay up to the ultimate upper lifetime limit of

$$\tau_N \simeq 10^{33}\ \mathrm{yr} \tag{5.10}$$

where any nucleon decay signal gets drowned in the background of neutrino interactions.

VI. Summary

In this report the theoretical picture of proton decay, as it emerges as a consequence of grand unification, has been drawn and subjected to the sentence of experiment.

The discussion of baryon number violation has been performed in an absolutely model-independent manner. Unfortunately, the existence of nucleon decay has not yet been established by experiment but all experimental findings are still consistent with a proton lifetime somewhere around 10^{32} years. Nevertheless, proton decay has not been found at the level predicted by the class of those grand unified theories which contain only two mass scales. Consequently, all conventional great desert GUTs (minimal SU(5) and the like) are inevitably disproved by the experiments carrying out the search for nucleon decay. This fact might be regarded as a hint that a crucial aspect of grand unified theories is not yet understood, that an important point or ingredient is still missing.

An (as I hope, convincing) reasoning in favour of supersymmetry has been given despite of the total lack of experimental evidence for it. Supersymmetric GUTs favour the nucleon decay modes into neutrinos,

$$N \to \nu_{\ell}^{c} + M , \qquad \ell = \tau,\mu,e \quad , \tag{6.1}$$

and, to a less amount, into strange mesons,

$$N \to \ell^{+} + K , \qquad \ell = \mu,e \; . \tag{6.2}$$

The experimental bounds for these decay modes are by far less restrictive than for the decay channels favoured by conventional GUTs, leaving the minimal supersymmetric GUTs alive. The experimental observation of the dominance of the decay modes (6.1) or (6.2) would not merely confirm baryon number violation but, simultaneously, point at supersymmetry and, eventually, even indicate an elementary Higgs boson [1.2].

The exchange of B violating gauge or Higgs bosons and their fermionic superpartners is not the only mechanism which can induce proton decay. There is a further source of nucleon instability in grand unified theories which should be mentioned: magnetic monopoles. Magnetic monopoles arise whenever a semi-simple GUT gauge group is spontaneously broken down to a subgroup containing an explicit U(1) factor [6.1],

$$\mathrm{GUT} \xrightarrow[m_X]{} H \times U(1) \; . \tag{6.3}$$

Their mass m_M is predicted to be superheavy,

$$m_M = O(\frac{m_X}{\alpha}) \gtrsim 10^{16}\ \mathrm{GeV}\ . \qquad (6.4)$$

These monopoles M may catalyze B violating processes like [6.3]

$$p + M \to e^+ + M + \text{mesons}\ . \qquad (6.5)$$

The corresponding cross sections are comparable to those for strong interactions [6.4,6.7]. In particular, they are not suppressed by the exponential factors $\exp(-2\pi/\alpha)$ characteristic of non-perturbative effects or by powers of the superheavy mass scale m_X [6.2,6.3,6.5]. Consequently, magnetic monopoles would reveal themselves by a chain of successive nucleon decays crossing the detector volume [6.6]. A search for monopole catalyzed nucleon decay has been performed by NUSEX, IMB, and KAMIOKANDE. However, until now there is no evidence of magnetic monopoles [5.6,5.11, 6.8].

Far beyond doubt, superstrings represent the most promising development taking place at present. Many of the ideas developed, investigated and pursued within the last decade or so - like grand unification, supersymmetry, supergravity and its hidden sector, extra space-time dimensions and their compactification (i.e. generalized Kaluza-Klein theories), and, last but not least, the concept of fundamental objects of finite extension (strings) - now seem to converge and fuse to a single theoretical concept, viz. superstring theories. (This observation adds another case for superstrings to the already well-known ones.) As time goes by, the superstring theory - completely determined only by the requirement to formulate a consistent quantum theory - might prove to have to say the final word on proton decay.

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Table 1

Approximate ranges of the branching ratios

$$B(N \to \ell^c + M) = \frac{\Gamma(N \to \ell^c + M)}{\Gamma(N \to 2\text{-body})}$$

for the two-body proton decay $p \to \ell^c + M$ in the minimal conventional SU(5) model.

Decay mode	Branching ratio [%]
$p \to e^+ \pi^0$	31 - 46
$p \to e^+ \eta$	0 - 8
$p \to e^+ \rho^0$	2 - 18
$p \to e^+ \omega$	15 - 29
$p \to \nu_e^c \pi^+$	11 - 17
$p \to \nu_e^c \rho^+$	1 - 7
$p \to \mu^+ K^0$	1 - 20
$p \to \nu_\mu^c K^+$	0 - 1

Table 2

Currently operational nucleon decay experiments.

Detector type	Experiment	Collaboration	Location	Depth [m.w.e.]	Total mass [t]	Fiducial mass [t]	Comments	Ref.
Fine grain tracking calorimeter	KGF	Tata-Osaka-Tokyo	Kolar Gold Fields (India)	7600	140	60	proportional counters	[5.1] [5.2]
	NUSEX (Nucleon Stability Experiment)	Frascati-Milano-Torino-CERN	Mont Blanc tunnel	5000	150	100	streamer tubes	[5.3] [5.6]
	Fréjus	Aachen-Orsay-Palaiseau-Saclay-Wuppertal	Fréjus tunnel	4400	912	600	flash chambers, Geiger tubes	
Imaging water Čerenkov	IMB	Irvine-Michigan-Brookhaven NL	Morton-Thiokol salt mine (Ohio)	1570	8000	3300	photomultiplier tubes	[5.5] [5.7] [5.8]
	KAMIOKANDE (Kamioka Nucleon Decay Experiment)	KEK-Tokyo-Niigata-Tsukuba	Kamioka mine (Japan)	2700	3000	880	photomultiplier tubes	
	HPW	Harvard-Purdue-Wisconsin	Silver King mine (Utah)	1500	780	560	photomultiplier tubes, proportional wire chambers	

Table 3

Number of fully contained nucleon decay candidate events observed by the currently operative nucleon decay experiments.

Experiment	Number of candidates
KGF	5
NUSEX	3
Fréjus	0
IMB	21
KAMIOKANDE	4
HPW	3

Table 4

Experimental lower bounds on the partial nucleon lifetimes

$$\tau(N \to \ell^c + M) = \frac{1}{\Gamma(N \to \ell^c + M)} = \frac{\tau_N}{B(N \to \ell^c + M)}$$

in units of 10^{31} yr (90% c.l., no background subtraction).

(a) Proton decay:

Decay mode	Experiment		
	NUSEX [5.9]	IMB [5.10]	KAMIOKANDE [5.11]
$p \to e^+ \pi^0$	1.3	25	5.1
$p \to e^+ \eta$	-	20	5.1
$p \to e^+ K^0$	-	7.7	3.0
$p \to e^+ \rho^0$	-	1.7	1.0
$p \to e^+ \omega$	-	3.7	2.0
$p \to e^+ K^{*0}$	-	-	0.8
$p \to \mu^+ \pi^0$	1.0	7.6	3.8
$p \to \mu^+ \eta$	-	4.6	2.1
$p \to \mu^+ K^0$	0.8	4.0	1.1
$p \to \mu^+ \rho^0$	-	1.6	0.6
$p \to \mu^+ \omega$	-	2.3	-
$p \to \nu^c \pi^+$	0.3	-	0.4
$p \to \nu^c K^+$	0.6	1.0	1.5
$p \to \nu^c \rho^+$	-	0.8	0.9
$p \to \nu^c K^{*+}$	-	1.0	1.7

(b) Bound neutron decay:

Decay mode	Experiment		
	NUSEX [5.9]	IMB [5.10]	KAMIOKANDE [5.11]
$n \to e^+ \pi^-$	2.1	3.1	2.6
$n \to e^+ \rho^-$	-	1.4	0.6
$n \to \mu^+ \pi^-$	0.6	2.3	2.0
$n \to \mu^+ \rho^-$	-	0.7	0.3
$n \to \nu^c \pi^0$	1.3	0.6	2.1
$n \to \nu^c \eta$	-	2.5	3.4
$n \to \nu^c K^0$	1.2	1.5	1.6
$n \to \nu^c \rho^0$	-	0.2	0.4
$n \to \nu^c \omega$	-	1.2	2.1
$n \to \nu^c K^{*0}$	-	0.5	0.4

Figure Captions

Fig. 1 (a) The Q^2-behaviour of the standard-model gauge coupling constants $\alpha_i(Q^2)$, i = 1,2,3, in great desert GUTs.

(b) The modification of this behaviour in supersymmetric GUTs.

Fig. 2 The interaction between three quarks and a lepton, generated by exchange of a generic superheavy, B and L violating gauge boson X.

Fig. 3 The components of the Yukawa coupling of a Higgs superfield to two matter superfields:

(a) ordinary Yukawa coupling $\overline{f^c} f\ H$,

(b) Higgsino-fermion-sfermion coupling $\overline{f^c \tilde{H}}(sf)$,

(c) two-sfermion-Higgs coupling (sf)(sf)H.

Fig. 4 The two-fermion-two-sfermion interaction, generated by exchange of

(a) a superheavy, B and L violating Higgs boson H_X, or

(b) its fermionic superpartner $\tilde{H}_X$.

Fig. 5 The effective dimension-five operator G_S(sf)(sf)(f)(f), obtained from the Higgs exchange diagrams of Fig. 4 in the limit $m_{H_X} \to \infty$.

Fig. 6 The effective four-fermion interaction G_S(f)(f)(f)(f), obtained from the dimension-five operator in Fig. 5 by gaugino dressing. The cross represents a Majorana mass of the gauge fermions.

Fig. 7 Contribution of $O_R^{(5)}$ to proton decay. The stars represent off-diagonal elements of the sfermion mass matrices, which mix left and right scalar fermions.

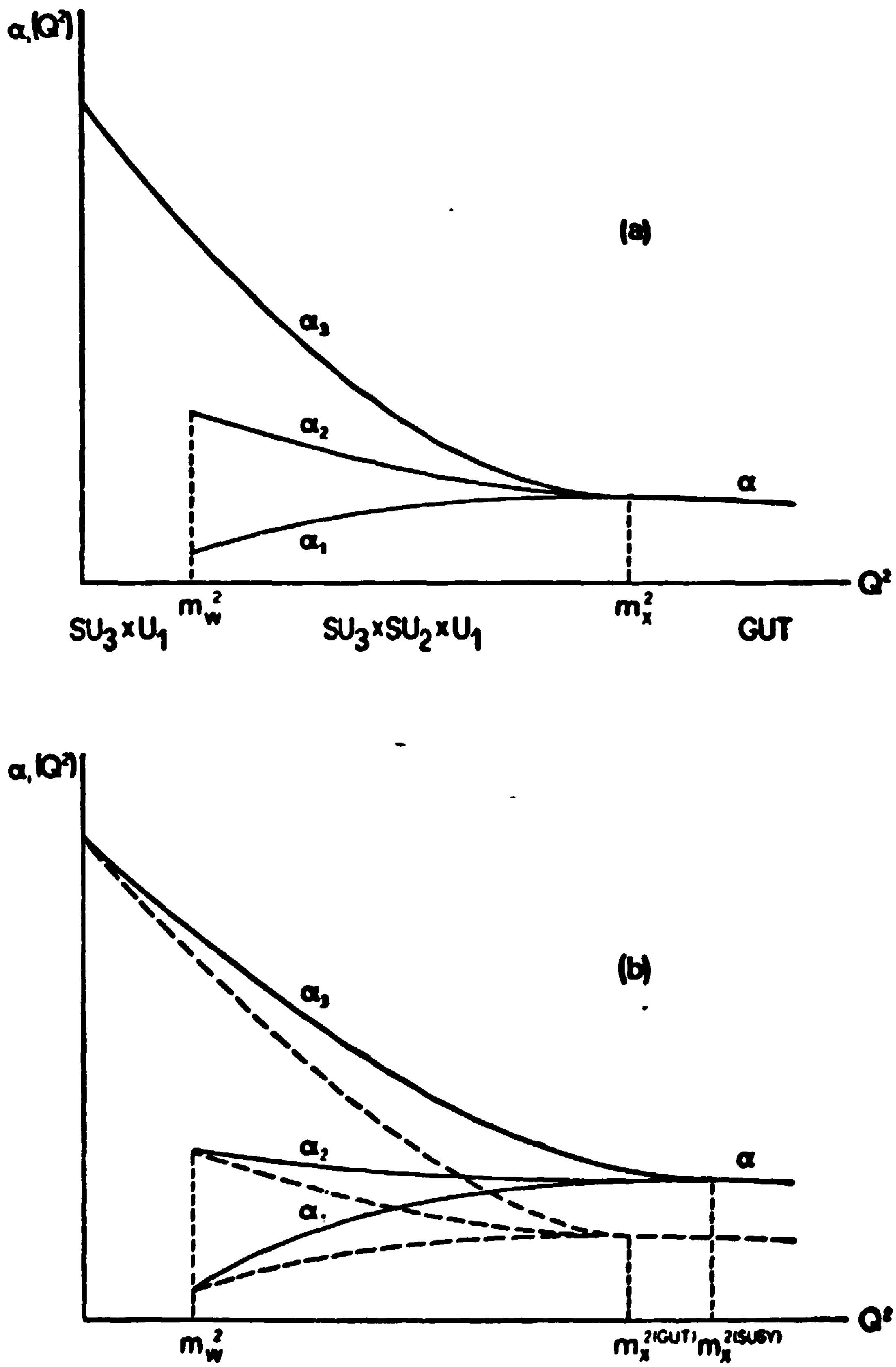
Fig. 8 The four-fermion operator $G(q\ q\ q\ \ell)$ responsible for proton decay.

Fig. 9 The three quark-level mechanisms by which the effective four-fermion interaction in Fig. 8 can induce proton decay:

(a) Two-quark annihilation: $q + q \to q^c + \ell^c$,

(b) Three-quark fusion: $q + q + q \to \ell^c$,

(c) Quark decay: $q \to q^c + q^c + \ell^c$.



αᵢ(Q²)
(a)
α₃
α₂
α
α₁
Q²
m_w²
m_x²
SU₃×U₁
SU₃×SU₂×U₁
GUT
αᵢ(Q²)
(b)
α₃
α₂
α
α₁
Q²
m_w²
m_x²(GUT)
m_x²(SUSY)

Fig. 1

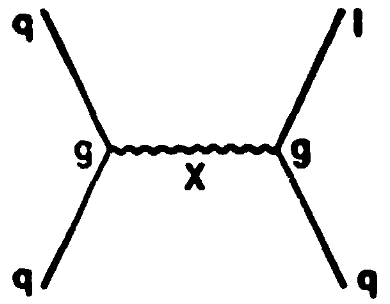


Fig. 2

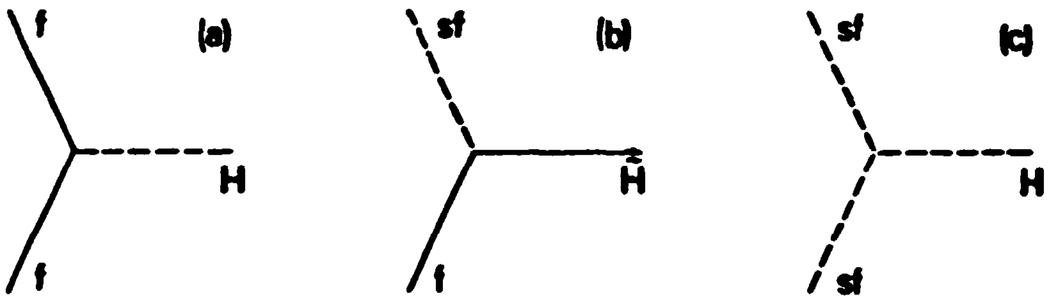


Fig. 3

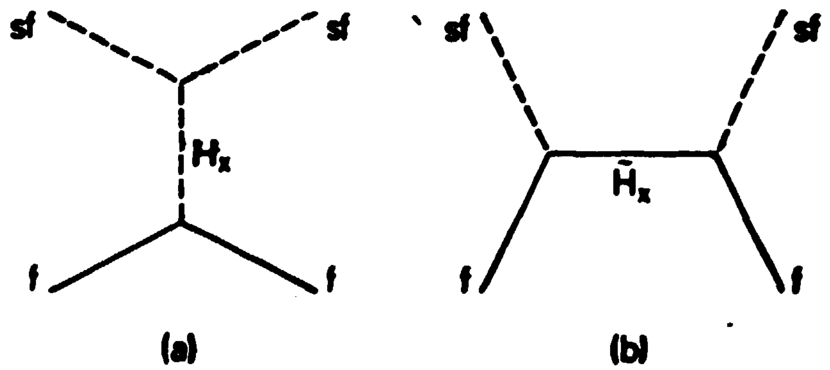


Fig. 4

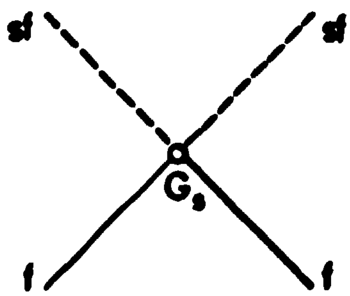


Fig. 5

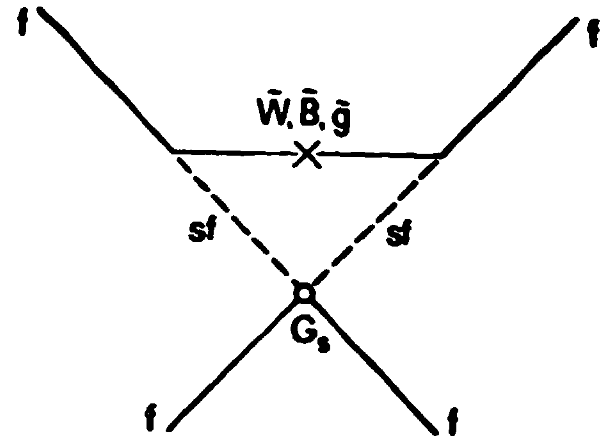


Fig. 6

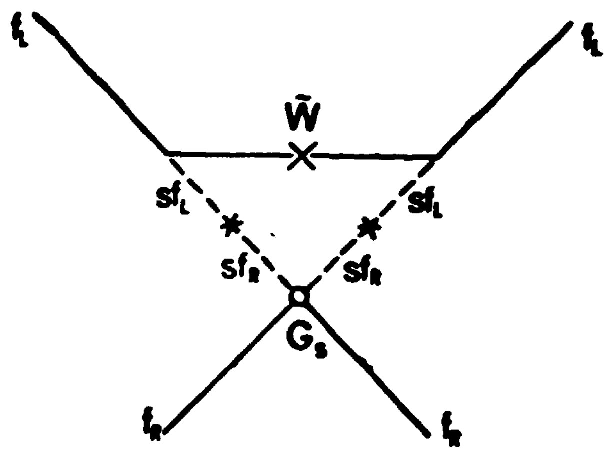


Fig. 7

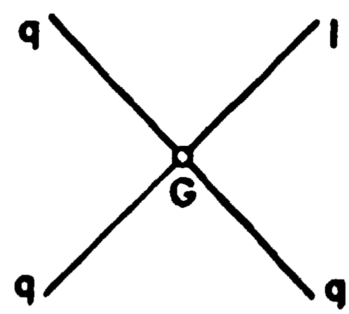


Fig. 8

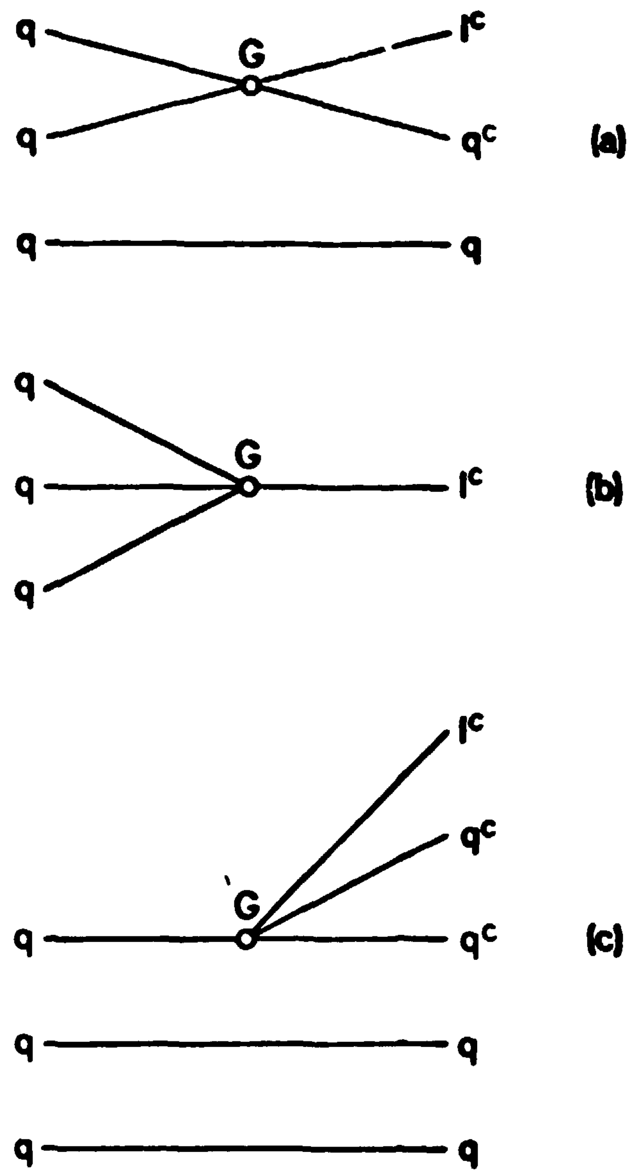


Fig. 9