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PROBINGS THROUGH PROTON DECAY AND $n-\bar{n}$ OSCILLATIONS *

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ABSTRACT

Violation of baryon, lepton and in general fermion number is central to the hypothesis of quark lepton unification in a gauge context. Three of its characteristic signatures are proton decay, $n-\bar{n}$ oscillation and neutrinoless double β decay. In 1974 and 1975 it was shown that within maximal gauging the proton may decay via four alternative modes (i.e. proton $\to$ one or three leptons or anti-leptons) satisfying $\Delta F = -2, 0, -4$ and -6 some of which may coexist; the deuteron may decay into pions and neutrinoless double β decay occur in the context of spontaneous gauge symmetry breaking. It is now observed that $n-\bar{n}$ oscillations (which are related to deuteron decays into protons) can coexist with proton decay especially of $\Delta F = -4$ variety ($p \to e^+\pi^0$) and both these processes may possess measurable strength so as to be amenable to forthcoming searches. We exhibit alternative routes for spontaneous breakdown of the maximal one-family symmetry SU(16) and show that the coexistence of alternative proton decay modes (even with $n-\bar{n}$ oscillations) does not pose any conflict with cosmological generation of baryon excess. Spontaneous rather than explicit violation of B, L and F plays an essential role in the realization of these features.

I. INTRODUCTION

We believe that in the context of gauge unification of particles and their forces [1)-4)], quantum numbers like baryon number, B, lepton number, L and fermion number F = 3B+L are best defined in the basic Lagrangian as parts of a non-abelian local symmetry G just like electric charge. In keeping with this point of view, we wish to examine in this note the consequences for proton decay of maximal gauging [4)] of quark-lepton unifying symmetries.

One characteristic feature of "maximal" gauging of symmetries is to ensure that baryon, lepton and fermion numbers are exact symmetries of the basic gauge Lagrangian; they are violated spontaneously when gauge particles acquire their masses. By contrast, for the case of non-maximal gauging, symmetries like B,L and F are violated, in general, explicitly.

The interest in maximal gauging of such symmetries stems from the fact that as a rule such gauging permits several intermediate mass scales [5)-7)] lying within the "grand plateau" between 10^2 and 10^{15} GeV. [8)] Consequently, it turns out that the proton may decay via four different modes [*)]

$$
\begin{aligned}
&(i) && p \to 3\ell + \text{mesons } (\Delta F = 0),\\
&(ii) && p \to \ell + \text{mesons } (\Delta F = -2),\\
&(iii) && p \to \bar{\ell} + \text{mesons } (\Delta F = -4) \text{ and}\\
&(iv) && p \to 3\bar{\ell} + \text{mesons } (\Delta F = -6) \quad .
\end{aligned}
\tag{1}
$$

While this had been shown in our earlier papers [4)] of 1975, here we show that not only can some of these decay modes coexist, but that they may also coexist with $\Delta B = 2$, $n-\bar{n}$ oscillations [9)] and possibly also with $\Delta L = 2$ neutrinoless double β decay with measurable strength. We show that such coexistence can occur without any conflict with the generation of baryon excess. [**)] A search for these effects would provide important clues on the possible existence of intermediate mass scales and thereby on the underlying, maximally gauged, design of grand unification.

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*) That proton may decay via these four alternative modes was first noted in Ref.4, where the maximal two family symmetry SU(32) was gauged.

**) A preliminary report of the results of this paper was given in Ref.7.

II. MAXIMAL GAUGED SYMMETRY SU(16) FOR ONE FAMILY AND ITS SPONTANEOUS DESCENT

Let us specify what we mean by "maximal" gauged symmetry. This corresponds to gauging all fermionic degrees of freedom [4] with fermions consisting of quarks <u>and</u> leptons. Thus with n two component left-handed fermions F_L plus n two component right-handed fermions F_R (which may be replaced by the left-handed charge conjugate fields $(F^c)_L$), the maximal symmetry is SU(2n). As an example, for a single family of eight left-handed fermions (six quarks and two leptons) plus their antiparticles, the maximal symmetry is SU(16). One word of qualification is in order. Such symmetries generate triangle anomalies, which are avoided by postulating that there exists a conjugate mirror set of fermions *) $F^m_{L,R}$ which couples to the gauge mesons through the helicity flip coupling $(F_{L,R} \leftrightarrow F^m_{R,L})$. Thus by "maximal" symmetries, we shall mean symmetries which are maximal upto the discrete mirror symmetry.

For maximal gauging of the three families (e, μ and τ) one would need to gauge [7),6)] $[SU(16)]^3$ or the still extended symmetry SU(48). These symmetries are no doubt gigantic, but if the quarks and leptons are proliferated, why not the associated gauge particles? We believe the real answer to proliferation must come from viewing quarks, leptons <u>and</u> also the associated gauge and Higgs particles as composites of more elementary objects – preons.[10),11)] From this point of view, extended maximal symmetries such as SU(48) or $[SU(16)]^3$ are only effective gauge symmetries generated from a much simpler and more economical basis of preons [12]. Spontaneous symmetry breaking can permit the descent of SU(48) or $[SU(16)]^3$ to the familiar low-energy symmetry $SU(2)_L \times U(1) \times SU(3)_C$ (via for example the diagonal symmetry $SU(16)_{e+\mu+\tau}$) in such a way that the interfamily universality ($e \leftrightarrow \mu \leftrightarrow \tau$) appears only below an energy scale of $\sim 10^5$ GeV. As shown elsewhere,[6),7)] such extended maximal symmetries permit signals for grand unification at low and intermediate mass scales ($\sim 10^4$ to 10^5 GeV and 10^8–10^{10} GeV) and thereby offer richer experimental possibilities than for example SU(5) or SO(10). [13] In what follows we shall (for simplicity) use SU(16), in much of our discussion, <u>as a language</u> for maximal symmetries **), though we shall ultimately view it as part of an extended maximal symmetry (such as SU(48) or $[SU(16)]^3$).

*) It is possible to avoid anomalies by introducing multiplets other than conjugate multiplets. An example is $\bar{5} + 10$ for SU(5) (Ref.3). However in such a theory no generator of the local symmetry can be associated with a linear combination of B, L or F, so that their breaking will necessarily be of the explicit variety.

**) For the diagonally summed SU(16) → $B = B_e + B_\mu + B_\tau$, $L = L_e + L_\mu + L_\tau$ and $F = F_e + F_\mu + F_\tau$.

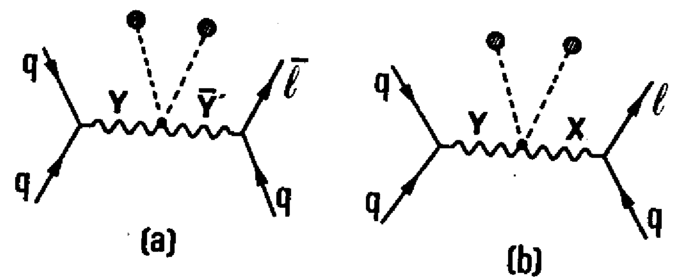
One immediate consequence of maximal gauging [4] is that a linear combination of baryon and lepton numbers (B-L for *) SU(16)) and, in addition the fermion number $F \equiv B_q + L = 3B + L$, get locally gauged; they are among the generators of G . (Here B_q denotes quark number, which is +1 for quarks and -1 for antiquarks; familiar baryon number B, which is +1 for proton, is $B_q/3$). As a consequence, B, L and F are conserved in the basic Lagrangian. They are violated spontaneously and unavoidably when the associated gauge particles acquire mass. The violations (1) come about for example through spontaneously <u>induced mixings</u> between gauge particles carrying different B, L and F quantum numbers. Specifically the mixing of an F = +2 diquark gauge particle Y coupled to $\bar{q}^c\gamma_\mu q$ current with a F = -2 quark-lepton gauge particle $\bar{Y}'$ coupled to the $\bar{q}\gamma_\mu \ell^c$ current (See Fig.1(a)) induces ΔB = ΔL (i.e. ΔF = -4) proton decay:

$$p \to \bar{\ell} + \text{mesons} \quad . \tag{2}$$

Likewise the spontaneously induced mixing of the F = 2 diquark gauge particle Y with an F = 0 leptoquark gauge particle X coupled to the $\bar{q}\gamma_\mu \ell$ current (see Fig.1(b)) gives rise to the ΔB = - ΔL (ΔF = -2) proton decay:

$$p \to \ell + \text{mesons} \quad . \tag{3}$$

The mechanism for spontaneous symmetry breaking leading to these and other B, L and F violating processes is elaborated in the next section.



<u>Fig.1</u>: Spontaneously induced proton decays satisfying ΔF = -4 and -2.

*) Note that SU(16) (as also SO(10)) contains $SU(4)_{colour}$ as a subgroup, with lepton number as the fourth colour. All such symmetries therefore contain $B_q - 3L = 3(B-L)$ as a generator, where B_q denotes quark number ($B_q \equiv 3B$).

This spontaneous violation of B, L and F should be distinguished from the case where the violations are intrinsic,that is to say explicit already in the basic gauge Lagrangian. For example for SU(5), neither B-L nor F is locally gauged and B, L as well as F are violated explicitly in the basic gauge interaction *). For SO(10) **) B-L is locally gauged, but not F and here too B, L and F are violated in the basic gauge Lagrangian. As explained elsewhere in detail [7], such explicit violations arise only provided one chooses to "squeeze" the gauges of the "maximal" symmetry (like SU(16)) such that one and the same gauge particle couples, for example, to the diquark $(\bar{q}^c\gamma_\mu q)$ as well as to the quark-lepton $(\bar{q}\gamma_\mu \ell^c)$ currents in the gauge Lagrangian. In other words, B, L and F are violated (explicitly if one chooses to gauge specific subgroups (like SU(5) or SO(10)) of the "maximal" group. The two cases of spontaneous versus explicit violations of B, L and F differ from each other conceptually as well as in their physical consequences. Spontaneous violation would in general disappear at high temperatures exceeding the masses of the relevant gauge particles, while explicit violation would acquire its maximum strength at such temperatures where the gauge particles would be massless.[14] This distinction would play its most obvious role in the early stage of the Universe as we indicate later.

It is now of interest to see the alternative routes for the spontaneous descent of SU(16) down to the low-energy symmetry $SU(2)_L \times U(1) \times SU(3)_C$. The three most obvious routes are via 1) SO(10) with respect to which the fundamental fermionic 16-plet remains irreducible, 2) the maximal chiral route $SU(8)_I \times SU(8)_{II} \times U(1)_F$, where the two SU(8)'s operate in the spaces of the eight fermions and the eight antifermions, respectively, and $U(1)_F$ represents fermion number and 3) the route through $SU(12)_q \times SU(4) \times U(1)_{|B|-|L|}$ where $SU(12)_q$ operates on six quarks plus six antiquarks, $SU(4)_\ell$ on two leptons plus two antileptons and U(1) corresponds to the symmetry $|B_q| - 3|L| \equiv 3(|B|-|L|)$. These three alternative routes are exhibited in Fig.2. Note that unlike the first two, the third route separates quarks from leptons at the very first stage of the spontaneous symmetry breaking. Low energy selection rules for B, L and F violations and proton decay as well as the strength of n-$\bar{n}$ oscillation would depend upon the route for spontaneous descent Nature chooses. Our attitude is that it would be premature at present to speculate on which of these routes if any is preferred. We exhibit below the Higgs system, which may permit descent via these alternative routes.

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*) For the minimal SU(5) model (with 5 and 24 of Higgs), while B,L and F are violated, B-L turns out to be a global symmetry of the basic Lagrangian.

**) This is because SO(10) contains $SU(4)_{colour}$ (Ref.2) as a subgroup.

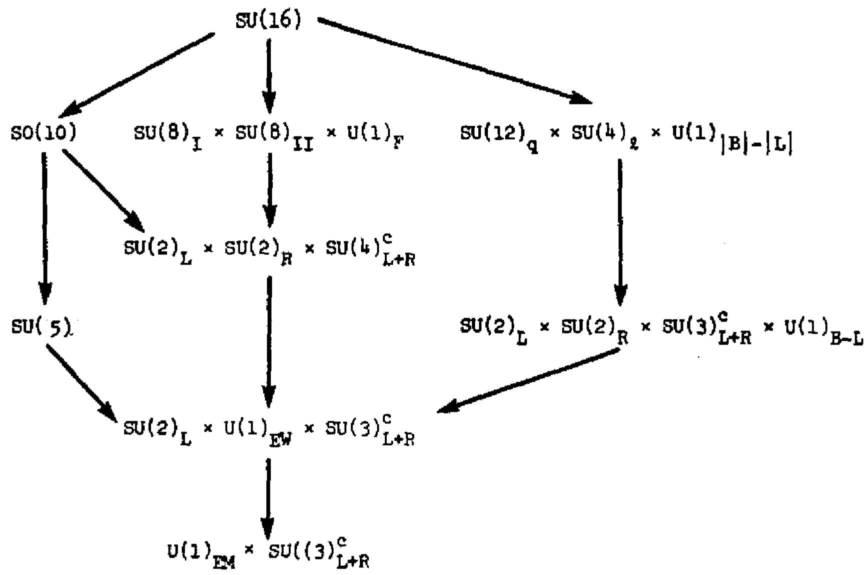

Fig.2: Alternative routes for spontaneous descent of SU(16). See text for definitions of the subsymmetries.

III. THE HIGGS SYSTEM

Even though we believe that symmetry violations have a dynamical role (perhaps in the fermion pairing forces), we shall employ the usual Higgs field mnemonic: local scalar operators multilinear in the fermion fields will be represented by independent scalar fields. Their non-vanishing vacuum expectation values, which are governed by an effective potential, determine the pattern of spontaneous symmetry breaking. Needless to say, we shall not be in a position to solve for the true minimum of the full effective potential, which in general involves the interplay of several scalar multiplets and thus a multitude of unknown mass and scalar quartic coupling parameters. Nor are we yet in a position to solve the gauge hierarchy problem [15]. We belive that a solution to both these problems, which are related,would arise by taking recourse to ideas similar to those of technicolour [16] (including grand unification), but where not only the Higgs, but also the relevant fermions and gauge particles are composites of preons or pre-preons [10)-12)]. In this note we shall content ourselves simply to spelling out alternative patterns of vacuum expectation values leading to distinctive experimental consequences. A choice between these patterns may soon be provided by ongoing experiments searching for proton decay [17)] and n-$\bar{n}$ oscillation [18)].

Let us denote the fermion 16-plet $\{u_r, u_y, u_b, \nu_e; d_r, d_y, d_b, e^- | d_r^c, d_y^c, d_b^c, e^+; u_r^c, u_y^c, u_b^c, {}^c_e\nu_L\}$ by ψ_A and its conjugate by $\bar{\psi}^A$, where c denotes charge conjugation and A = 1,2,...,16. Assuming that the Higgs are effectively composites of even numbers of fermions (including antifermions), we are led to consider Higgs fields such as

$$\begin{aligned} \psi^T_{(A} \; C^{-1} \; \psi_{B)} &\sim \phi_{(AB)} \\ \bar{\psi}^B \; \psi_A &\sim \phi^B_A \\ \bar{\psi}^{(C} \; \bar{\psi}^{D)} \; \psi_{(A} \; \psi_{B)} &\sim \phi^{(CD)}_{(AB)} \quad , \quad \text{etc.} \end{aligned} \tag{4}$$

These include the symmetric second rank tensorial representation $\phi_{(AB)} = \underline{136}$, the adjoint representation $\phi^B_A = \underline{255}$ and the fourth rank tensorial representation symmetric in two upper and symmetric in two lower indices $\phi^{(AB)}_{(CD)} = \underline{18,240}$ of SU(16). The last two multiplets provide the alternative chains of SU(16)-breaking, while the VEV of $\underline{136}$ breaks SU(2) × SU(2) × SU(4) in addition to providing masses for the fermions. It is helpful to decompose SU(16) representations in terms of its subgroups*) $SU(8)_I \times SU(8)_{II} \times U(1)_F$ and $SU(2)_L \times SU(2)_R \times SU(4)^c_{L+R}$ corresponding to the chain

$$\begin{aligned} SU(16) &\to SU(8)_I \times SU(8)_{II} \times U(1)_F \\ &\to SU(2)_L \times SU(2)_R \times SU(4)^c_{L,R} \; . \end{aligned} \tag{5}$$

As mentioned before, $SU(8)_I$ and $SU(8)_{II}$ operate in the spaces of eight left-handed fermions and eight left-handed antifermions, respectively, while $U(1)_F$ denotes fermion number.

Under $SU(8)_I \times SU(8)_{II} \times U(1)_F$, the 16-component objects ψ_A and $\bar{\psi}_A$ decompose as $\underline{16} = (8,1)_1 + (1,\bar{8})_{-1}$ and $\overline{16} = (\bar{8},1)_{-1} + (1,8)_{+1}$, respectively, while under $SU(2)_L \times SU(2)_R \times SU(4)_{L+R} \times U(1)_F$ they decompose as:

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*) The spontaneous breakdown of SU(8) × SU(8) subgroup (<u>without</u> the fermion number symmetry $U(1)_F$) has been considered in a different context in Ref.19.

$$\psi_A = \begin{pmatrix} \psi_{a\alpha} \\ \psi_{\bar{a}\bar{\alpha}} \end{pmatrix} = \left[\frac{(2,1,4)_1}{(1,2,\bar{4})_{-1}} \right] \qquad \begin{aligned} a &= 1,2 \\ \alpha &= 1,\ldots,4 \end{aligned}$$

$$\bar{\psi}^A = \begin{pmatrix} \bar{\psi}^{a\alpha} \\ \bar{\psi}^{\bar{a}\bar{\alpha}} \end{pmatrix} = \left[\frac{(2,1,\bar{4})_{-1}}{(1,2,4)_{+1}} \right] \quad . \tag{6}$$

Here the subscripts ±1 denote fermion numbers. Both $\psi_{a\alpha}$ and $\psi_{\bar{a}\bar{\alpha}}$ are left-handed spinors, $\psi_{a\alpha}$ represent eight fermions in two flavours and four <u>colours</u>, while $\psi_{\bar{a}\bar{\alpha}}$ represents their antiparticles. Note that both $\psi_{a\alpha}$ and $\bar{\psi}^{\bar{a}\bar{\alpha}} = (\psi_{\bar{a}\bar{\alpha}})^\dagger \gamma^0$ transform in the same way under $SU(4)^c$. In short ψ^A is composed of $(f_L, (\bar{f})_L)$, while $\bar{\psi}^A$ is composed of $((\bar{f})_R, f_R)$, where $f_{L,R}$ are the octets of left-handed and right-handed fermions.

The 255 gauge particles of SU(16) decompose under $SU(8)_I \times SU(8)_{II} \times U(1)_F$ as follows:

$$\begin{aligned} 255 = {} & (1,1)_0 + (63,1)_0 + (1,63)_0 \\ & + (8,8)_{+2} + (\bar{8},\bar{8})_{-2} \quad , \end{aligned} \tag{7a}$$

while under $SU(2)_L \times SU(2)_R \times SU(4)^c \times U(1)_F$, the decomposition is as follows:

$$\begin{aligned} 255 = {} & (1,1,1)_0 + (3,1,1+15)_0 + (1,3,1+15)_0 \\ & + (1,1,15+15)_0 + (2,2,6+10)_2 + (2,2,6+\overline{10})_{-2} \quad . \end{aligned} \tag{7b}$$

The singlet $(1,1,1)_0$ which is also a singlet of $SU(8)_I \times SU(8)_{II}$ couples to the fermion number current. The remaining F = 0 gauge particles couple to the chiral currents of SU(8) × SU(8). The components with F = ±2 couple to fermion number changing currents $\bar{f}^c_L \gamma_\mu f_L$ and $\bar{f}_L \gamma_\mu f^c_L$. The components $(2,2,6+10)_2$ are hermitian conjugates of $(2,2,6+\overline{10})_{-2}$. The gauge particles may symbolically be represented by the following 16 × 16 matrix:

$$V = \left[\begin{array}{c|c} \begin{matrix} (63,1)_0 \\ +(1,1)_0 \end{matrix} & (8,8)_2 \\ \hline (\bar{8},\bar{8})_{-2} & \begin{matrix} (1,63)_0 \\ -(1,1)_0 \end{matrix} \end{array} \right] = \left[\begin{array}{l|l} W_L, V_L \text{ (gluons)} & Y \sim \bar{q}^c_L \gamma_\mu q_L \\ X_L \text{(leptoquark)} & Y' \sim \bar{\ell}^c_L \gamma_\mu q_L , \; \bar{q}^c_L \gamma_\mu \ell_L \\ (V_{B-L}, V_F)_L & Y'' \sim \bar{\ell}^c_L \gamma_\mu \ell_L \\ \hline \bar{Y}, \bar{Y}', \bar{Y}'' & L \to R \end{array} \right] \tag{8}$$

[It is worth noticing that the gauge particles of the subgroup SO(10) include the 24 real symmetric combinations "Y_s" $\equiv (1/\sqrt{2})\,\{(2,2,6)_2 + (2,2,6)_{-2}\}$ but not the corresponding antisymmetric combinations "Y_a" $\equiv 1/\sqrt{2}\,\{(2,2,6)_{-2} - (2,2,6)_{-2}\}$. Since any member belonging to the set "Y_s" couples to a mixture of $F = +2$ and $F = -2$ currents, its exchange leads to an explicit violation of fermion number as well as of baryon and/or lepton numbers for SO(10).]

Of the Higgs multiplets mentioned earlier (see (4)), the <u>136</u> decomposes as

$$\underline{136} = (3,1,10)_{+2} + (1,3,\overline{10})_{-2} + (1,1,6)_{+2} + (1,1,6)_{-2} + (2,2,1+15)_0 \,. \tag{9}$$

Thus it contains no SU(2) × SU(2) × SU(4) singlet. <u>Hence at the level where $SU(2)_L \times SU(2)_R \times SU(4)^c$ is preserved</u>, the <u>136</u> <u>should develop no VEV's</u>. The <u>255</u> possesses one singlet (see (8)); thus

$$\left\langle \phi^{b\beta}_{a\alpha} \right\rangle = \mathcal{V}_0\, \delta^b_a\, \delta^\beta_\alpha \;, \quad \left\langle \phi^{\bar b\bar\beta}_{\bar a\bar\alpha} \right\rangle = -\mathcal{V}_0\, \delta^{\bar b}_{\bar a} - \delta^{\bar\beta}_{\bar\alpha}$$
$$\left\langle \phi^{\bar b\bar\beta}_{a\alpha} \right\rangle = \left\langle \phi^{b\beta}_{\bar a\bar\alpha} \right\rangle = 0 \quad . \tag{10}$$

Such a pattern breaks SU(16) into $SU(8) \times SU(8) \times U(1)_F$. Note that VEV of <u>255</u> preserves fermion number F as also B and L. Thus $\mathcal{V}_0$ gives masses only to those gauge particles which carry fermion number ±2. These are the $(8,8)_{+2}$ and $(\bar 8,\bar 8)_{-2}$ appearing in Eq.(7). In the notation of earlier papers $(8,8)_{+2}$ contains the diquark (Y), the lepto-antiquark (Y') and the dilepton (Y") gauge particles coupled, respectively, to the currents $\bar q^c_L \gamma_\mu q_L$, $\bar\ell^c_L \gamma_\mu q_L$ and $\bar\ell^c_L \gamma_\mu \ell_L$. The $(\bar 8,\bar 8)_{-2}$ are the hermitian conjugates of $(8,8)_{+2}$. (At the lower stages of the hierarchy of SSB, components of <u>136</u> as well as <u>additional</u> components of <u>255</u> may acquire non-zero VEV with magnitudes much smaller than that of $\mathcal{V}_0$. These would break $SU(2)_L \times SU(2)_R \times SU(4)^c_{L+R}$.)

Vacuum expectation values of the 4th rank tensorial field $\phi^{\{AB\}}_{\{CD\}}$ can take SU(16) into $SU(2)_L \times SU(2)_R \times SU(4)_{L+R}$ violating F, B and L. For special combinations of these, the SU(16) may break down to SO(10). To demonstrate this we first note that $\phi^{\{AB\}}_{\{CD\}}$ possesses <u>six independent</u> SU(2) × SU(2) × SU(4) singlets. This may seen as follows: Under $SU(8) \times SU(8) \times U_F(1)$, this multiplet decomposes as:

$$\begin{aligned}\phi^{\{AB\}}_{\{CD\}} = {} & (36,36)_4 + (\overline{36},\overline{36})_{-4} \\ & + (280,8)_2 + (8,280)_2 + (8,8)_2 \\ & + (\overline{280},\bar 8)_{-2} + (\bar 8,\overline{280})_{-2} + (\bar 8,\bar 8)_{-2} \\ & + (1232,1)_0 + (1,1232)_0 + (1,1)_0 + (1,63)_0 + (63,1)_0 \\ & + (63,63)_0 \quad .\end{aligned} \tag{11}$$

Under $SU(2)_L \times SU(2)_R \times SU(4)^c_{L+R} \times U(1)_F$ the above multiplets further decompose into

$$\begin{aligned}(36,36)_4 &= (1,1,1+15+20')_4 + (3,1,15+45)_4 + (1,3,15+\overline{45})_4 \\ &\quad + (3,3,1+15+84)_4 \\ (1232,1)_0 &= (1,1,1+15+20'+84)_0 + (3,1,15+15+45+\overline{45}+84)_0 + (5,1,1+15+84)_0 \\ (8,8)_2 &= (2,2,6+10)_2 \\ (8,280)_2 &= (2,2(4\times 20''))_2 + (2,2+4(4\times 4)+(4\times 36))_2 \\ (63,63)_0 &= (1,1(15\times 15))_0 + (3,1,15+(15\times 15))_0 + (1,3,15+(15\times 15))_0 \\ &\quad + (3,3,1+15+15+(15\times 15))_0 \quad .\end{aligned} \tag{12}$$

Thus the invariants of SU(2) × SU(2) × SU(4) in $\phi^{\{AB\}}_{\{CD\}}$ lie in the six sub-multiplets of SU(8) × SU(8) × U(1); viz: $(1,1232)_0$, $(1232,1)_0$, $(1,1)_0$, $(63,63)_0$, $(36,36)_4$ and $(\overline{36},\overline{36})_{-4}$. Of these (36,36) and $(\overline{36},\overline{36})$ are hermitian conjugates of each other. Accordingly, at the level that SU(2) × SU(2) × SU(4) is preserved, the VEV of $\phi^{\{AB\}}_{\{CD\}}$ may be represented by five parameters:

$$\begin{aligned}\left\langle \phi^{c\gamma\;\; d\delta}_{a\alpha\;\; b\beta} \right\rangle &= (\mathcal{V}_{1,1} + 5\mathcal{V}_{1232,1})\,(\delta^c_a\delta^d_b - \delta^d_a\delta^c_b)\,(\delta^\gamma_\alpha\delta^\delta_\beta - \delta^\delta_\alpha\delta^\gamma_\beta) \\ &\quad + (\mathcal{V}_{1,1} - \mathcal{V}_{1232,1})\,(\delta^c_a\delta^d_b + \delta^d_a\delta^c_b)\,(\delta^\gamma_\alpha\delta^\delta_\beta + \delta^\delta_\alpha\delta^\gamma_\beta) \\ \left\langle \phi^{c\gamma\;\; d\delta}_{a\alpha\;\; b\beta} \right\rangle &= -\tfrac{1}{4}\,(9\mathcal{V}_{1,1} + \mathcal{V}_{63,63})\,\delta^c_a\,\delta^{\bar d}_{\bar b}\,\delta^\gamma_\alpha\,\delta^\beta_\delta + \mathcal{V}_{63,63}\,\delta^c_a\,\delta^{\bar d}_{\bar b}\,\delta^\beta_\alpha\,\delta^\gamma_\delta \\ \left\langle \phi^{\bar c\bar\gamma\;\; \bar d\bar\delta}_{\bar a\bar\alpha\;\; \bar b\bar\beta} \right\rangle &= (\mathcal{V}_{1,1} + 5\mathcal{V}_{1,1232})\,(\delta^{\bar c}_{\bar a}\delta^{\bar d}_{\bar b} - \delta^{\bar d}_{\bar a}\delta^{\bar c}_{\bar b})\,(\delta^\alpha_\gamma\delta^\beta_\delta - \delta^\beta_\gamma\delta^\alpha_\delta) \\ &\quad + (\mathcal{V}_{1,1} - \mathcal{V}_{1,1232})\,(\delta^{\bar c}_{\bar a}\delta^{\bar d}_{\bar b} + \delta^{\bar d}_{\bar a}\delta^{\bar c}_{\bar b})\,(\delta^\alpha_\gamma\delta^\beta_\delta + \delta^\beta_\gamma\delta^\alpha_\delta) \\ \left\langle \phi^{\bar c\bar\gamma\;\; \bar d\bar\delta}_{a\alpha\;\; b\beta} \right\rangle &= \left\langle \phi^{a\alpha\;\; b\beta}_{\bar c\bar\gamma\;\; \bar d\bar\delta} \right\rangle^* = \varepsilon_{ab}\,\epsilon^{\bar c\bar d}\,\epsilon_{\alpha\beta\gamma\delta}\,\mathcal{V}_{36,36} \quad .\end{aligned} \tag{13}$$

Of these, note that only $\vartheta_{36,36}$, corresponding to the VEV of the submultiplets $(36,36)_4$ and $(\overline{36},\overline{36})_{-4}$ violates fermion number by ±4 units, characteristic of SO(10).

It is a straightforward though tedious task to evaluate the contributions of these Higgs fields to the vector mass terms. Denote the components of the gauge field by $W_{\mu A}{}^{B}$. They enter the covariant derivatives in the form

$$\nabla_\mu \phi_A^B = \partial_\mu \phi_A^B - ig\,(W_{\mu A}{}^{C}\,\phi_C^B - W_{\mu C}{}^{B}\,\phi_A^C)\,, \tag{14}$$

$$\nabla_\mu \phi_{AB}^{CD} = \partial_\mu \phi_{AB}^{CD} - ig(W_{\mu A}{}^{E}\,\phi_{EB}^{CD} + W_{\mu B}{}^{E}\,\phi_{AE}^{CD} - W_{\mu E}{}^{C}\,\phi_{AB}^{ED} - W_{\mu E}{}^{D}\,\phi_{AB}^{CE}). \tag{15}$$

It follows that the vector mass matrix is contained in the epxression

$$\begin{aligned}&\tfrac{1}{2}\left(\nabla_\mu\langle\phi_A^B\rangle\right)\left(\nabla_\mu\langle\phi_B^A\rangle\right) + \tfrac{1}{4}\left(\nabla_\mu\langle\phi_{AB}^{CD}\rangle\right)\left(\nabla_\mu\langle\phi_{CD}^{AB}\rangle\right)\\ &= -g^2\Big[\langle\phi_B^C\rangle\langle\phi_D^A\rangle - \tfrac{1}{2}\delta_B^C\langle\phi_D^B\rangle\langle\phi_B^A\rangle - \tfrac{1}{2}\delta_D^A\langle\phi_B^D\rangle\langle\phi_D^C\rangle\\ &\quad + 2\langle\phi_{BR}^{CS}\rangle\langle\phi_{DS}^{AR}\rangle - \langle\phi_{BD}^{RS}\rangle\langle\phi_{RS}^{AC}\rangle\\ &\quad - \tfrac{1}{2}\delta_B^C\langle\phi_{DR}^{ST}\rangle\langle\phi_{ST}^{AR}\rangle - \tfrac{1}{2}\delta_D^A\langle\phi_{BR}^{ST}\rangle\langle\phi_{ST}^{CR}\rangle\Big]\,W_A^B\,W_C^D\,.\end{aligned} \tag{16}$$

To study the structure of this, first assume that the SU(8) × SU(8) singlet $\vartheta_{(1,1)}$ dominates over all the other non-singlet VEV parameters of $\langle\phi_{\{CD\}}^{\{AB\}}\rangle$, i.e. $\vartheta_{(1,1)} \gg \vartheta_{(1232,1)}, \vartheta_{(1,1232)}, \vartheta_{(63,63)}$ and $\vartheta_{(36,36)}$. The singlet parameters $\vartheta_{(1,1)}$ will contribute only to the masses of the F = ±2 gauge particles Y, Y', Y" (see Eq.(8)) belonging to (8,8) and its conjugate, which in addition receive contribution from ϑ_0 (the VEV of 255, see Eq.(10)). Their net masses (in the approximation that SU(8) × SU(8) × $U(1)_F$ is preserved) are given by

$$m^2_{Y,Y',Y''} \approx g^2\,(\vartheta^2_{(1,1)} + \vartheta^2_0) \tag{17}$$

which multiplies $W_{a\alpha}^{\bar b\bar\beta}\,W_{\bar b\bar\beta}^{a\alpha}$.

This mass term will be split by the VEV of the non-singlet components of the scalar field and thereby Y, Y', Y" would receive differing masses. In general, such splitting is not of much interest, but the contributions of $\vartheta_{(36,36)}$ are distinguished by the fact that they yield fermion number changing mass mixing term

$$(\Delta m^2)\big|_{|\Delta F|=4} = g^2\,\vartheta_{(1,1)}\,\vartheta_{(36,36)}\,W_{\bar a\bar\alpha}{}^{b\beta}\,W_{\bar c\bar\gamma}{}^{d\delta}\,\epsilon_{\alpha\beta\gamma\delta}\,\epsilon_{bd}\,\epsilon^{\bar a\bar c} + \text{h.c.} \tag{18}$$

which induces Y-$\overline{Y}$' and Y"-$\overline{Y}$" ΔF = ±4 gauge mixings. *) The SU(8) × SU(8) singlet gauge vector V_F whose source is the fermion number current of course acquires a mass at this stage

$$M_F^2\,V_F^2 = g^2\,|\vartheta_{(36,36)}|^2\,(W_{a\alpha}{}^{a\alpha} - W_{\bar a\bar\alpha}{}^{\bar a\bar\alpha})^2\,. \tag{19}$$

Of the remaining 63 + 63 = 126 gauge particles belonging to SU(8) × SU(8), twenty-one, which gauge the subgroup H ≡ SU(2) × SU(2) × $SU(4)^c$, remain massless at this stage, while the rest 105 lying in the coset space SU(8) × SU(8)/H will acquire masses through $\vartheta_{(63,63)}$, $\vartheta_{(1232,1)}$, $\vartheta_{(1,1232)}$ and $\vartheta_{(36,36)}$. These 105 gauge particles will split into three distinct submultiplets $(1,1,15)_{L-R}$, (3,1,15) and (1,3,15) of the residual symmetry H. Their mass terms, under the simplifying assumption that $|\vartheta_{(36,36)}|$ is small compared to the other three parameters, are:

$$\begin{aligned}&g^2\,\vartheta^2_{(63,63)}\;A_\alpha{}^\beta\,A_\beta{}^\alpha \;\to\; (1,1,15)_{L-R}\\ &g^2\,(\vartheta^2_{(63,63)} + \vartheta^2_{(1232,1)})\;B_{La\alpha}{}^{b\beta}\,B_{Lb\beta}{}^{a\alpha} \;\to\; (3,1,15)\\ &g^2\,(\vartheta_{(63,63)} + \vartheta_{(1,1232)})\;B_{R\bar a\bar\alpha}{}^{\bar b\bar\beta}\,B_{R\bar b\bar\beta}{}^{\bar a\bar\alpha} \;\to\; (1,3,15)\,,\end{aligned} \tag{20}$$

where

$$\begin{aligned}A_\alpha{}^\beta &= W_{a\alpha}{}^{a\beta} - W_{\bar a\bar\beta}{}^{\bar a\bar\alpha} - \tfrac{1}{4}\,\delta_\alpha^\beta\,(W_{a\gamma}{}^{a\gamma} - W_{\bar a\bar\gamma}{}^{\bar a\bar\gamma})\\ B_{La\alpha}{}^{b\beta} &= W_{a\alpha}{}^{b\beta} - \tfrac{1}{2}\,\delta_a^b\,W_{c\alpha}{}^{c\beta} - \tfrac{1}{4}\,\delta_\alpha^\beta\,W_{a\gamma}{}^{b\gamma} + \tfrac{1}{8}\,\delta_a^b\,\delta_\alpha^\beta\,W_{c\gamma}{}^{c\gamma}\\ B_{R\bar a\bar\alpha}{}^{\bar b\bar\beta} &= W_{\bar a\bar\alpha}{}^{\bar b\bar\beta} - \tfrac{1}{2}\,\delta_{\bar a}^{\bar b}\,W_{\bar c\bar\alpha}{}^{\bar c\bar\beta} - \tfrac{1}{4}\,\delta_{\bar\alpha}^{\bar\beta}\,W_{\bar a\bar\gamma}{}^{\bar b\bar\gamma} + \tfrac{1}{8}\,\delta_{\bar a}^{\bar b}\,\delta_{\bar\alpha}^{\bar\beta}\,W_{\bar c\bar\gamma}{}^{\bar c\bar\gamma}\,.\end{aligned} \tag{21}$$

We stress that there is no a priori theoretical reason – hierarchical or otherwise – for the simplifying assumption made above regarding the relative smallness of $|\vartheta_{(36,36)}|$ and that in general $\vartheta(36,36)$ may even exceed the magnitudes of the parameters retained in (20). In this case SU(8) × SU(8) × $U(1)_F$ would descend to SU(2) × SU(2) × SU(4) through the single VEV parameter $\vartheta(36,36)$.

*) The Y-$\overline{Y}$' induces ΔF = −4 proton decays ($p \to e^+$ + mesons, etc.) as shown in Fig.1(a), while the Y'-$\overline{Y}$" induces ΔL = ±4 processes $\nu + \nu \to \bar\nu + \bar\nu$.

The next stage of symmetry breaking

$$SU(2)_L \times SU(2)_R \times SU(4)_{L+R} \rightarrow SU(2)_L \times U(1)_Y \times SU(3)^c \;,$$

where $U(1)_Y$ is defined by the sum $I_{3R} + \sqrt{2/3}\; F'_{15} = I_{3R} + (B-L)/2$ may be effected by the VEV of other components of the scalar multiplet $\phi^{\{AB\}}_{\{CD\}}$ transforming non-trivially under $SU(2)_R \times SU(4)_{L+R}$ (see (12)), as well as by the VEV of the scalar multiplet $\underline{136}$, which contains the submultiplets $(3,1,10)_{+2}$, $(1,3,\overline{10})_{-2}$, $(1,1,6)_{+2}$ + $(1,1,6)_{-2}$ + $(2,2,1+15)_0$.

For example, one might have

$$\left\langle \phi^{\overline{c\gamma}\;\overline{d\delta}}_{\overline{a\alpha}\;\overline{b\beta}} \right\rangle = \left(\delta^{\bar c}_{\bar a}\delta^{\bar d}_{\bar b} \pm \delta^{\bar d}_{\bar a}\delta^{\bar c}_{\bar b}\right)\left(\delta^{\bar\gamma}_{\bar\alpha}\delta^{\bar 4}_{\bar\beta}\delta^{\bar\delta}_{\bar 4} \pm \delta^{\bar\gamma}_{\bar\beta}\delta^{\bar 4}_{\bar\alpha}\delta^{\bar\delta}_{\bar 4} \pm \delta^{\bar\delta}_{\bar\alpha}\delta^{\bar 4}_{\bar\beta}\delta^{\bar\gamma}_{\bar 4} + \delta^{\bar\delta}_{\bar\beta}\delta^{\bar 4}_{\bar\alpha}\delta^{\bar\gamma}_{\bar 4}\right) u_{15\pm}$$

$$+ \left(\delta^c_a(\tau_3)^d_b \pm \delta^d_a(\tau_3)^c_b \pm \delta^c_b(\tau_3)^d_a + \delta^d_b(\tau_3)^c_a\right) \times$$

$$\times \;\left(\delta^\gamma_\alpha\,\delta^\delta_\beta \pm \delta^\delta_\alpha\,\delta^\gamma_\beta\right) u_{3\pm} \qquad (22)$$

with non-vanishing vacuum expectation values u_{15} and u_3 of the $SU(2)_L \times SU(2)_R \times SU(4)$ representations (1,1,15) and (1,3,1). The parameters u_{15} and u_3, respectively, give masses to $W^\pm_R$ of $SU(2)_R$ and to the leptoquark gauge particles (the X's of $SU(4)_{colour}$). An alternative attractive possibility is that the VEV of the multiplet $\underline{136}$ symbolically denoted by the pattern *)

$$\left\langle (1,3,\overline{10})_{-2} \right\rangle = v_R^{136} \gg 10^4\text{-}10^5 \text{ GeV}, \;\left\langle (3,1,\overline{10})_{-2} \right\rangle = 0, \;\left\langle (1,1,6)_{+2} \right\rangle = 0$$

$$\left\langle (1,1,6)_{-2} \right\rangle = 0, \;\left\langle (2,2,15) \right\rangle \ll v_R \qquad (23)$$

gives heavy masses $\gtrsim 10^4$-10^5 GeV to $W^\pm_R$ as well as to the leptoquark gauge particles X's and thereby breaks (through the VEV of a single submultiplet $(1,3,\overline{10})_{-2}$) left ⟷ right symmetry **), quark-lepton unification and fermion number simultaneously.

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*) In tensorial notation this corresponds to

$$\left\langle \phi_{\{AB\}} \right\rangle = \left\langle \phi_{\{\bar a\bar\alpha, \bar b\bar\beta\}} \right\rangle = (\tau^-)_{ab}\; \delta^4_\alpha\, \delta^4_\beta\, v_R^{136} \;.$$

**) An analogous suggestion has been made by Mohapatra and Marshak (Ref.20) in the context of the symmetry structure [2)] $SU(2) \times SU(2) \times SU(4)^c$.

Of particular interest are the leptoquark gauge particles

$$X_\alpha = W_{a\alpha}^{\;\;a4} - W_{\overline{a4}}^{\;\;\overline{a\alpha}} \;. \qquad (24)$$

These can induce *) (via a loop diagram involving $W_{L,R}\, X\, \phi \langle\phi\rangle$ gauge coupling (see Fig.3) or through a tree diagram (involving an $X_\mu \partial_\mu \phi \langle\phi\rangle$ coupling)) effective Yukawa transitions of the type:

$$(q_i)_{L,R} \rightarrow (\ell_j)_{R,L} + \phi_k \qquad (25)$$

with i,j,k taking appropriate group indices. These transitions have amplitudes proportional to $1/m_X^2$. Such effective transitions, taken in the third order, followed by quartic scalar coupling $\lambda\phi^4$, subject to one component of ϕ_ℓ having non-zero VEV, induce $\Delta F = 0$ proton decays to three leptons (see Fig.3)

$$p \sim 3q \rightarrow 3 \text{ leptons} + (\text{mesons}) \;(\Delta F = 0) \;. \qquad (26)$$

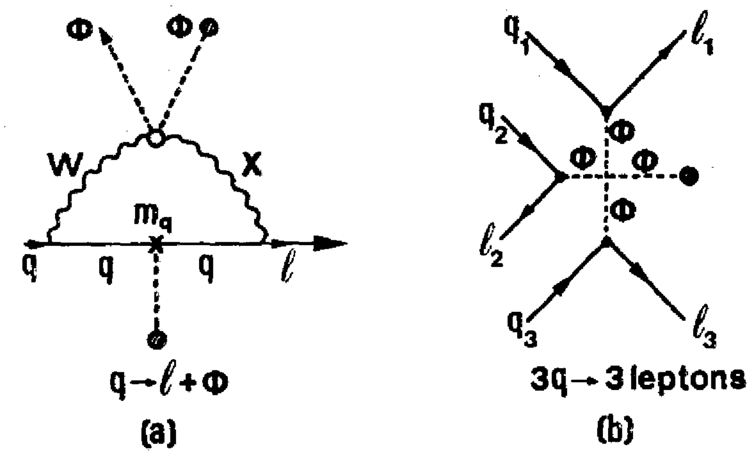

Fig.3: Spontaneously induced $q \rightarrow \ell + \phi$ and $3q \rightarrow 3\ell$ transitions.

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*) See Ref.21 for details.

We have discussed elsewhere that for reasonable values of the quartic coupling constant $\lambda \sim \mathcal{O}(\alpha)$ and for relevant Higgs masses $\sim$10-100 GeV, such amplitudes would lead to a partial proton decay rate $\sim (10^{31\pm 3}$ years$)^{-1}$ for a leptoquark mass m_X in the range of 10^4-10^5 GeV.

The leptoquarks X_α are also important in that their mixings with the superheavy diquark Y gauge particles induce $\Delta F = -2$ proton decays shown in Fig.1(b)

$$p \rightarrow (e^- \text{ or } \nu) + \text{mesons} \quad .$$

The diquarks Y coupling to $\bar{q}_L^c \gamma_\mu q_L$ are doublets of $SU(2)_L$, while the leptoquarks X_α belonging to $SU(4)_{colour}$ are singlets of $SU(2)_L$. Thus a mixing between them inevitably violates $SU(2)_L$ consistent with the results of Ref. 22 . Such a mixing can therefore occur only at the last stage of spontaneous symmetry breaking when $SU(2)_L \times U(1)$ breaks to $U(1)_{em}$. It could arise*) from the VEV of $|F| = 2$ components of $\Phi_{\{CD\}}^{\{AB\}}$ for example

$$\left\langle \Phi_{a\alpha\ b\beta}^{c\gamma\ \bar{d}\bar{\delta}} \right\rangle = \epsilon_{\alpha\beta\delta}\, \delta_4^\gamma\, \epsilon_{ab}\, \delta_1^c\, \delta_1^{\bar{d}}\, \omega \quad , \tag{27}$$

where it is understood that $\epsilon_{\alpha\beta\delta}$ vanishes unless α, β, δ is a permutation of 1,2,3. The magnitude of ω is necessarily limited to $O(10^2$ GeV) since this component contributes to the masses of the weak bosons W_L and Z. The mixing term it generates is

$$(\Delta m^2)_{|\Delta F|=2}\ g^2 \mathcal{V}_{(1,1)}\, \omega\, \left(W_{14}^{\ a\alpha} - W_{\bar{1}\bar{\alpha}}^{\ \bar{a}\bar{4}}\right) W_{\bar{1}\bar{\delta}}^{\ b\beta}\, \epsilon_{\alpha\beta\delta}\, \epsilon_{ab} + \text{h.c.} \tag{28}$$

We can now compare the orders of magnitude of the $\Delta F = -2$ and $\Delta F = -4$ transitions associated with this scheme. The leading contributions to the respective amplitudes would be

$$\text{Amp}(\Delta F = -2) \approx \frac{g^2 \mathcal{V}_{(1,1)}\omega}{M_X^2 M_Y^2} \sim \frac{M_W}{M_X^2 M_Y} \quad ,$$

$$\text{Amp}(\Delta F = -4) \approx \frac{g^2 \mathcal{V}_{(1,1)} \mathcal{V}_{(36,36)}}{M_Y^2 M_{Y'}^2} \quad , \tag{29}$$

where $M_Y \sim M_{Y'} \sim g(\mathcal{V}_{(1,1)} + \mathcal{V}_0)$ and $M_F \sim g\,\mathcal{V}_{(36,36)}$ denote typical masses of the vector diquark and the neutral vector V_F which couples to fermion number.

*) Alternatively, it may arise through VEV of an additional scalar multiplet such as $\Phi_{[CD]}^{[AB]} \sim [16 \times 16]_{antisym} \times [16^* \times 16^*]_{antisym}$.

It is clear that depending upon the magnitudes of $\mathcal{V}_0$, $\mathcal{V}_{(1,1)}$, $\mathcal{V}_{(36,36)}$ as well as the masses of the leptoquark gauge particles, the X's (which receive contributions from u_3, u_{15} and $\mathcal{V}_R^{136}$, see Eqs.(22) and (23)), several alternative scenarios regarding relative importance of $\Delta F = 0$ versus $\Delta F = -2$ versus $\Delta F = -4$ modes are a priori permissible. In particular three alternatives are worth noting:

$$\text{(A)} \quad \left.\begin{aligned} &\mathcal{V}_{(1,1)} \sim \mathcal{V}_{(36,36)} \sim 10^{14}\ \text{GeV} \gtrsim \mathcal{V}_0 \\ &m_X \gtrsim 10^8\ \text{GeV} \end{aligned}\right\} . \tag{30}$$

Here $m_Y^2 \sim m_{Y'}^2 \sim (\Delta m^2)_{|\Delta F|=4} \sim 10^{14}$ GeV^2; thus the $\Delta F = -4$ amplitude has the canonical strength $\sim 10^{-29}$ GeV^{-2}; this would lead to decays of the type $p \rightarrow e^+\pi^0$ etc., with a partial rate $\sim (10^{30}$ years$)^{-1}$. The $\Delta F = -2$ amplitude would have a strength $\lesssim 10^{-29}$ GeV^{-2} for $m_X \gtrsim 10^8$ GeV. In this case <u>$\Delta F = -4$ and $\Delta F = -2$ decay modes of the proton can coexist, but the $\Delta F = 0$ mode would be severely damped</u> *). (If m_X far exceeds 10^8 GeV, $\Delta F = -2$ as well as $\Delta F = 0$ modes and $\Delta F = -4$ modes (i.e. $p \rightarrow e^+\pi^0$ etc.) would be the sole decay modes of the proton.)

$$\text{(B)} \quad \left.\begin{aligned} &\mathcal{V}_0 \sim 10^{14} - 10^{15}\ \text{GeV} \gg \mathcal{V}_{(1,1)} \sim \mathcal{V}_{(36,36)} \sim 10^8\text{-}10^{10}\ \text{GeV} \\ &m_X \sim 10^4\text{-}10^5\ \text{GeV} \end{aligned}\right\} . \tag{31}$$

For this case $\Delta F = -2$ and also the $\Delta F = 0$ amplitudes can each lead to partial decay rates for proton of order $(10^{30}$ years$)^{-1}$, but the $\Delta F = -4$ amplitude would have a strength $\sim 10^{-40}$ GeV^{-2} and thus the corresponding decay modes ($p \rightarrow e^+\pi^0$ etc.) would be severely damped.

$$\text{(C)} \quad \left.\begin{aligned} &\mathcal{V}_0 \sim 10^{13}\ \text{GeV} \gg (\mathcal{V}_{(1,1)}\ \mathcal{V}_{(36,36)})^{1/2} \sim 10^{11}\ \text{GeV} \\ &m_X \sim 10^4\text{-}10^5\ \text{GeV},\ (\delta m)^2_{X\bar{Y}} \sim m_{W_L} m_X \sim 10^6\text{-}10^8\ \text{GeV}^2 \end{aligned}\right\} . \tag{32}$$

This case with $(\delta m^2)_{XY}$ given by $m_{W_L} m_X$ (rather than by $g^2 \mathcal{V}(1,1)\omega \approx g\,\mathcal{V}(1,1) m_{W_L}$) can arise provided the corresponding $X \leftrightarrow \bar{Y}$ mixing arises through the VEV of a multiplet other than **) $\Phi_{\{CD\}}^{\{AB\}}$ with ω (defined by Eq.(27)) set to zero. Here $\Delta F = 0$, -2 as well as -4 would have comparable strengths with partial rates of order $(10^{30}$-10^{32} years$)^{-1}$.

*) Note that in this case, one cannot permit m_X to be much lower than 10^8 GeV, or else the $\Delta F = -2$ amplitude alone would lead to a proton lifetime much shorter than 10^{30} years.

**) As mentioned before, X-$\bar{Y}$ mixing can also arise through VEV of the multiplet $\Phi_{[CD]}^{[AB]}$.

The pattern of symmetry breaking we have so far considered essentially corresponds to the first route involving a descent for SU(16) via SU(8) × SU(8) × $U(1)_F$ (see Fig.2). The relevant steps may be summarized as follows:

$$\text{(I) } SU(16) \xrightarrow{\mathcal{V}_0^{255} \text{ and/or } \mathcal{V}_{(1,1)}} \text{(II) } SU(8)\times SU(8)\times U(1)_F$$

$$\Big\downarrow \mathcal{V}_{(36,36)} \text{ with or without } \{\mathcal{V}_{(1232,1)}, \mathcal{V}_{(1,1232)}, \mathcal{V}_{(63,63)}\}$$

$$\text{(III) } SU(2)_L \times SU(2)_R \times SU(4)_{L+R}$$

$$\Big\downarrow \mathcal{V}_R^{136} \text{ and/or } (u_3, u_{15})$$

$$\text{(IV) } SU(2)_L \times U(1) \times SU(3)^c$$

$$\Big\downarrow \omega,\ldots$$

$$\text{(V) } U(1)_{em} \times SU(3)^c \qquad (33)$$

Needless to say, one single vacuum expectation value parameter can suffice to yield the desired descent at any given stage of SSB exhibited above. For example the pattern

$$\mathcal{V}_0^{255} \gg \mathcal{V}(36,36) \subset \phi^{\{AB\}}_{\{CD\}} \gg \mathcal{V}_R^{136} \gg (\omega,\ldots) \qquad (34)$$

corresponding to just one single scalar multiplet being operative at any one stage of SSB would suffice to provide the descent depicted above.

The SO(10) route

At this point we may mention that an alternative chain, which emphasises the SO(10) subgroup raher than SU(8) × SU(8) can be realized through the dominance of the VEV of the SO(10) singlet in the multiplet $\phi^{\{AB\}}_{\{CD\}}$.

In terms of the SU(8) × SU(8) notation for distinguishing $SU(2)_L \times SU(2)_R \times SU(4)$ singlets, the SO(10) chain would be singled out if the minimum of the effective scalar potential would guarantee the relations (upto perhaps terms of order α or α^2):

$$\mathcal{V}(36,36) = \mathcal{V}(63,63) = 6\mathcal{V}(1232,1) = 6\mathcal{V}(1,1232)$$

$$= \frac{51}{2}\mathcal{V}(1,1) \gg \mathcal{V}_0^{255} \qquad (35)$$

In this case, neither $\mathcal{V}_0$ nor $\mathcal{V}(1,1)$ dominate and as is expected of SO(10), the $\Delta F = -4$ transitions (i.e. $p \to e^+\pi^0$ etc.) are the sole decay modes *) of the proton, the corresponding amplitudes being of order $g^2/m^2_{Y_s}$. Here "Y_s" $\equiv 1/\sqrt{2}\ \{(2,2,6)_2 + (2,2,6)_{-2}\}$ are the gauge particles belonging to SO(10). Note that the special relations (35) between the VEV parameters keep the linear combinations "Y_s" massless, while giving masses to the orthogonal combinations "Y_a" $\equiv 1/\sqrt{2}\ \{(2,2,6)_2 - (2,2,6)_{-2}\}$.

SO(10) can break down to SU(2) × SU(2) × SU(4) at the second stage of spontaneous symmetry breaking via either **) $\mathcal{V}_0^{255}$, or via small departures of order α or α^2 from the special relations (35) or both. The symmetry SU(2) × SU(2) × SU(4) can break down subsequently to SU(2) × U(1) × $SU(3)^c$ as before through $\mathcal{V}_R^{136}$. The second route of spontaneous symmetry breaking may be summarized by the following steps:

$$\text{(II) } SU(16) \xrightarrow[\text{with Eq.(35)}]{\phi^{\{AB\}}_{\{CD\}}} SO(10) \xrightarrow{\mathcal{V}_0^{255}} SU(2)\times SU(2)\times SU(4)^c$$

$$\xrightarrow{\mathcal{V}_R^{136}} SU(2)\times U(1)\times SU(3)^c \xrightarrow{\theta(\omega)} U(1)\times SU(3)^c\,. \qquad (36)$$

The $SU(12)_q \times SU(4)_\ell \times U(1)$ chain and the coexistence of proton decay with n-$\bar{n}$ oscillations

The third alternative route $SU(16) \to SU(12)_q \times SU(4)_\ell \times U(1)$ with $SU(12)_q$ operating on six quarks plus six antiquarks, $SU(4)_\ell$ operating on two leptons plus two antileptons and U(1) gauging $|B| - |L|$, could materialize if the adjoint ***) 255 in the presence of other scalar multiplets acquires a VEV of the form diag A(1,1,1,1,1,1,1,1,1,1,1,1,-3,-3,-3,-3) with A far exceeding the VEV of the

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*) If the grand unification mass is permitted to be as high as Planck mass, then leptoquark X may be permitted to be light enough [23] ($\sim 10^5$ GeV) for SO(10) in accordance with renormalization group equations and this would make $\Delta F = 0$ mode probable. One does not know of course whether perturbative use of renormalization group equations upto Planck mass, where quantum gravity should be important, is meaningful.

**) Note that $\mathcal{V}_0^{255}$ by itself would take SU(16) to SU(8) × SU(8) × U(1). Since SU(8) × SU(8) × U(1) and SO(10) overlap as SU(2) × SU(2) × $SU(4)^c$, the parameter $\mathcal{V}_0^{255}$ and the VEV parameters of $\phi^{\{AB\}}_{\{CD\}}$ together always take SU(16) to SU(2) × SU(2) × $SU(4)^c$.

***) Or alternatively, if the $SU(12)_q \times SU(4)_\ell \times U(1)$ singlet within $\phi^{\{AB\}}_{\{CD\}}$ acquires the largest VEV of all.

other multiplets. This would give superheavy masses ($\gg 10^4$ GeV) to the lepto-quarks X's and the lepto-antiquarks Y''s, but would keep the diquark Y's as well as the dilepton Y" massless. In other words, in this third route, quark-lepton unification is lost at the very first stage of SSB. This is in sharp contrast to the first two routes which retain quark-lepton unification symbolized by $SU(4)_{colour}$ and (thus masslessness of X) till the second or third stages of SSB.

Note also that at this stage the quark number ($B_q = 3B$) as well as the lepton number L are separately conserved and are coupled to <u>massless</u> gauge particles.

At the next stage of SSB, which could arise through the VEV parameters $\mathcal{V}(1,1)$, $\mathcal{V}(36,36)$, etc. belonging to $\Phi^{\{AB\}}_{\{CD\}}$, the symmetry $SU(12) \times SU(4) \times U(1)_{|B|-|L|}$ would descend to $SU(2)_L \times SU(2)_R \times SU(3)^c \times U(1)_{B-L}$. This is the intersection *) of the two symmetries $SU(12)_q \times SU(4)_\ell \times U(1)_{|B|-|L|}$ and $SU(2)_L \times SU(2)_R \times SU(4)^c_{L+R}$. At this stage the diquark Y's acquire masses and simultaneously fermion number violations involving for example Y-$\bar{Y}$' and Y-$\bar{X}$ gauge mixings occurs. The symmetry can reduce further at a third stage through $\mathcal{V}^{136}_R$ and/or u_3, u_{15} (see Eqs.(22) and (23)) to $SU(2)_L \times U(1) \times SU(3)^c$, which can finally reduce to $U(1)_{em} \times SU(3)^c$ via VEV of the type $\omega \sim 10^2$ GeV as mentioned before.

The various stages of SSB corresponding to this third route are summarized below:

$$
\begin{aligned}
&(\mathrm{I})\ SU(16) \xrightarrow{\mathcal{V}^{255}_0 \text{ and/or } "\Phi^{\{AB\}}_{\{CD\}}"} (\mathrm{II})\ SU(12)_q \times SU(4)_\ell \times U(1)_{|B|-|L|} \\
&\qquad\qquad \Big\downarrow\ \mathcal{V}(36,36),\ \mathcal{V}(1,1) \subset \Phi^{\{AB\}}_{\{CD\}} \\
&(\mathrm{III})\ SU(2)_L \times SU(2)_R \times SU(3)^c_{L+R} \times U(1)_{B-L} \\
&\qquad\qquad \Big\downarrow\ \mathcal{V}^{136}_R \text{ and/or } (u_{15}, u_3) \\
&(\mathrm{IV})\ SU(2)_L \times U(1) \times SU(3)^c \\
&\qquad\qquad \Big\downarrow\ \omega, \ldots \\
&(\mathrm{V})\ U(1)_{em} \times SU(3)^c
\end{aligned}
\tag{37}
$$

*) Note that the flavour $SU(2)_{L,R}$ of quarks are disjoint from those of the leptons within $SU(12)_q \times SU(4)_\ell \times U(1)$. The union of the SU(2) of quarks with that of leptons occurs in the present scheme at the second stage of SSB for example through the VEV of $\Phi^{\{AB\}}_{\{CD\}}$.

The striking feature of this route is that the diquark Y as well as the dilepton Y" are relatively light, while Y' and X are superheavy. We now observe that this route can lead to proton decays of the $\Delta F = -4$ variety (i.e. $p \to e^+\pi^0$ etc.) coexisting with n-$\bar{n}$ oscillations of measurable strength.

To construct an effective Lagrangian for n-$\bar{n}$ oscillations, we find it necessary to utilize propagating scalar particles. For simplicity, we restrict the discussion to a <u>136</u> of SU(16). This is the only representation, which can form an SU(16) invariant Yukawa coupling with the 16-plet of fermions. Some of its components develop vacuum expectation values and thereby give rise to the fermion mass terms:

$$
\begin{aligned}
h\ \Psi^T_A C^{-1} \Psi_B \left\langle \bar{\Phi}^{\{AB\}} \right\rangle &= m_u\, u^T_\alpha C^{-1} u^c_{\bar{\alpha}} \\
&+ m_d\, d^T_\alpha C^{-1} d^c_{\bar{\alpha}} + m_e\, e^T C^{-1} e^c + m_\nu\, \nu^T C^{-1} \nu^c \\
&+ \mathcal{M}(\nu_L)\, \nu^T C^{-1} \nu + \mathcal{M}(\nu_R)\, \nu^{cT} C^{-1} \nu^c \ .
\end{aligned}
\tag{38}
$$

Both Dirac (m_ν) and Majorana ($\mathcal{M}(\nu_{L,R})$) mass terms for the neutrino are included for generality. The Dirac mass terms arise from the VEV of the F = 0 components $(2,2,1)_0$ and $(2,2,15)_0$, while the $|\Delta F| = |\Delta L| = 2$ Majorana mass terms $\mathcal{M}(\nu_L)$ and $\mathcal{M}(\nu_R)$ arise respectively from the VEV of the $|F| = 2$ components $(3,1,10)_2$ and $(1,3,\overline{10})_{-2}$, <u>all belonging to the single multiplet 136 of SU(16).</u> The Majorana mass for the left-handed neutrino $\mathcal{M}(\nu_L)$ arising from the VEV of $(3,1,10)_2$, if non-vanishing, must be sufficiently small relative to the other parameters, since VEV of $(3,1,10)_2$ involving a <u>triplet</u> of $SU(2)_L$ leads to departures from unity of the ratio *) $\rho = m_{W_L}/(m_Z \cos\theta_W)$. The Majorana mass of ν_R can however be large ($\gg 100$ GeV) corresponding to a large VEV of $(1,3,\overline{10})_{-2}$ consistent with the suggestion of several authors [24)] in the context of subsymmetries like SO(10) and $SU(2) \times SU(2) \times SU(4)$ and with our discussion in earlier sections regarding the hierarchical pattern of gauge masses. This will keep the left-handed neutrino light with a mass $\sim (1 \text{ to few MeV})^2/\mathcal{M}(\nu_R)$ ~ 1 to 10 eV for $\mathcal{M}(\nu_R) \sim 10^3$ GeV, where (1 to few) MeV represents a typical Dirac mass for the first generation of quarks and leptons. With Majorana masses for neutrinos, neutrinoless double β decay must occur at a level depending upon these masses [4),24)].

*) The present experimental value of $\rho = 1.00 \pm 0.02$.

We now observe that the diquark Y gauge particles can induce, via a loop diagram involving $W_{L,R}\, Y\, \phi' \langle \phi' \rangle$ gauge coupling (see Figs.4(a) and (b)) or through a tree diagram involving an $Y_\mu \partial_\mu \phi' \langle \phi' \rangle$ coupling (Fig.4(c)), <u>effective</u> Yukawa transitions *) satisfying $|\Delta F| = 2$ of the type

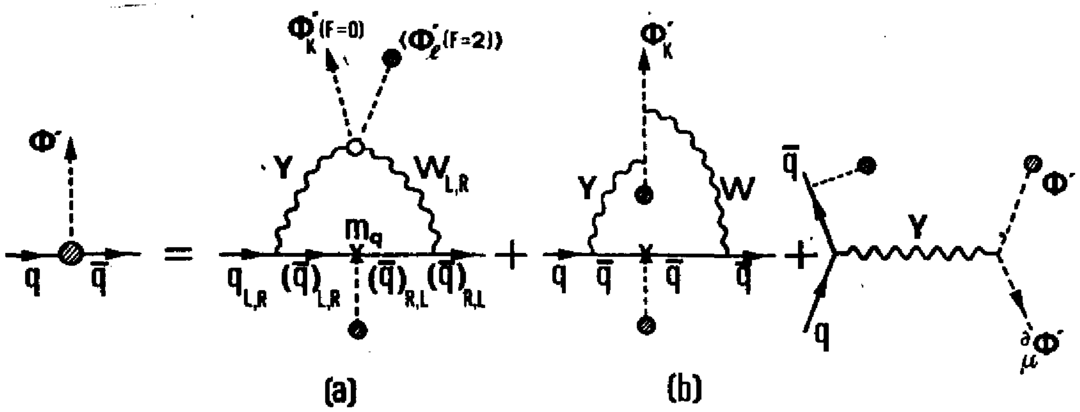
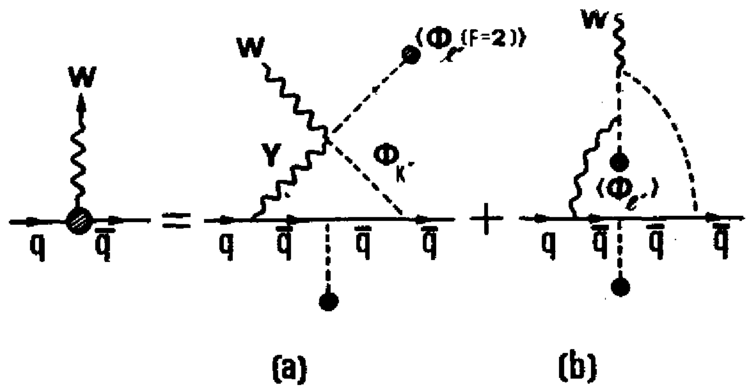


Fig.4: Spontaneously induced $\Delta F = -2$ $q \to \bar{q} + \phi'$ transition. ϕ' belongs to <u>136</u>.



<u>Fig.5</u>: Spontaneously induced $\Delta F = -2$ $q \to \bar{q} + W$ transition.

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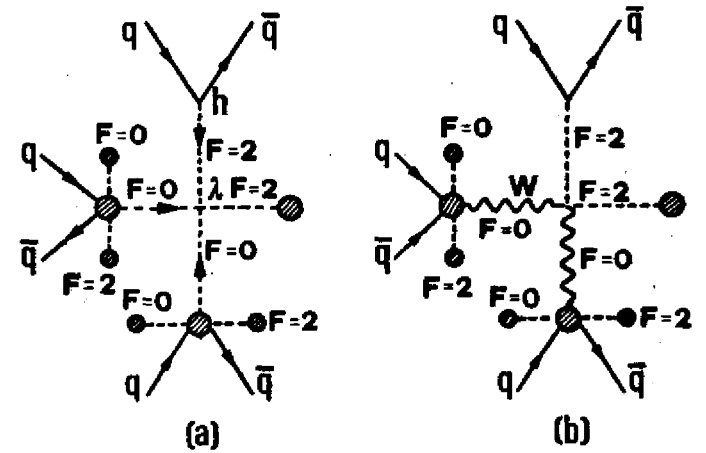
*) Throughout this paper, we have worked with gauge couplings or Higgs self-couplings. No direct Yukawa couplings of fermions with Higgs have been introduced for use in proton decay or n-$\bar{n}$ transition calculations. The gauge couplings of fermions and Higgs particles together with Higgs <u>self-couplings</u> (which in turn determine the Higgs VEV's) then allow us to determine <u>effective</u> Higgs-Yukawa couplings in terms of the gauge parameters and the VEV's . We avoid in this way the arbitrariness of the Higgs-Yukawa coupling, both in respect of their magnitudes as well as in respect of their group-theoretic complexion.

$$(q_i)_{L,R} + (\bar{q}_j)_{R,L} + \phi'_k \,(F = 0) + \langle \phi'_\ell \,(F = 2) \rangle \quad . \qquad (39)$$

Here i,j,k,ℓ take appropriate group indices; ϕ' denotes the scalar multiplet <u>136</u> introduced before; ϕ'_k belongs to the submultiplet $(2,2,1)_0$ or $(2,2,15)_0$ of <u>136</u> possessing $F = 0$, while ϕ'_ℓ belongs to the submultiplet $(3,1,10)_{+2}$ or $(1,3,10)_{+2}$, with $F = +2$, or 136 or 136* carrying $F = 2$. Recall that the VEV of $(1,3,10)_{+2}$ and (possibly also) $(3,1,10)_{+2}$, which are now utilized, are the ones which give Majorana masses to ν_R and ν_L respectively.

The Yukawa transitions listed above are indeed completely analogous to those of the $\Delta F = 0$ $q \to \ell + \phi$ type discussed before (see Fig.3(a)) except that the mediating particle now is the diquark Y rather than the leptoquark X.

The $|\Delta F| = 2$ transitions at the quark level may also arise through effective vertices of the type $q \to \bar{q} + W$ involving the emission of an $F = 0$ gauge meson W (rather than a Higgs ϕ') via loop diagrams shown in Figs.5(a) and (b). In each case the VEV of an appropriate $|F| = 2$ component of ϕ' is utilized.



<u>Fig.6</u>: Mechanisms for $3q \to 3\bar{q}$ transitions. Selection rules force two of the propagating scalar fields in Fig.6(a) to have $F = 0$ and the third to have $F = 2$. The Yukawa vertices denoted by a blob can only arise through a combination of invariant gauge plus Yukawa couplings <u>together</u> with a non-vanishing VEV of an $F = 2$ field.

The effective operator for $n \leftrightarrow \bar{n}$ oscillation consists of several terms of the form uudddd. These can now come about (in a manner completely analogous to $\Delta F = 0$ proton decay, see Fig.3(b)) through third order of Yukawa transitions $q \to \bar{q} + \phi'$ followed by an invariant quartic $\lambda\phi'^4$ coupling, subject to one $F = 2$ component (i.e. $(3,1,10)_{+2}$ or $(1,3,10)_{+2}$) of ϕ' having non-zero VEV. The anticommuting nature of the quark field operators together with fermion number conservation in the basic Lagrangian imply *) that two of the Yukawa transitions must involve the emission of $F = 0$ components (i.e. $(2,2,15)_0$) of ϕ'

*) This follows by noting that the process $3q \to 3\bar{q}$ for the present model (where fermion number is a local symmetry and thus conserved in the gauge Lagrangian) can come about only if we utilize VEV of $F = 2$ components of ϕ' thrice. This says that in Fig.6(a), either (i) all three of these $|\Delta F| = 2$ VEV parameters are utilized in the three Yukawa transitions, each involving $q \to \bar{q} + \phi'$, or (ii) two of these $\Delta F = 2$ VEV parameters are utilized in the Yukawa transitions, while the third $|\Delta F| = 2$ VEV parameter is invoked at the $\lambda(\phi')^4$ vertex. The first case (i) can be eliminated as follows:

Case(i)- Consider the emission of an $F = 0$ component of ϕ' (i.e. either $(2,2,1)_0$ or $(2,2,15)_0$) at all three Yukawa vertices $q \to \bar{q} + \phi'$. To obtain effective uudddd interaction, at least one Yukawa vertex must be of the type $d \to \bar{d} + \phi'$ with a corresponding Lorentz and $SU(3)_{\text{colour}}$ invariant effective interaction $\propto d^T_{L\alpha}(x)\, C^{-1}\, d_{L\beta}(x)\, \bar{\Phi}^{\{AB\}}(x)$; where α, β are SU(2) colour indices running from 1 to 3. Since the d's anticommute, only the antisymmetric combination in the product $(\underset{\sim}{3})_{2,1} \times (\underset{\sim}{3})_{2,1}$ with $\underset{\sim}{3}$ pertaining to SU(3) colour and the subscript (2,1) to Lorentz indices can contribute (i.e. either $(\underset{\sim}{6})_{1,1}$ or $(\underset{\sim}{\bar{3}})_{3,1}$). But for spin-0, only $(\underset{\sim}{6})_{(1,1)}$ is relevant. This however is not contained in either $(2,2,1)_0$ or $(2,2,15)_0$. This eliminates Case (i). The $(\underset{\sim}{6})_{(1,1)}$ is contained, on the other hand, in an $F = 2$ component of ϕ' like $(1,3,10)_2$. Furthermore, observe that two of the Yukawa transitions can be of the type $u \to \bar{d} + \phi'$ for which the symmetric combination $(\bar{3})_{1,1}$, which occurs in the $F = 0$ component $(2,2,15)_0$ can contribute. Thus Case (ii) involving the emission of $F = 0$ components at two of the three Yukawa vertices (i.e. $u \to \bar{d} + \phi'$) and the emission of the $F = 2$ component (like $(1,3,10)_2$) at the third vertex ($d \to \bar{d} + \phi'$) is allowed and uniquely selected out. This uniqueness is a consequence of fermion number being gauged as local symmetry. Note by contrast that if one worked within the subunification symmetry $SU(2)_L \times SU(2)_R \times SU(4)^c_{L+R}$ (Ref.2), which contains B-L but not fermion number as a generator, there would be greater flexibility including the possibility that all four Higgs participating at the $\lambda\phi'^4$ vertex can be (1,3,10) or (3,1,10) particles. Such a vertex utilized in Ref.20 in the basic Lagrangian is simply not permissible within SU(16) invariant structure, as it violates fermion number. For this reason Fig.6(a), while resembling certain features of the mechanism presented in Ref.20, differs in its intrinsic structure.

(each producing a change in fermion number by two units at the vertex concerned, see Eq.(39)), while the third must involve the emission of an $F = 2$ component (i.e. $(1,3,10)_2$ or $(3,1,10)_2$) utilizing the normal SU(16) invariant Yukawa interaction of ϕ'.

An analogous mechanism with two of the propagating $F = 0$ particles being gauge rather than Higgs mesons can also induce $3q \to 3\bar{q}$ transition as shown in Fig.6(b).

Noting that the amplitude for the effective Yukawa transitions $q \to \bar{q} + \phi'$ ($F = 0$) arising via a typical loop diagram (Fig.4(a)) is $\sim \alpha^2 (m_q \langle\phi'(F = 2)\rangle / m_Y^2)\, \ell n(m_Y^2/m_W^2) \equiv h_Y$, the amplitude for $3q \to 3\bar{q}$ (vide Fig.6(a)) is:

$$A(3q \to 3\bar{q}) \sim h\lambda [h_Y]^2 \;\langle\phi'(F = 2)\rangle / [m^2_{\phi'(F=2)}\; m^4_{\phi'(F = 0}] \,. \tag{40}$$

Taking $m_Y \approx 10^4\text{-}10^5$ GeV, $\langle\phi'(F = 2)\rangle \approx 10^4$ GeV, $\lambda \approx \mathcal{O}(\alpha^2)$, $h \approx \mathcal{O}(\alpha)$, $m_{\phi'(F=2)} \approx m_{\phi'(F=0)} \sim \sqrt{\lambda}\, \langle\phi'(F=2)\rangle \sim 10^2$ GeV and $(m_q)_{\text{eff}} \sim \frac{1}{3}$ GeV, we obtain $A(3q \to 3\bar{q}) \sim (10^{-29}\text{-}10^{-33})\ \text{GeV}^{-5}$. This would correspond to an $n\text{-}\bar{n}$ mixing mass $\delta m_{(n\bar{n})} \sim (10^{-29}\text{-}10^{-33})$ GeV, the corresponding $n\text{-}\bar{n}$ oscillation period (for free neutrons) being $\tau_{n\bar{n}} \approx (10^{+5}\text{-}10^{+9})$ sec. with a deuteron lifetime $\approx 10^{30}\text{-}10^{38}$ years. (The $\Delta B = 2$, $\Delta L = 0$ deuteron decay into pions was considered in Ref.4.) Needless to say, the above should be regarded only as a rough estimate of the orders of magnitudes *). Similar estimates come about from Fig.6(b).

The important point to note is that for a mass scale $\approx 10^4\text{-}10^5$ GeV, characterizing the breakdown of $SU(12)_q \times SU(4)_\ell \times U(1)$ to the lower stage (see Eq.(37)) and in particular the mass of the diquark Y, $n\text{-}\bar{n}$ oscillation begins to acquire a strength which should be measurable in the near future[18].

We now face the questions: does the proton decay into leptons and if so does the ultralight mass ($\sim 10^4$ GeV) of the diquark Y make it decay too rapidly? We observe that as $SU(12)_q \times SU(4)_\ell \times U(1)$ descends to $SU(2)_L \times SU(2)_R \times SU(3)^c \times U(1)_{B-L}$ through the VEV parameters $\mathcal{V}_{(36,36)}$ and $\mathcal{V}_{(1,1)} \subset \Phi^{\{AB\}}_{\{CD\}}$ (see Eq.(13)). The diquark Y acquires a mass; simultaneously the Y mixes with the lepto-antiquark $\bar{Y}'$ with a mixing $(\text{mass})^2$ $(\delta m)^2_{YY'} = g^2 \mathcal{V}(1,1)\, \mathcal{V}(36,36) \approx m_Y^2$ (see Eq.(18)).

*) In particular one might consider $\langle\phi'(F=2)\rangle \sim 10^3$ GeV and therefore a relatively light W_R of about 10^3 GeV. Allowing $\lambda \sim \mathcal{O}(\alpha)$, $m(\phi')$ is still $\sim 10^2$ GeV; such a mass will not much alter the estimate above.

This has the important consequences that (i) the $\Delta F = -4$ proton decay ($p \to e^+\pi^0$ etc.) must occur as in Fig.1(a) and (ii) the amplitude for this decay (see Eq.(29)) being $g^2(\delta m)^2_{YY'}/[m_Y^2 m_{Y'}^2]$ is just $(g^2/m_{Y'}^2)$. Thus with $m_{Y'} \approx 10^{14}$ GeV (a feature consistent with the fact that the lepto-antiquark Y' acquires mass at the first stage of spontaneous breakdown of SU(16) into $SU(12)_q \times SU(4)_\ell \times U(1)$), the $\Delta F = -4$ proton decay has the canonical strength $\approx 10^{-29}\ \mathrm{GeV}^{-2}$. In other words, the proton will decay (via $\Delta F = -4$, $\Delta(B-L) = 0$ mode), but not unduly rapidly, <u>despite</u> the ultralight diquark Y ($m_Y \sim 10^4$ GeV). We stress that this is something which could be possible only because of the spontaneous nature of B, L, F violations occurring within our approach (for which Y-$\bar{Y}'$ mixing vanishes as $m_Y \to 0$).

In this third route of spontaneous descent, not only lepto-antiquark Y' but also the leptoquark X acquires a mass at the first stage of SSB. With $m_{Y'} \gtrsim 10^{14}$ GeV (see above) it follows that $m_X \approx m_{Y'} \gtrsim 10^{14}$ GeV. Thus for this third route, it is clear (see discussion presented before) that $\Delta F = 0$ proton decay (i.e. $p \to 3$ leptons + mesons) would be strongly suppressed. The $\Delta F = -2$ mode, with amplitude $\propto g^2(m_{W_L} m_Y)/(m_X^2 m_Y^2) \sim g^2(m_{W_L}/m_Y)(1/m_X^2)$ would also be suppressed relative to the $\Delta F = -4$ mode. To summarize, the third route of spontaneous descent with a mass hierarchy of the type summarized below

$$m_X \approx m_{Y'} \approx 10^{14}\ \mathrm{GeV},\ m_Y \approx m_{Y''} \approx (\delta m)_{YY'} \approx 10^4\text{-}10^5\ \mathrm{GeV} \qquad (41)$$

<u>gives rise to proton decays of the $\Delta F = -4$ variety, i.e. $p \to e^+\pi^0$ etc. coexisting with n-$\bar{n}$ oscillations of measurable strength even though $\Delta F = 0$, -2 and -6 proton decay modes would be suppressed.</u>

We regard this as one of the most interesting features of physics based on SU(16). If proton decay of the $\Delta F = -4$ variety as well as n-$\bar{n}$ oscillations are indeed found to occur, we would interpret this as indicative of a maximal gauging as in SU(16) with a relatively light diquark Y ($m_Y \sim 10^4$-10^5 GeV), together with a superheavy leptoquark X ($m_X \approx 10^{14}$ GeV). Contrast this with the model of Ref.20, based on the subunification symmetry [2)] $SU(2) \times SU(2) \times SU(4)$, in which n-$\bar{n}$ oscillation occurs but without proton decay and it is the leptoquark X, which is relatively light ($m_X \sim 10^4$-10^6 GeV). We trace this distinction primarily to the fact that the subunification symmetry does not contain fermion number as a local gauge symmetry, while SU(16) does.

Finally, we wish to note that the $\Delta F = -6$ mode (i.e. $p \to 3\bar{\ell}$ + mesons), though not explicitly discussed so far, can also be prominent. Just as the $\Delta F = 0$ mode ($p \to 3\ell$ + mesons) is mediated by the leptoquark X coupled to $\bar{q}\gamma_\mu \ell$ current (see Figs.3(a) and (b)) likewise the $\Delta F = -6$ mode ($p \to 3\bar{\ell}$ + mesons)

can be mediated by the lepto-antiquark Y' coupled to $\bar{q}\gamma_\mu \ell^c$ current and would be prominent if $m_{Y'} \approx 10^4$-10^5 GeV. The three routes of spontaneous descent exhibited in Sec.I do not however permit the lepto-antiquark to be this light and thus suppress the $\Delta F = -6$ mode. One amusing possibility nevertheless is that SU(16) may descend via a fourth alternative route utilizing the VEV of <u>255</u> into $SU(8)'_I \times SU(8)'_{II} \times U(1)$, where $SU(8)'_I$ operates on six quarks plus two antileptons (rather than leptons) and $SU(8)'_{II}$ operates on six antiquarks plus two leptons (rather than antileptons); the symmetry $SU(8)'_I \times SU(8)'_{II} \times U(1)$ descends subsequently to the <u>familiar</u> lower symmetry $SU(2)_L \times SU(2)_R \times SU(3)^c_L \times U(1)_{B-L}$ and still further to $SU(2)_L \times U(1) \times SU(3)_C$ utilizing the VEV of $\phi^{\{AB\}}_{\{CD\}}$ and <u>136</u> respectively. In this way the Y' can be light ($\sim 10^4$-10^5 GeV) if $\langle \phi^{\{AB\}}_{\{CD\}} \rangle \sim 10^4$-$10^5$ GeV. The diquark Y would be heavy ($\gg 10^4$-10^5 GeV), as for route I, but simultaneously the leptoquark X' would be heavy. Thus the $\Delta F = -6$ mode ($p \to e^+ + \bar{\nu} + \bar{\nu}$ etc) can coexist with the $\Delta F = -4$ mode *) ($p \to e^+$ + mesons) but the $\Delta F = 0$ ($p \to 3$ leptons + mesons) and $\Delta F = -2$ ($p \to e^-$ + mesons) modes (requiring relatively light X) as well as $\Delta B = 2$ n-$\bar{n}$ oscillation (requiring relatively light Y) would be suppressed.

The features of the four alternative routes of spontaneous descent are summarized at the end in Table I.

The different complexions for proton decay and $\Delta B = 2$ n-$\bar{n}$ oscillation if found to occur with measurable strengths for the present level of experimentation would symbolize different characteristic mass scales (M_c) for the hierarchy of grand unification. These are listed below **)

*) The $\Delta F = -4$ mode would be induced by Y-$\bar{Y}'$ mixing as in Fig.1(a), the mixing $(\mathrm{mass})^2$ is bounded above in the present case by the lighter mass scale $\sim m_{Y'} \sim 10^4$-10^5 GeV.

**) These mass scales derived through explicit mechanisms coincide approximately with the estimates of Ref.25 which are based on simplifying operator and dimensional analysis. Some differences (as in the case of $\Delta F = -4$ permitting $M_{Y,Y'} \ll 10^{14}$ GeV with $\Delta_{YY'} \ll m_{Y,Y'}$ for the present case versus $M_c \sim 10^{14}$ GeV of Ref.26) have their origin in spontaneous rather than explicit violations of the B, L and F. We stress that the characteristic mass scales exhibited in Eq.(42) are derived under the assumption that quarks and leptons are elementary. These estimates may change radically for a preonic or pre-preonic composite model of quarks, leptons, gauge and Higgs particles.

$$\Delta F = -4 \ (p \to e^+\pi^0 \text{ etc}) \cdots\cdots \to \quad M_c \sim 10^{14} \text{ GeV (\underline{if} } m_Y \approx m_{Y'} \approx \Delta_{YY'} \text{)}$$

$$\Delta F = -2 \ (p \to e^- + \text{mesons}) \cdots\cdots \to \quad M_c \sim 10^{8}\text{-}10^{12} \text{ GeV (see text)}$$

$$\Delta F = 0 \ (p \to 3 \text{ leptons} + \text{mesons}) \cdots\cdots \to \quad M_c \sim 10^{4}\text{-}10^{5} \text{ GeV}$$

$$\Delta F = -6 \ (p \to 3 \text{ antileptons} + \text{mesons}) \cdots \to \quad M_c \sim 10^{4}\text{-}10^{5} \text{ GeV}$$

$$\Delta F = 2 \ (n \leftrightarrow \bar{n} \text{ oscillation}) \cdots\cdots \to \quad M_c \sim 10^{4\text{-}5} \text{ GeV} \quad . \tag{42}$$

We now face two important questions:

__Question 1:__ Can the vastly different mass scales of the type exhibited in Eqs.(30)-(32) and (41) and Table I coexist in accordance with the constraints of renormalization group equations for the running coupling constants and the observed values of $\sin^2\theta$ and α_s?

It has been shown elsewhere [6),7)] that the answer to the above question is in the affirmative. In particular, the gauge mass patterns of the type exhibited in Eqs.(30), (31) or (32) and therefore the coexistence of some or all of the proton decay modes satisfying $\Delta F = 0, -2$ and -4 are indeed permissible if we allow distinctions between the three known families to appear at the basic level by gauging their maximal symmetries like $[SU(16)]^3$ or $SU(48)$. The major consequence of such maximal symmetries from an experimental point of view is that they fill the much advocated "desert" between 10^2 and 10^{15} GeV with intermediate mass scales and thus new physics [*)] appears at these energy scales.

Finally, the realization of n-$\bar{n}$ oscillations together with the $\Delta F = -4$ proton decay ($p \to e^+\pi^0$ etc.) at an observable level requires a mass pattern $m_{Y'} \sim m_X \sim 10^{14}$ GeV and $m_Y \approx \Delta_{YY'} \approx 10^4$-$10^5$ GeV (Eq.(41)). This also appears to be permissible for the three family symmetry $[SU(16)]^3$ descending in a

*) One example of such new physics is the breakdown of (e,μ,τ) family universality. Such a breakdown would arise from contributions of gauge particles belonging to $[SU(16)]^3/SU(16)_{e+\mu+\tau}$, which couple to __differences__ of electronic, muonic and τ currents. These contributions can be prominent at an energy scale which need be no higher than 10^4-10^5 GeV (see Refs.6 and 7).

manner analogous [*)] to $SU(16) \to SU(12)_q \times SU(4)_\ell \times U(1)_{|B|-|L|}$.

__Question 2:__ Does the coexistence of the alternative proton decay modes ($\Delta F = -4, -2, 0, -6$) with or without $n \leftrightarrow \bar{n}$ oscillation (symbolizing diverse characteristic mass scales) conflict with the cosmological generation of baryon excess [27)] for a Universe, assumed initially to be matter-antimatter symmetric?

Weinberg [25)] has made the interesting observation that the answer to this question is in the affirmative, if violations of B, L and F are explicit. The argument is briefly this: since the $\Delta F = 0, -2$ and -6 proton decays as also $\Delta B = 2$ n-$\bar{n}$ oscillation are mediated by intermediate mass scales M_I ($\sim 10^4$-10^{11} GeV) $\ll M \sim 10^{14}$ GeV, such processes would have relatively fast rates $\sim \alpha T$ or $\alpha^2 T$ in the early Universe at temperatures in the range $M_I \ll T \ll M$, where the gauge particle masses vanish. Thus these processes (with rates exceeding the rate of expansion of the Universe) would be in thermodynamic equilibrium and would wipe out any baryon excess generated in earlier epochs at temperatures $\sim M$ due to the $\Delta F = -4$ process, unless a specific linear combination $B + aL$ is conserved ($a \neq 0$). This would imply that only one of the three modes $\Delta F = 0, -2$ __or__ -6 can coexist with the $\Delta F = -4$ mode - not more than one - and furthermore there should be no $\Delta B = 2$ n-$\bar{n}$ oscillation.

*) For $[SU(16)]^p$, where $p = 1,2,3$ denotes one, two, three families, the relevant descents may be as follows:

$$\text{(i)} \quad [SU(16)]^p \xrightarrow{M_1} SU(16)_{e+\mu+\tau} \to SU(12) \times SU(4) \times U(1)$$

or

$$\text{(ii)} \quad [SU(16)]^p \xrightarrow{M_1} SU(12)^p \times SU(4)^p \times U(1)^p \to SU(12) \times SU(4) \times U(1)$$

or

$$\text{(iii)} \quad [SU(16)]^p \xrightarrow{M_1} SU(12) \times SU(4)^p \times U(1)^p \to SU(12) \times SU(4) \times U(1)$$

$$\text{(iv)} \quad [SU(16)]^p \to [SU(6)_q]^p_L \times [SU(6)_q]^p_R \times [SU(4)_\ell]^p \times U(1) \ .$$

Descents via (iii) or (iv) lead to desirable solutions for gauge masses permitting coexistence of proton decay with n-$\bar{n}$ oscillations of measurable strength. The renormalization group equations for these cases will be presented in a separate note.

It has been remarked [7] that these arguments apply, however, only if the violations of B, L and F are explicit rather than spontaneous. For the latter case the violations ($\Delta F = 0, -2, -6$ or $\Delta B = 2$) associated with a mass scale M_I disappear for temperatures $> M_I$. These violations appear only at temperatures $T \lesssim M_I$. However, the associated gauge particles acquire their masses at the same time that the violations appear. The rates of B, L and F violating processes ($\Delta F = 0, -2, -6$, etc.) damped by the associated gauge masses are now lower than the expansion rate of the Universe at this epoch. This keeps all these processes out of equilibrium whenever they are operative. Consequently, baryon excess generated in an earlier epoch $T > M_I$ by for example the $\Delta F = -4$ process is not wiped out by processes occurring at later times. We conclude that there is no conflict between the coexistence of $\Delta F = 0, -2, -4$ and -6 proton decay modes (with even $\Delta B = 2$ n-$\bar{n}$ oscillation) and the generation of the baryon excess, if the violations of B, L, F are spontaneous. This is one of the crucial differences between explicit versus spontaneous violations of B, L, F.

IV. CONCLUDING REMARKS

(1) Violation of baryon, lepton and, in general, fermion number is central to the hypothesis of quark lepton-unification in a gauge context. Three of its characteristic signatures are proton decay, n-$\bar{n}$ oscillation and neutrinoless double β decay. In 1974 and 1975 it was already shown [4] that within maximal gauging, the proton may decay via alternative modes satisfying $\Delta F = 0, -2, -4$ and -6 some of which may coexist; the deuteron may decay into pions and neutrinoless double β decay may occur in the context of spontaneous gauge symmetry breaking. Here we have shown that n-$\bar{n}$ oscillations (which are related to deuteron decays into pions) can coexist with proton decay of especially $\Delta F = -4$ variety ($p \to e^+ \pi^0$) and both these processes may possess measurable strength so as to be amenable to forthcoming searches. This paper has sharpened the formalism, given the details of Higgs breakings and also shown that the existence of alternative decay modes of the proton pose no conflict with the cosmological generation of baryon excess.

(2) Non-maximal symmetries like SO(10), $SU(8) \times SU(8) \times U(1)_F$ and SU(5) are contained within SU(16) and arise effectively as special cases through spontaneous descent of the maximal symmetry SU(16). Their predictions on proton decay are therefore naturally contained within the predictions of SU(16).

(3) It is a reasonable expectation [21),26)] within most models that the lifetime of the proton will lie within the range of 10^{28}-10^{33} years, making forthcoming searches fully sensitive to all these models.

(4) Neither proton decay into the modes enumerator in this paper nor n-$\bar{n}$ oscillation are tied in any way to the nature of quark charges.

(5) Quark-lepton unification symmetry gauged in its maximal form, permits in general several intermediate mass scales [6),7)] ($\sim 10^4$ to 10^{12} GeV) filling the grand plateau between 10^2 and 10^{15} GeV. Proton decay and n-$\bar{n}$ oscillation can provide a window to these intermediate scales. In particular, the observation of the $\Delta F = 0$, or -6 mode (i.e. $p \to 3\ell$ or $3\bar{\ell}$ + mesons) and/or n-$\bar{n}$ oscillation at any level within conceivable future will strongly suggest the existence of new physics at 10 to 100 TeV region and thereby motivate building high energy machines in this range. For this reason, second and third generation experiments for proton decay and n-$\bar{n}$ oscillation must be planned to look for all modes listed above as possible rare processes, in case they are not found in the first generation experiments.

(6) Coexistence of any two of the proton decay modes $\Delta F = 0, -2$ or -6 or the mere existence of $\Delta B = 2$ n-$\bar{n}$ oscillation with proton decay of the $\Delta F = -4$ variety would strongly suggest that the associated violations of B, L, F are spontaneous rather than explicit. This is in order that cosmological generation of baryon excess [27)] may survive.

(7) Observation of proton decay will strongly support the idea that quark matter and leptonic matter are similar [2)] in their composition, though this has no bearing on the question of whether quarks and leptons represent the ultimate constituents of matter.

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TABLE I

A Sample of Alternative Selection Rules

Routes of SSB (see Fig. 2 and text)	Patterns of VEV	Selection Rules
(I) $SU(16) \to [SU(8)_I \times SU(8)_{II}] \times U(1)_F$	(A) $v_{(1,1)} \sim v_{(36,36)} \sim 10^{14} \gtrsim v_0^{255}$ $m_X \gtrsim 10^8$	$\Delta F = -4, -2$ Yes $\Delta F = 0, -6$, No $\Delta B = 2$ $n-\bar{n}$ No
	(B) $v_0^{255} \sim 10^{14} \gg v_{(1,1)} \sim v_{(36,36)} \sim 10^8 - 10^{10}$ $m_X \sim 10^4 - 10^5$	$\Delta F = 0, -2$ Yes $\Delta F = -4, -6$ No $\Delta B = 2$ $n-\bar{n}$ No
	(C) $v_0^{255} \sim 10^{13} \gg [v_{(1,1)} v_{(36,36)}]^{1/2} \sim 10^{11}$ $m_X \sim 10^4 - 10^5, (\delta m)^2_{XY} \sim m_{W_L} m_X \sim 10^6 - 10^7$	$\Delta F = 0,-2,-4$ Yes $\Delta F = -6$ No $\Delta B = 2$ $n-\bar{n}$ No
(II) $SU(16) \to SO(10)$	VEV of $\phi^{\{AB\}}_{\{CD\}} \gg v_0^{255}$ subject to the special relations eq. (35)	$\Delta F = -4$ Yes $\Delta F = -2,-6$ No $\Delta F = 0$ No (see text) $\Delta B = 2$ $n-\bar{n}$ No
(III) $SU(16) \to SU(12)_q \times SU(4)_\ell \times U(1)_{\lvert B\rvert - \lvert L\rvert}$	$v_0^{255} \sim 10^{14}\,\mathrm{GeV} \sim m_{X,Y'}$ $v_{(36,36)} \sim v_{(1,1)} \sim 10^4 - 10^5 \sim m_{Y,Y''} \sim (\delta m)_{YY'}$	$\Delta B = 2$ $n-\bar{n}$ Yes $\Delta F = -4$ Yes $\Delta F = 0,-2,-6$ No
(IV) $SU(16) \to SU(8)'_I \times SU(8)'_{II} \times U(1)$	$v_0^{255} \sim 10^{14}$ $m_{Y,X,Y''}$ $v_{(1,1)} \sim v_{(36,36)} \sim 10^4 - 10^5 \sim m_{Y'}$	$\Delta F = -4,-6$ Yes $\Delta F = 0,-2$ No $\Delta B = 2$ $n-\bar{n}$ No

Table I: $\Delta F = 0, -2, -4$ and -6 proton-decays correspond to $p \to 3$ leptons + mesons, $p \to$ lepton (e^- or ν) + mesons, $p \to$ antilepton (e^+ or $\bar{\nu}$) + mesons and $p \to 3$ antileptons + mesons respectively. The VEV-parameters and masses are given in units of GeV. For explanations of the alternative patterns of VEV and the definitions of VEV-parameters see text. For a definition of the different gauge particles X, Y, Y' and Y" see Eq.(8). The characterization "yes" for any proton-decay mode corresponds to a partial rate $\sim (10^{29}-10^{33}$ years$)^{-1}$, while that for $n-\bar{n}$ oscillation corresponds to free neutron oscillation period $\sim 10^{+5} = 10^{+9}$ sec. with "no" implying that the corresponding modes have lower rates.

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